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Electroweak and Strong Penguins in $B^{\pm,0} \to \pi\pi, \pi K$ and KK Decays

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Abstract

We calculate CP-violating rates and asymmetry parameters in charged and neutral $B \to \pi\pi, \pi K$ and $\bar{K}K$ decays arising from the interference of tree and penguin (strong and electroweak) amplitudes with different strong and CKM phases. The perturbative strong (electroweak) phases develop at order α_s (α_{em}) from absorptive parts of one-loop matrix elements of the next-to-leading (leading) logarithm corrected effective Hamiltonian. The BSW model is used to estimate the hadronic matrix elements. Based on this model, we find that the effect of strong phases and penguins is substantial in most channels, drastic in many. However, a measurement of the time dependence parameter $a_{\epsilon+\epsilon'}$ in the $\pi^+\pi^-$ channel is only influenced at the 20% level by the complication of the penguins. Recent flavor sum rules developed for $B^{0,\pm} \to \pi\pi, \pi K, K\bar{K}$ amplitudes are tested in this model. Some are well satisfied, others badly violated, when electroweak penguins are included.

I. INTRODUCTION

So far CP violation [1] has been detected only in processes related to $K^0 - \bar{K}^0$ mixing [2] but considerable efforts are being made to find it in B decays. Partial rate and time dependence asymmetries are the leading signals if the CKM [3] model of CP violation is the correct guide. The next generation of B experiments will then make detailed probes of rates and time dependence from which, in principle, a definitive test of the CKM scheme can be made in a model independent way. This amounts to testing whether the unitarity triangle closes, that is, whether three generations and a single phase suffice to describe CP violation in the K and B systems.

CP violation in the CKM model is either 'direct,' in which two amplitudes for the same process have different weak and strong phases, or 'indirect,' in which one of the two interfering amplitudes proceeds through the mixing of the neutral B and $\bar{B}$ mesons. Rate asymmetries between charge conjugate $B^{\pm}$ exclusive channels are purely direct. Their advantage is that they do not require complicated time dependence measurements, involving tagging, and they are definitive signals of CP violation if seen. Their disadvantage is that a model is needed to calculate the strong phase if CKM phase information is to be extracted. Typically these asymmetries arise when tree and penguin amplitudes interfere. A further complication is that direct and indirect CP violation occur simultaneously in neutral decays when there are penguin as well as tree amplitudes in addition to mixing. A recent treatment of this complication can be found in reference [4].

Various authors have shown how to extract CKM phases from certain sets of measurements in a fairly model independent way. Gronau and London [5] have shown that measurements of all charge states in $B \to \pi\pi$, as well as the time dependence of $B \to \pi^+\pi^-$ are required to obtain the CKM phases. Nir and Quinn [6] have extended this analysis to the $K\pi$ system. Finally, Gronau et al. [7] have argued that by using SU(3) relations, there is enough information from rates alone to determine the CKM phases without measurements of time dependence. Unfortunately this analysis relied heavily on the assumption that elec-

troweak penguins were negligible. Subsequently Deshpande and He [8, 9] showed that this analysis which relied on SU(3) symmetry and the particular structure of the strong penguins ceases to be true when electroweak penguins are included. Recently the authors of reference [7] have modified and extended their approach to include electroweak (EW) penguins and B_s transitions [10].

In practice a full determination of all rates and time dependence is a formidable experimental task. Thus in the first stages input from models will be useful to guide the analysis to the measurements which are most easily made and which can give the best early answers. This requires a consistent modelling of tree and penguin amplitudes including strong phases. The strong phase is generated by final state interactions. At the quark level the strong interaction effects can be modeled perturbatively, following Bander, Silverman and Soni [11], by the absorptive part of penguin diagrams.

Recently we have developed a model for the strong phases of the penguins and applied it to rate asymmetries in charged $B \to PP, PV, VV$ channels [12, 13]. In order to systematically take into account the $O(\alpha_s)$ penguin matrix elements, we base our treatment on the next-to-leading logarithmic short distance corrections evaluated by Buras et al. [14]. In this work we have also included the $O(\alpha_{em})$ electroweak penguins in order to investigate their effect on rates, asymmetries, and the SU(3) sum rules.

Having modeled the tree and penguin operators, we use factorization and the BSW current matrix elements to calculate rates and asymmetries. In this note we will apply this model to neutral and charged $\pi\pi, \pi K, KK$ channels. Of course, in order to make definitive predictions we need the CKM parameters as input. A recent study of the experimental data constraining these parameters is the work of Ali and London [15] which shows a preferred solution of $\rho = -0.12, \eta = 0.34$ in the Wolfenstein representation. We use this solution as an illustrative example of how significant strong and electroweak penguins can be in this model. Of course other still possible values for ρ and η will lead to somewhat different predictions.

The remainder of this paper is organized as follows. In Sect. 2 we describe the effective weak Hamiltonian and the evaluation of the hadronic matrix elements. In Sect. 3 we give a

short account of ref. [4] needed for our calculation. The final results for the branching ratios and rate differences are discussed in Sect. 4. In this section we also analyse the SU(3) sum rules. In Sect. 5 we end with a short summary and draw some conclusions.

II. THE EFFECTIVE HAMILTONIAN

In the next two subsections we present the short distance Hamiltonian and the quark-level matrix elements. In subsection 2.3 we describe the evaluation of the hadronic matrix elements which are relevant for the PP final states.

A. Short Distance QCD Corrections

For calculations of CP-violating observables it is most convenient to exploit the unitarity of the CKM matrix and split the effective weak Hamiltonian into two pieces, one proportional to $v_u \equiv V_{ub}V_{us}^*$ (or $V_{ub}V_{ud}^*$ in the case of $b \to d$ transitions) and the other one proportional to $v_c \equiv V_{cb}V_{cs}^*$ (or $V_{cb}V_{cd}^*$ correspondingly),

$$\mathcal{H}_{\rm eff} = 4\frac{G_F}{\sqrt{2}}\left(v_u\mathcal{H}_{\rm eff}^{(u)} + v_c\mathcal{H}_{\rm eff}^{(c)}\right) \ . \tag{1}$$

The two terms ($q = u, c$)

$$\mathcal{H}_{\rm eff}^{(q)} = \sum_i c_i(\mu)\cdot O_i^{(q)} \ , \tag{2}$$

differ only by the quark content of the local operators, and for our purposes it is sufficient to consider only the following four-quark operators:

$$\begin{aligned}
O_1^{(q)} &= \bar{s}_\alpha\gamma^\mu Lq_\beta\cdot\bar{q}_\beta\gamma_\mu Lb_\alpha \ , & O_2^{(q)} &= \bar{s}_\alpha\gamma^\mu Lq_\alpha\cdot\bar{q}_\beta\gamma_\mu Lb_\beta \ , \\
O_3 &= \bar{s}_\alpha\gamma^\mu Lb_\alpha\cdot\sum_{q'}\bar{q}'_\beta\gamma_\mu Lq'_\beta \ , & O_4 &= \bar{s}_\alpha\gamma^\mu Lb_\beta\cdot\sum_{q'}\bar{q}'_\beta\gamma_\mu Lq'_\alpha \ , \\
O_5 &= \bar{s}_\alpha\gamma^\mu Lb_\alpha\cdot\sum_{q'}\bar{q}'_\beta\gamma_\mu Rq'_\beta \ , & O_6 &= \bar{s}_\alpha\gamma^\mu Lb_\beta\cdot\sum_{q'}\bar{q}'_\beta\gamma_\mu Rq'_\alpha \ , \\
O_7 &= \bar{s}_\alpha\gamma^\mu Lb_\alpha\cdot\sum_{q'}e_{q'}\bar{q}'_\beta\gamma_\mu Rq'_\beta \ , & O_8 &= \bar{s}_\alpha\gamma^\mu Lb_\beta\cdot\sum_{q'}e_{q'}\bar{q}'_\beta\gamma_\mu Rq'_\alpha \ , \\
O_9 &= \bar{s}_\alpha\gamma^\mu Lb_\alpha\cdot\sum_{q'}e_{q'}\bar{q}'_\beta\gamma_\mu Lq'_\beta \ , & O_{10} &= \bar{s}_\alpha\gamma^\mu Lb_\beta\cdot\sum_{q'}e_{q'}\bar{q}'_\beta\gamma_\mu Lq'_\alpha \ .
\end{aligned} \tag{3}$$

where L and R are the left- and right-handed projection operators. The q' run over the quark operators that are active at the scale $\mu = O(m_b), (q' \epsilon \{u, d, s, c, b\})$. The operators $O_3, \ldots, O_6$ arise from (QCD) penguin diagrams which contribute at order α_s to the initial values of the coefficients at $\mu \approx M_W$[14] or through operator mixing during the renormalization group summation of short distance QCD corrections [16]. The usual tree-level W-exchange corresponds to O_2 (with $c_2(M_W) = 1 + O(\alpha_s)$). Similarly, $O_7, \ldots, O_{10}$ arise from electroweak penguin diagrams.

The renormalization group evolution from $\mu \approx M_W$ to $\mu \approx m_b$ has been evaluated in leading logarithmic (LL) order in the electromagnetic coupling and in next-to-leading logarithmic (NLL) precision in the strong coupling α_s. The full 10x10 matrices for the anomalous dimensions γ_i are given by Buras et al. in references [17, 18, 19]. The Wilson coefficients c_i obtained depend on the renormalization scheme used. This renormalization scheme dependence can be isolated in terms of the matrices $\mathbf{r}_{ij}^{s,ew}$ by writing [14, 19, 20]

$$c_j(\mu) = \sum_i \bar{c}_i(\mu) \left[\delta_{ij} - \frac{\alpha_s(\mu)}{4\pi} \mathbf{r}_{ij}^{s} - \frac{\alpha_{em}(\mu)}{4\pi} \mathbf{r}_{ij}^{ew} \right] \ , \tag{4}$$

where the coefficients $\bar{c}_j$ are scheme independent at this order. The matrix elements $\mathbf{r}_{ij}^{s,ew}$ have been evaluated in references [20, 14, 19, 8]. To obtain numerical values for the $\bar{c}_i$ we must specify the input. We choose $\alpha_s(M_z) = 0.118$, $\alpha_{em}(M_z) = 1/128$ and $\mu = m_b = 4.8\ GeV$. Then we have [14, 9]

$$\begin{aligned}
\bar{c}_1 &= -0.324 \ , & \bar{c}_2 &= 1.15 \ , \\
\bar{c}_3 &= 0.017 \ , & \bar{c}_4 &= -0.038 \ , \\
\bar{c}_5 &= 0.011 \ , & \bar{c}_6 &= -0.047 \ , \\
\bar{c}_7 &= -1.05 \times 10^{-5} \ , & \bar{c}_8 &= -3.84 \times 10^{-4} \ , \\
\bar{c}_9 &= -0.0101 \ , & \bar{c}_{10} &= 1.96 \times 10^{-3} \ ,
\end{aligned} \tag{5}$$

Other values using slightly different input can be found in [20, 14]. The electroweak coefficient $\bar{c}_9$, as noted in reference [9], is not much smaller than the strong penguins; its major contribution arises from the Z penguin.

B. Quark-level Matrix Elements

Working consistently at NLL precision, the matrix elements of $\mathcal{H}_{\rm eff}$ are to be treated at the one-loop level in order to cancel the scheme dependence from the renormalization group evolution. The one-loop matrix elements can be rewritten in terms of the tree-level matrix elements of the effective operators [20, 21, 12, 13]

$$\langle sq'\bar{q}'|\mathcal{H}_{\rm eff}^{(q)}|b\rangle = \sum_{i,j} c_i(\mu)\left[\delta_{ij} + \frac{\alpha_s(\mu)}{4\pi}\mathbf{m}_{ij}^{s}(\mu,\ldots) + \frac{\alpha_{em}(\mu)}{4\pi}\mathbf{m}_{ij}^{ew}(\mu,\ldots)\right]\langle sq'\bar{q}'|O_j^{(q)}|b\rangle^{\rm tree} \ . \quad (6)$$

The functions $\mathbf{m}_{ij}^{s,ew}$ are determined by the corresponding renormalized one-loop diagrams and depend in general on the scale μ, on the quark masses and momenta, *and* on the renormalization scheme. The various one-loop diagrams can be grouped into two classes: *vertex-corrections*, where a gluon connects two of the outgoing quark lines, and *penguin* diagrams, where a quark-antiquark line closes a loop and emits a gluon, which itself decays finally into a quark-antiquark pair. The full matrices $\mathbf{m}_{ij}^{s,ew}$ have been worked out by Fleischer [20, 21] and Deshpande and He [8].

When expressing the rhs of (6) in terms of the renormalization scheme independent coefficients $\bar{c}_i$, the effective coefficients multiplying the matrix elements $\langle sq'\bar{q}'|O_j^{(q)}|b\rangle^{\rm tree}$ become

$$c_j^{\rm eff} \equiv \bar{c}_j + \sum_i \bar{c}_i \cdot [\frac{\alpha_s}{4\pi}\left(\mathbf{m}_{ij}^{s} - \mathbf{r}_{ij}^{s}\right) + \frac{\alpha_{em}}{4\pi}\left(\mathbf{m}_{ij}^{ew} - \mathbf{r}_{ij}^{ew}\right)] \ . \quad (7)$$

The renormalization scheme dependence, which is present in $\mathbf{m}_{ij}$ and $\mathbf{r}_{ij}$, explicitly cancels in the combinations $\mathbf{m}_{ij}^{s,ew} - \mathbf{r}_{ij}^{s,ew}$ [14, 20, 21].

The effective coefficients multiplying the matrix elements $\langle sq'\bar{q}'|O_j^{(q)}|b\rangle^{\rm tree}$ become

$$\begin{aligned}
c_3^{\rm eff} &= \bar{c}_3 - \frac{1}{N}\frac{\alpha_s}{8\pi}(c_t + c_p) \\
c_4^{\rm eff} &= \bar{c}_4 + \frac{\alpha_s}{8\pi}(c_t + c_p) \\
c_5^{\rm eff} &= \bar{c}_5 - \frac{1}{N}\frac{\alpha_s}{8\pi}(c_t + c_p) \\
c_6^{\rm eff} &= \bar{c}_6 + \frac{\alpha_s}{8\pi}(c_t + c_p) \\
c_7^{\rm eff} &= \bar{c}_7 + \frac{\alpha_{em}}{8\pi}c_e
\end{aligned}$$

$$
\begin{aligned}
c_8^{\rm eff} &= \bar{c}_8 \\
c_9^{\rm eff} &= \bar{c}_9 + \frac{\alpha_{em}}{8\pi} c_e \\
c_{10}^{\rm eff} &= \bar{c}_{10},
\end{aligned}
\tag{8}
$$

where we have separated the contributions c_t and c_p from the "tree" operators $O_{1,2}$ and from the strong penguin operators $O_{3\ldots6}$, respectively. c_e, given below, comes from the electroweak penguins.

In addition to the contributions from penguin diagrams with insertions of the tree operators $O_{1,2}^{(q)}$

$$
c_t = \bar{c}_2 \cdot \left[\frac{10}{9} + \frac{2}{3} \ell n \frac{m_q^2}{\mu^2} - \Delta F_1\Big(\frac{k^2}{m_q^2}\Big) \right] \ , \tag{9}
$$

where ΔF_1 is defined in [12], we have evaluated the penguin diagrams for the matrix elements of the penguin operators [12]:

$$
\begin{aligned}
c_p = \bar{c}_3 \cdot & \left[\frac{280}{9} + \frac{2}{3} \ell n \frac{m_s^2}{\mu^2} + \frac{2}{3} \ell n \frac{m_b^2}{\mu^2} - \Delta F_1\Big(\frac{k^2}{m_s^2}\Big) - \Delta F_1\Big(\frac{k^2}{m_b^2}\Big) \right] \\
& + (\bar{c}_4 + \bar{c}_6) \cdot \sum_{j=u,d,s,\ldots} \left[\frac{10}{9} + \frac{2}{3} \ell n \frac{m_j^2}{\mu^2} - \Delta F_1\Big(\frac{k^2}{m_j^2}\Big) \right] \ ,
\end{aligned}
\tag{10}
$$

For the electroweak penguins we consider only those arising from the insertion of the tree operators. The corresponding coefficient is given by

$$
c_e = \frac{8}{9}(3\bar{c}_1 + \bar{c}_2)\Big(\frac{10}{9} + \frac{2}{3} \ell n \frac{m_q^2}{\mu^2} - \Delta F_1\Big(\frac{k^2}{m_q^2}\Big), \tag{11}
$$

Note that the coefficients $c_i^{\rm eff}$ depend on k^2 and, as we shall see later, on the $q\bar{q}$ states that are included in the sum over the intermediate states.

C. Hadronic Matrix Elements in the BSW Model

To take into account long distance QCD effects which build up the hadronic final states, we follow Bauer, Stech and Wirbel [22]: With the help of the factorization hypothesis the three-hadron matrix elements are split into vacuum-meson and meson-meson matrix elements of the quark currents entering in $O_1, \ldots, O_{10}$. In addition, OZI suppressed form

factors and annihilation terms are neglected. In the BSW model, the meson-meson matrix elements of the currents are evaluated by overlap integrals of the corresponding wave functions and the dependence on the momentum transfer (which is equal to the mass of the factorized meson) is modeled by a single-pole ansatz. As a first approximation, this calculational scheme provides a reasonable method for estimating the relative size and phase of the tree and penguin terms that give rise to the CP-violating signals.

When pseudoscalars are involved there are additional contributions from the $(V+A)$ penguin operator O_6 and O_8: After Fierz reordering and factorization they contribute terms which involve a matrix element of the quark-density operators between a pseudoscalar meson and the vacuum. For O_6, for example, this is given by

$$\langle P_1 P_2|O_6|B\rangle = -2\sum_q \Big(\langle P_1\,|\bar{q}b_L|0\rangle\langle P_2\,|\bar{s}q_R|B\rangle + \langle P_2\,|\bar{q}b_L|0\rangle\langle P_1\,|\bar{s}q_R|B\rangle\Big)\ . \tag{12}$$

Using the Dirac equation, the matrix elements entering here can be rewritten in terms of those involving usual $(V-A)$ currents,

$$\langle P_1 P_2|O_6|B\rangle = R[P_1,P_2]\langle P_1 P_2|O_4|B\rangle\ , \tag{13}$$

with

$$R[P_1,P_2] \equiv \frac{2M_{P_1}^2}{(m_{q1}+m_{\bar{q}1})(m_b-m_{q2})}\ . \tag{14}$$

Here, m_{q1} ($m_{\bar{q}1}$) and m_{q2} are the current masses of the (anti-)quark in the mesons P_1 and P_2, respectively. We use the quark masses $m_u = m_d = 10$ MeV, $m_s = 200$ MeV, $m_c = 1.5$ GeV and $m_b = 4.8$ GeV. The same relations work for O_8.

Finally, one arrives at the form

$$\begin{aligned}\langle P_1\,P_2|\mathcal{H}_{\rm eff}^{(q)}|B\rangle &= Z_1^{(q)}\langle P_1|j^\mu|0\rangle\langle P_2|j_\mu|B\rangle \\ &\quad + Z_2^{(q)}\langle P_2|j'^\mu|0\rangle\langle P_1|j'_\mu|B\rangle\ ,\end{aligned} \tag{15}$$

where j_μ and j'_μ are the corresponding (neutral or charged) $V-A$ currents. The factorization coefficients $Z_1^{(q)}$ and $Z_2^{(q)}$ are listed in the appendix. In terms of the form factors F_0 for the current matrix elements defined by BSW [22], this yields

$$\langle P_1 P_2 | \mathcal{H}_{\rm eff}^{(q)} | B \rangle = Z_1^{(q)} (M_B^2 - M_2^2) f_{P_1} F_0(M_1^2) + Z_2^{(q)} (M_B^2 - M_1^2) f_{P_2} F_0(M_2^2) \tag{16}$$

M_B is the mass of the decaying B meson and f_P is the decay constant of the pseudoscalar mesons in the final state.

Concerning how $1/N$ terms are treated in the coefficients (see (8) and (18)), it is well known [23] that this model has problems accounting for the decays with branching ratios which are proportional to the combination $\bar{c}_1 + \bar{c}_2/N$. This is due to the rather small absolute value of this particular combination when using the short-distance QCD corrected coefficients. An analogous effect is also known in nonleptonic D decays [22], and several authors advocated a modified procedure to evaluate the factorized amplitudes [22, 24]: There, only terms which are dominant in the $1/N$ expansion are taken into account. Recently there has been much discussion in the literature concerning these issues. Fits to measured branching ratios indicate that the effective a_1 coefficient has the opposite sign from that expected if the $1/N$ terms are completely cancelled by non-factorizable terms, as if the cancellation were incomplete, an effect mimicked by taking $N = 2$. In this work we shall quote results for $N = \infty$ and $N = 2$.

The strong phase shifts are generated in our model only by the absorptive parts (hard final state interactions) of the quark-level matrix elements of the effective Hamiltonian. Of course, when factorizing the hadronic matrix elements, all information on the crucial value of the momentum transfer k^2 of the gluon in the penguin diagram is lost. While there has been an attempt [25] to model a more realistic momentum distribution by taking into account the exchange of a hard gluon, we will use here for simplicity only a fixed value of k^2. From simple two body kinematics [26] or from the investigations in ref. [25] one expects k^2 to be typically in the range

$$\frac{m_b^2}{4} \lesssim k^2 \lesssim \frac{m_b^2}{2} \, . \tag{17}$$

The results we shall present are sensitive to k^2 in this range because the $c\bar{c}$ threshold lies between these limits. Arguments have been made that the lower limit is a more appro-

priate choice [27]. In this work we follow [12] and choose the upper limit for our numerical presentation in the tables.

The factorization coefficients $Z_{1,2}^{(q)}$ defined in (15) are listed in Tab. 1 a,b for the $B \to PP$ channels of interest. Here we add tree and penguin contributions. It is understood that they have different CKM factors which must be inserted. Colour suppressed terms may readily be included in the coefficients

$$a_i \equiv c_i^{\rm eff} + \frac{1}{N} c_j^{\rm eff} \tag{18}$$

where $\{i,j\}$ is any of the pairs $\{1,2\}$, $\{3,4\}$, $\{5,6\}$, $\{7,8\}$,or $\{9,10\}$. In Tab. 1 a,b we have adopted the convention of including factors of $\sqrt{2}$ associated with a neutral meson P_2. They arise either from current matrix elements between P_2 and B (Tab. 1 a), or from the definition of the decay constants for P_2 (Tab. 1 b). Care should be taken with the latter since these factors are sometimes absorbed into the decay constants (e.g. as tabulated in [22]).

Our phase convention for the mesons states relative to the weak Hamiltonian is defined by:

$$[\pi^+, \pi^0, \pi^-] = [\bar{u}d, \frac{d\bar{d} - u\bar{u}}{\sqrt{2}}, d\bar{u}] \tag{19}$$

$$(K^+, K^0) = (u\bar{s}, \quad d\bar{s}), \quad (K^-, K^0) = (s\bar{d}, \quad s\bar{u}) \tag{20}$$

$$(B^-, B^0) = (b\bar{u}, \quad b\bar{d}), \quad (B^+, B^0) = (\bar{b}u, \quad \bar{b}d) \tag{21}$$

III. CP-VIOLATING OBSERVABLES

In this section we define our notation and review the CP violating observables, following the treatment of reference [4].

Let M^0 be the neutral meson (i.e. B^0) and $\bar{M}^0$ its antiparticle. M^0 and $\bar{M}^0$ can mix with each other and form two physical mass eigenstates

$$M_1 = p|M^0> +q|\bar{M}^0>, \qquad M_2 = p|M^0> -q|\bar{M}^0> \tag{22}$$

The CP-violating parameter ϵ_M is introduced via

$$\epsilon_M = \frac{1-q/p}{1+q/p}\ , \qquad \frac{q}{p} \equiv \sqrt{\frac{H_{21}}{H_{12}}} \tag{23}$$

where $H_{12} \equiv M_{12} - \frac{i}{2}\Gamma_{12} =< M^0|H_{eff}|\bar{M}^0 >$.

Let f denote the final decay state of the neutral meson and $\bar{f}$ its charge conjugate state. The decay amplitudes of M^0 and $\bar{M}^0$ are denoted by

$$g \equiv< f|H_{eff}|M^0 >,\ h \equiv< f|H_{eff}|\bar{M}^0 >;\ \bar{g} \equiv< \bar{f}|H_{eff}|\bar{M}^0 >,\ \bar{h} \equiv< \bar{f}|H_{eff}|M^0 > \tag{24}$$

Parameters containing direct CP violation are defined by

$$\epsilon'_M \equiv \frac{1-h/g}{1+h/g},\ \bar{\epsilon}'_M \equiv \frac{1-\bar{g}/\bar{h}}{1+\bar{g}/\bar{h}};\ \epsilon''_M \equiv \frac{1-\bar{g}/g}{1+\bar{g}/g},\ \bar{\epsilon}''_M \equiv \frac{1-h/\bar{h}}{1+h/\bar{h}} \tag{25}$$

In terms of these, the rephase-invariant observables are:

$$\begin{aligned}
a_\epsilon &= \frac{1-|q/p|^2}{1+|q/p|^2} = \frac{2Re\epsilon_M}{1+|\epsilon_M|^2}\ , \qquad a_{\epsilon'} = \frac{1-|h/g|^2}{1+|h/g|^2} = \frac{2Re\epsilon'_M}{1+|\epsilon'_M|^2}\ ; \\
a_{\epsilon+\epsilon'} &= \frac{-4Im(qh/pg)}{(1+|q/p|^2)(1+|h/g|^2)} = \frac{2Im\epsilon_M(1-|\epsilon'_M|^2)+2Im\epsilon'_M(1-|\epsilon_M|^2)}{(1+|\epsilon_M|^2)(1+|\epsilon'_M|^2)} \\
a_{\epsilon\epsilon'} &= \frac{4Re(qh/pg)}{(1+|q/p|^2)(1+|h/g|^2)} - 1 = \frac{4Im\epsilon_M\ Im\epsilon'_M - 2(|\epsilon_M|^2+|\epsilon'_M|^2)}{(1+|\epsilon_M|^2)(1+|\epsilon'_M|^2)}
\end{aligned} \tag{26}$$

Only three of them are independent as $(1-a_\epsilon^2)(1-a_{\epsilon'}^2) = a_{\epsilon+\epsilon'}^2 + (1+a_{\epsilon\epsilon'})^2$. Analogously, one has observables $a_{\bar{\epsilon}'}$, $a_{\epsilon+\bar{\epsilon}'}$ and $a_{\epsilon\bar{\epsilon}'}$ similar to $a_{\epsilon'}$, $a_{\epsilon+\epsilon'}$ and $a_{\epsilon\epsilon'}$ but with ϵ'_M being replaced by $\bar{\epsilon}'_M$.

Two additional rephase-invariant quantities complete the set of observables,

$$a_{\epsilon''} = \frac{1-|\bar{g}/g|^2}{1+|\bar{g}/g|^2} = \frac{2Re\epsilon''_M}{1+|\epsilon''_M|^2}, \qquad a_{\bar{\epsilon}''} = \frac{1-|\bar{h}/h|^2}{1+|\bar{h}/h|^2} = \frac{2Re\bar{\epsilon}''_M}{1+|\bar{\epsilon}''_M|^2} \tag{27}$$

To apply this general analysis to the specific processes considered in this work, consider the following two cases:

i) $M^0 \to f$ $(M^0 \not\to \bar{f})$, $\overline{M}^0 \to \bar{f}$ $(\overline{M}^0 \not\to f)$, i.e., f or $\bar{f}$ is not a common final state of M^0 and $\overline{M}^0$. This applies to the channel $\bar{B}^0 \to K^-\pi^+$.

ii) $M^0 \to (f = \bar{f},\ f^{CP} = f) \leftarrow \overline{M}^0$, i.e., final states are CP eigenstates. This applies to the channels $\bar{B}^0 \to \pi^+\pi^-,\ \pi^0\pi^0, K_s\pi^0, \bar{K}K$.

In the following we will make the good approximation that $a_\epsilon = 0$. In the scenario i), one has: $a_{\epsilon'} = -a_{\bar{\epsilon}'} = 1$, $a_{\epsilon+\epsilon'} = 0 = a_{\epsilon+\bar{\epsilon}'}$ and $a_{\epsilon\epsilon'} = -1 = a_{\epsilon\bar{\epsilon}'}$. The time-dependent rate asymmetry is:

$$A_{CP}(t) = \frac{\Gamma(M^0(t) \to f) - \Gamma(\overline{M}^0(t) \to \bar{f}}{\Gamma(M^0(t) \to f) + \Gamma(\overline{M}^0(t) \to \bar{f})} = a_{\epsilon''} \tag{28}$$

There is a second asymmetry corresponding to the last expression but with $\overline{M}^0$ replaced by M^0. When $a_\epsilon = 0$ these asymmetries are equal. This asymmetry also applies to charged decays.

In the scenario ii) in which $a_{\epsilon'} = a_{\epsilon''} = a_{\bar{\epsilon}'} = a_{\bar{\epsilon}''}$ and $a_{\epsilon+\epsilon'} = a_{\epsilon+\bar{\epsilon}'}$, the time-dependent CP asymmetry is:

$$A_{CP}(t) \simeq \frac{a_{\epsilon'}\cos(\Delta m t) + a_{\epsilon+\epsilon'}\sin(\Delta m t)}{\cosh(\Delta\Gamma t) + (1 + a_{\epsilon\epsilon'})\sinh(\Delta\Gamma t)} \tag{29}$$

If $|\Delta\Gamma| \ll |\Delta m|$ and $|\Delta\Gamma/\Gamma| \ll 1$ then, $A_{CP}(t)$ further simplifies

$$A_{CP}(t) \simeq a_{\epsilon'}\cos(\Delta m t) + a_{\epsilon+\epsilon'}\sin(\Delta m t) \tag{30}$$

which may be applied, in a good approximation, to the $B^0 - \bar{B}^0$ system.

Because of the notation in this section, a'_ϵ is the rate asymmetry for the CP eigenstates, case (ii) decays, whereas $a_{\epsilon''}$ is the rate asymmetry for case (i) and charged decays. To simplify the tables of the next section, the rate asymmetry always appears labeled by a'_ϵ in the tables.

IV. RESULTS AND DISCUSSION OF RATES, ASYMMETRIES, AND SUM RULES

In Tab. 2 and 3 we present the decay parameters for the $\pi\pi$, $K\pi$ and KK channels. In both tables we use the CKM parameters of ref. [15] with ρ negative and small, $\rho = -0.12, \eta = 0.34$. Tab. 2 is for $N = \infty$ and Tab. 3 is for $N = 2$. For each channel (first column) we present results for the four cases: (t+p+e), (t+p), (t+p') and (t), second column. By (t) we denote tree contributions, color dominant or color suppressed. By (p) we denote strong penguin contributions, including the effect of the strong phase of our model. By (p') we denote the penguin without the strong phase. By (e) we denote electroweak penguin contributions including the absorptive parts. One purpose of this work is to investigate the influence of the penguin contributions with absorptive parts on the time dependent asymmetries for the decay of neutral B mesons into CP eigenstates $f = \bar{f} = f_{CP}$. Such cases are the decays $\bar{B}^0 \to \pi^0\pi^0, \pi^+\pi^-, \bar{K}^0\pi^0, (K_s\pi^0)$ and $K^0\bar{K}^0$. The asymmetry $A_{CP}(t)$ (see (31)) has two terms, one proportional to $cos(\Delta mt)$ with the coefficient a'_ϵ, the other one proportional to $sin(\Delta mt)$ with coefficient $a_{\epsilon+\epsilon'}$. These two coefficients are given in the third and fourth columns of Tab. 2 and 3, respectively. For the channels that are not CP eigenstates, there is no entry because there is no time dependent asymmetry. In the fifth column we give the rates averaged over B and $\bar{B}$. In the last two columns we present the real and imaginary part of the amplitude as defined in (15).

The rates and asymmetry parameters are quite different for $N = \infty$ and $N = 2$ so we consider these cases separately, first in the approximation of neglecting the electroweak penguins.

A. $N = \infty$, Without Electroweak Penguins

Let us consider first the results for $\bar{B}^0 \to \pi^0\pi^0$ with $1/N = 0$. Here the influence of the penguins is very large, since C (color suppressed tree) and strong penguin P are of the same

order of magnitude and have the same sign. The amplitude of this decay is $\sim (P-C)$, which is small so that with strong penguins the sign of the real part of the amplitude changes. This has an effect on $a_{\epsilon+\epsilon'}$, which for t + p' also changes sign, so that in this case $a_{\epsilon+\epsilon'}$ changes dramatically compared to the tree value $a_{\epsilon+\epsilon'} = -sin2\alpha$. With absorptive parts added (t + p) we generate a non-zero value for a'_ϵ, which now is of the same order of magnitude as $a_{\epsilon+\epsilon'}$ completely changing the time dependence of $A_{CP}(t)$. We see that $a_{\epsilon+\epsilon'}$ is influenced very little by the absorptive part. The branching ratios are small in all cases, about $5 \cdot 10^{-7}$ even when penguin terms are included.

The pattern is similar for the decay $\bar{B}^0 \to \bar{K}^0\pi^0$. The influence of the penguin terms is even stronger as one sees in the real part of the amplitude which changes by a factor of 40, increasing the branching ratio by two orders of magnitude. The influence on $a_{\epsilon+\epsilon'}$ is also strong, as in the previous case. The other parameter a'_ϵ is non-vanishing when we add the absorptive parts, but is only about -3% so that it has little influence on A_{CP}.

Next we discuss the decay $\bar{B}^0 \to \pi^+\pi^-$ which is most interesting since it is considered as the decay channel by which the angle α could be measured from the asymmetry $A_{CP}(t)$. This is governed by the two parameters a'_ϵ and $a_{\epsilon+\epsilon'}$. The second parameter $a_{\epsilon+\epsilon'} = -sin2\alpha$ for a pure tree amplitude. In Tab. 2 we observe that $a_{\epsilon+\epsilon'}$ is decreased by 20% by the strong penguin contributions. The absorptive parts change this additionally by only the negligible amount of 2% in absolute value. The parameter a'_ϵ which determines the $cos(\Delta mt)$ contribution of $A_{CP}(t)$ is non-negligible, $a'_\epsilon/a_{\epsilon+\epsilon'} \sim -0.1$.

The last decay with a time dependent asymmetry is the pure penguin mode $\bar{B}^0 \to K^0\bar{K}^0$. The parameters $a_{\epsilon+\epsilon'}$ and a_ϵ are almost equal when absorptive parts are included. They are of the order of 10% so that the total time dependence is not very large. The change of $a_{\epsilon+\epsilon'}$ by the absorptive parts of the penguins is only about 20%.

In Tab. 2 we also present the results for the decay $\bar{B}^0 \to K^-\pi^+$. This final state is not a CP eigenstate and we must apply the formalism of case (i) above. The only physical parameters are the branching ratio and the parameter $a_{\epsilon''}$ defined in (25) above. This parameter is listed in the third column of Tab. 2 and 3 respectively with the label $a_{\epsilon'}$.

The asymmetry is of the order of 10% and is negative. As we see, this asymmetry is time independent and not affected by mixing. We notice that $a_{\epsilon'}(\pi^+\pi^-) \simeq -a_{\epsilon'}(K^-\pi^+)$ to a very good approximation. This result can be easily explained from the formulas and also through the recent work of Deshpande and He [9].

The decays $B^- \to \pi^0\pi^-, K^-\pi^0, \bar{K}^0\pi^-, K^0K^-$ have been considered in our earlier work [12]. Of interest are the results for the branching ratios and the rate asymmetry parameter $a_{\epsilon'}$. Since in this work we assumed different values of the CKM parameters ρ and η the asymmetry has changed. The different sign as compared to our earlier work has to do with the fact the the asymmetry is the difference $B^0 - \bar{B}^0$ [4] rather the opposite used in ref [12].

B. N=2, Without Electroweak Penguins

When we now look back at the more realistic case with $N = 2$, which accounts for the sign change in the QCD coefficient a_1 (See(18)) the pattern of the results does not change very much, but with two exceptions. In the decay $\bar{B}^0 \to \pi^0\pi^0$ the interference between tree and penguin contributions in the amplitude proportional to (P-C) is now different. This has the effect that $a_{\epsilon'}$ is smaller now and has the opposite sign. Furthermore the effect of the absorptive penguin terms on $a_{\epsilon+\epsilon'}$ is decreased and the branching ratio increases now due to penguin contributions. In the decay $\bar{B}^0 \to \bar{K}^0\pi^0$ the a'_ϵ is also changed; it is smaller and has the opposite sign. $a_{\epsilon+\epsilon'}$ is increased and the branching ratio is somewhat smaller than in the $N = \infty$ case. Of course the result $a_{\epsilon'}(\pi^+\pi^-) \simeq -a_{\epsilon'}(K^-\pi^+)$ is still valid. In ref. [9] one can see explicitly that this relation is independent of the details of the QCD coefficients.

We conclude that independent of details of the QCD coefficients the influence of the penguins as compared to the tree contributions is very significant in the channels $\pi^0\pi^0, \bar{K}^0\pi^0$. The penguins influence rates and also the time dependence through $a_{\epsilon+\epsilon'}$. The influence of the absorptive parts on the time dependence is not very important. They do influence $a_{\epsilon+\epsilon'}$ and add an additional time dependent term in the decay $\bar{B}^0 \to \pi^0\pi^0$ in the $N = \infty$ version where the C and P terms interfere destructively.

C. Electroweak Penguins

The most significant electroweak penguin operator is the term proportional to c_9 which influences coefficient a_9 and a_{10} (much less) in (18). Therefore we expect channels in Tab. 1 with a_9 coefficients to be most strongly affected by the electroweak penguins. This is evident from Tab. 2 and 3 for the $\pi^0\pi^0, \pi^0\pi^-, \bar{K}^0\pi^0$ and $K^-\pi^0$ channels if we look at amplitudes. The change is stronger in ReA. The change of the imaginary part is only 30% of the change of the real part for the $\pi^0\pi^0$ and $\pi^0\pi^-$ and negligible in the imaginary part for the $\bar{K}^0\pi^0$ and $K^-\pi^0$ channels. This is because the CKM phase is real in the $K\pi$ channels. It is interesting to note that for $N = \infty$ the time dependence of the asymmetry for the decay $\bar{B}^0 \to \pi^0\pi^0$ is significantly influenced by the change of $a_{\epsilon+\epsilon'}$. This is not the case for $N = 2$, where $a_{\epsilon'}$ is also smaller. From Tab. 2 and 3 it is evident that the electromagnetic penguins can influence branching ratios substantially as shown by the $K^-\pi^0$ and $\bar{K}^0\pi^0$ channels.

D. Effect of Penguins on Flavor Sum Rules

In this section we shall examine the effects of the SU(3) breaking in the matrix elements and the influence of the EW penguins on the sum rules of [5, 6, 7]. We shall start with the sum rule for $B \to \pi\pi$ amplitudes based on SU(2) symmetry [5], which with our definition of quark content in (19-21) is:

$$A = B \tag{31}$$

where

$$A = \sqrt{2}A(B^- \to \pi^-\pi^0) \tag{32}$$

$$B = \sqrt{2}A(\bar{B}^0 \to \pi^0\pi^0) - A(\bar{B}^0 \to \pi^+\pi^-) \tag{33}$$

If we neglect the $\pi^\pm - \pi^0$ mass differences etc, the amplitude A and B have according to Tab. 1 the following form in terms of the coefficients $\hat{a}_i$ which contain in addition to the Wilson coefficients a_i the appropriate CKM factors for tree and penguin contributions,

$$A = B = M\{-(\hat{a}_1 + \hat{a}_2) + \frac{3}{2}(\hat{a}_7 - \hat{a}_9) - \frac{3}{2}(\hat{a}_{10} + \hat{a}_8 R[\pi\pi])\} \tag{34}$$

where M is the reduced matrix element assuming SU(2) symmetry and factorization which can be read off from (16) and Tab. 1. This means that the sum rule $A = B$ is not spoiled by penguin contributions even when the EW penguins are included. In (34) $\hat{a}_2(\hat{a}_1)$ can be identified with the tree (color suppressed tree) and $\frac{3}{2}(\hat{a}_7 - \hat{a}_9)[\frac{3}{2}(\hat{a}_{10} + \hat{a}_8 R)]$ with the EW penguin [color suppressed EW penguin]. That $A = B$ remains valid with the EW penguin terms stems from the fact that the additional contributions proportional to $O_{7...10}$ transform as the same mixture of I=1/2 and 3/2 operators as the tree operators $O_{1,2}$. The strong penguins on the other hand have only a single isospin component. This explains why the amplitudes A and B have no strong penguin contributions. In Tab. 4 and 5 we give the numerical results of the amplitudes A and B in our model. As we see $A = B$ is very well satisfied in all four versions of the model calculation: t, t+p', t + p and t+p+e. By comparing the results for t+p+e with t+p it is clear that the electroweak penguins make substantial contributions ($\sim 15\%$ in $ReA, \sim 2\%$ in ImA) but preserve the sum rule for the reason given above. There are therefore substantial I=3/2 pieces in the amplitudes A and B which do not originate from the tree operators $O_{1,2}$.

Similarly, we have for the $B \to K\pi$ transitions the sum rule

$$C = D \tag{35}$$

where

$$C = [\sqrt{2}A(\bar{B}^0 \to \bar{K}^0\pi^0) - A(\bar{B}^0 \to K^-\pi^+)]/r_u \tag{36}$$

$$D = [\sqrt{2}A(B^- \to K^-\pi^0) + A(B^- \to \bar{K}^0\pi^-)]/r_u \tag{37}$$

where $r_u = V_{us}/V_{ud}$ has been inserted for later use.

Since for $B \to K\pi$ two different reduced matrix elements occur in (16), the amplitudes C and D are now given by

$$C = D = M_2/r_u\{-\hat{a}'_1 + \frac{3}{2}(\hat{a}'_7 - \hat{a}'_9)\} + M_1/r_u\{-\hat{a}'_2 - \frac{3}{2}(\hat{a}'_{10} + \hat{a}'_8 R[K,\pi])\} \tag{38}$$

where the $M_1(M_2)$ are the reduced matrix elements in (16) proportional to the $Z_1^{(q)}(Z_2^{(q)})$ and the $\hat{a}'_i$ are the Wilson coefficients multiplied by the appropriate CKM matrix elements. As in the case $A = B$ the sum rule $C = D$ is broken only by small SU(2) mass differences. Futhermore the strong penguin amplitudes drop out in the amplitudes C and D for the same reason they did in A and B. In Tab. 4 and 5 we can see that $C = D$ is quite well satisfied in all four model versions except perhaps for the ReA in t+p which is caused by the symmetry breaking effects in connection with $O(\alpha_s)$ corrections. By comparing the values for t+p+e with t+p we observe that the EW penguin effects are very large in ReA (factor ~ 3) but again preserve the sum rule $C = D$.

An important sum rule based on SU(3) symmetry and absence of EW penguins in reference [7] is

$$A = C \tag{39}$$

In contrast to the previously considered sum rules, different matrix elements are now involved, primarily those with different factors $f_K/f_\pi \neq 1$. If we compare (34) with (38) it is clear that the sum rule $A = C$ is valid if $M = M_1/r_u = M_2/r_u$ and if there are no EW penguins. Thus when EW penguins are neglected, we would expect $(f_K/f_\pi)A = C$ which is approximately valid in Tab. 4 and 5. The factor $1/r_u$ corrects for the different CKM factors in the $b \to d$ versus $b \to s$ transitions in the tree amplitudes. Because the electroweak penguin terms in (39) have a different CKM factor this sum rule is strongly violated by the EW penguin. Indeed if we compare $(f_K/f_\pi)A$ with C for t+p+e we observe a factor of 3.5 difference in ReA.

The last relation considered in reference [7] is the quadrangle relation $E = F$. For our model including strong and EW penguins we have

$$E = (\hat{a}_2 + \hat{a}_4 + \hat{a}_6 R[\pi,\pi] + \hat{a}_{10} + \hat{a}_8 R[\pi,\pi])M - (\hat{a}_4 + \hat{a}_6 R[K,K] - \frac{1}{2}(\hat{a}_{10} + \hat{a}_8 R[K,K]))M_K \tag{40}$$

and

$$F = [\hat{a}'_2 + \frac{1}{2}(\hat{a}'_{10} + \hat{a}'_8 R[K,\pi])]M_1/r_u \tag{41}$$

It is clear from (40,41) that the relation $E = F$ is badly spoiled by the SU(3) symmetry breaking in the reduced matrix elements, i.e. $M \neq M_K \neq M_1$ and $R[\pi,\pi] \neq R[K,K] \neq R[K,\pi]$. The EW penguins in the amplitude E are very small since they are color suppressed and Cabibbo suppressed compared to the tree term (compare line t+p+e with t+p in Tab. 4 and 5 for case E). In contrast the EW penguin of the amplitude F while color suppressed is not Cabibbo suppressed, which explains the larger difference by almost a factor of two. This has the consequence that apart from the SU(3) breaking the sum rule $E = F$ is very badly violated for the ReA by a factor of 2.5 ($N = \infty$) or 1.7 ($N = 2$), respectively.

This calculation supports the criticism of reference [8, 9] to the scheme for extracting weak phases advocated in reference [7]. The latter authors have subsequently modified their analysis to include EW penguins [10]. By making additional measurements of B_s decays, they argue that it is still possible to extract the CKM angles from rate measurements alone. A somewhat different scheme has been proposed by Deshpande and He [28] involving the additional measurement of the rate for $B^- \to K^-\eta_8$ and the assumption of no EW penguin contributions in $B^- \to \pi^0\pi^-$; as we see in our model this is valid up to 15% for ReA and 2% in ImA, as shown in Tab. 4 and 5.

V. SUMMARY AND CONCLUSIONS

We have reported the results of a model of trees, strong and EW penguins with absorptive corrections. This model is subject to various well known uncertainties associated with soft physics. In addition we had to use badly determined CKM parameters. Nevertheless we believe our results give the right order of magnitude of the effect of strong and EW penguins and absorptive corrections. They indicate that certain flavor sum rules are spoiled by the EW penguins and SU(3) symmetry violation. As expected, SU(2) sum rules are preserved,

even those with EW penguin effects. Strong and weak penguins have important effects on the branching ratios and the CP violation parameters, with details depending on the model of the effective Wilson coefficients ($N = \infty$ or $N = 2$). By calculating more rates one could test the procedures of Deshpande and He [8, 28] and Gronau et. al. [10] for determining the CP violating phase γ.

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TABLE CAPTIONS

Tab. 1a,b: Factorization coefficients $Z_{1,2}^{(q)}$ for various $B \to P_1 P_2$ decays. The short distance coefficients a_i are defined in (18) and the factor R is given in (14). The coefficients do not include the appropriate CKM phases as in (1) and (2).

Tab. 2: Decay and CP violating parameters in $B \to \pi\pi, \pi K, K\bar{K}$. The amplitudes are calculated in a model with tree (t) and strong penguins p (p') with (without) absorptive parts and electroweak penguins (e) with absorptive parts. The model parameters are NLL QCD Coefficients, BSW model for matrix elements, Wolfenstein parameters: $(\rho, \eta) = (-0.12, 0.34)$, $N = \infty$.

Tab. 3: Decay and CP violating parameters in $B \to \pi\pi, \pi K, K\bar{K}$. The amplitudes are calculated in a model with tree (t) and strong penguins p (p') with (without) absorptive parts and electroweak penguins (e) with absorptive parts. The model parameters are NLL QCD Coefficients, BSW model for matrix elements, Wolfenstein parameters: $(\rho, \eta) = (-0.12, 0.34)$, $N = 2$.

Tab. 4: Sum rules for reduced $B \to \pi\pi, \pi K, K\bar{K}$ amplitudes. The amplitudes are calculated in a model with tree (t) and strong penguins p (p') with (without) absorptive parts and electroweak penguins (e) with absorptive parts. The model parameters are NLL QCD Coefficients, BSW model for matrix elements, Wolfenstein parameters: $(\rho, \eta) = (-0.12, 0.34)$, $N = \infty$.

Tab. 5: Sum rules for reduced $B \to \pi\pi, \pi K, K\bar{K}$ amplitudes. The amplitudes are calculated in a model with tree (t) and strong penguins p (p') with (without) absorptive parts and electroweak penguins (e) with absorptive parts. The model parameters are NLL QCD Coefficients, BSW model for matrix elements, Wolfenstein parameters: $(\rho, \eta) = (-0.12, 0.34)$, $N = 2$.

Tab. 1a

P_1	P_2	$Z_1^{(q)}$
$\bar{K}^0$	π^0	$(a_4 - a_{10}/2 + (a_6 - a_8/2)\, R[\bar{K}^0, \pi^0])/\sqrt{2}$
K^-	π^+	$a_2\delta_{qu} + a_4 + a_{10} + (a_6 + a_8)\, R[K^-, \pi^+]$
K^-	π^0	$-(a_2\delta_{qu} + a_4 + a_{10} + (a_6 + a_8)\, R[K^-, \pi^0])/\sqrt{2}$
$\bar{K}^0$	π^-	$a_4 - a_{10}/2 + (a_6 - a_8/2)\, R[K^0, \pi^-]$
π_1^0	π_2^0	$(-a_1\delta_{qu} + a_4 - a_{10}/2 + 3a_7/2 - 3a_9/2 + (a_6 - a_8/2)\, R[\pi_1^0, \pi_2^0])/2$
π^+	π^-	0
π^-	π^0	$(-a_2\delta_{qu} - a_4 - a_{10} - (a_6 + a_8)\, R[\pi^-, \pi^0])/\sqrt{2}$
K^0	K^-	$a_4 - a_{10}/2 + (a_6 - a_8/2)\, R[K^0, K^-]$
K^0	$\bar{K}^0$	$a_4 - a_{10}/2 + (a_6 - a_8/2)\, R[K^0, \bar{K}^0]$

Tab. 1b

P_1	P_2	$Z_2^{(q)}$
$\bar{K}^0$	π^0	$(-a_1\delta_{qu} + 3a_7/2 - 3a_9/2)\sqrt{2}$
K^-	π^+	0
K^-	π^0	$(-a_1\delta_{qu} + 3a_7/2 - 3a_9/2)/\sqrt{2}$
$\bar{K}^0$	π^-	0
π_1^0	π_2^0	$(-a_1\delta_{qu} + a_4 - a_{10}/2 + 3a_7/2 - 3a_9/2 + (a_6 - a_8/2)\, R[\pi_2^0, \pi_1^0])/2$
π^+	π^-	$a_2\delta_{qu} + a_4 + a_{10} + (a_6 + a_8)\, R[\pi^-, \pi^+]$
π^-	π^0	$(-a_1\delta_{qu} + a_4 + 3a_7/2 - 3a_9/2 - a_{10}/2 + (a_6 - a_8/2)\, R[\pi^0, \pi^-])/\sqrt{2}$
K^0	K^-	0
K^0	$\bar{K}^0$	0

Tab. 2

Decay Parameters and CP Violation in $B \to \pi\pi, \pi K, K\bar{K}$ Tree, Strong and EW Penguins t,p (p'), e with (without) Absorptive Parts $N = \infty$ NLL QCD Coefficients, BSW model, $(\rho, eta) = (-0.12, 0.34)$						
		Decay Parameters				
Channel	Amp	$a_{\epsilon'}$	$a_{\epsilon+\epsilon'}$	$< BR >$	Re A$\times 10^3$	Im A$\times 10^3$
	t + p + e	0.308	-0.126	4.22×10^{-7}	0.256	-1.40
	t + p	0.316	0.201	4.25×10^{-7}	0.528	-1.32
$\pi^0\pi^0$	t + p'	0.0	0.216	4.14×10^{-7}	0.654	-1.56
	t	0.0	-0.951	5.66×10^{-7}	-0.658	-1.86
	t + p + e	0.0703	-0.784	1.19×10^{-5}	-1.65	-8.59
	t + p	0.0708	-0.766	1.19×10^{-5}	-1.59	-8.57
$\pi^+\pi^-$	t + p'	0.0	-0.777	1.18×10^{-5}	-1.41	-8.93
	t	0.0	-0.951	1.43×10^{-5}	-3.31	-9.37
	t + p + e	0.0000	–	3.50×10^{-6}	1.46	4.69
	t + p	0.0	–	3.69×10^{-6}	1.68	4.76
$\pi^0\pi^-$	t + p'	0.0	–	3.69×10^{-6}	1.68	4.76
	t	0.0	–	3.69×10^{-6}	1.68	4.76
	t + p + e	-0.0365	0.435	5.49×10^{-6}	-6.09	-1.56
	t + p	-0.0265	0.457	7.78×10^{-6}	-7.29	-1.54
$\bar{K}^0\pi^0$	t + p'	0.0	0.454	7.58×10^{-6}	-7.25	-0.384
	t	0.0	-0.951	3.68×10^{-8}	-0.168	-0.476
	t + p +e	-0.0760	–	1.74×10^{-5}	-10.6	-4.11
	t + p	-0.0725	–	1.82×10^{-5}	-10.9	-4.10
$K^-\pi^+$	t + p'	0.0	–	1.79×10^{-5}	-10.8	-2.47
	T	0.0	–	9.28×10^{-7}	-0.909	-2.60
	t + p + e	-0.0458	–	11.1×10^{-6}	8.66	2.41
	t + p	-0.0587	–	8.56×10^{-6}	7.56	2.43
$K^-\pi^0$	t + p'	0.0	–	8.37×10^{-6}	7.52	1.27
	t	0.0	–	3.01×10^{-7}	0.481	1.36
Pure Penguin Modes						
	p + e	0.0922	0.0984	1.27×10^{-6}	2.58	1.18
	p	0.0937	0.0951	1.23×10^{-6}	2.54	1.17
$K^0\bar{K}^0$	p'	0.0	0.118	1.18×10^{-6}	2.81	0.654
	p + e	0.0920	–	1.30×10^{-6}	2.62	1.20
	p	0.0936	–	1.27×10^{-6}	2.58	1.19
K^0K^-	p'	0.0	–	1.22×10^{-6}	2.85	0.664
	p + e	-0.00606	–	1.52×10^{-5}	-10.2	-1.51
	p	-0.00616	–	1.48×10^{-5}	-10.1	-1.51
$\bar{K}^0\pi^-$	p'	0.0	–	1.44×10^{-5}	-10.0	0.131

Tab. 3

Decay Parameters and CP Violation in $B \to \pi\pi, \pi K, K\bar{K}$ Tree, Strong and EW Penguins t,p (p'), e with (without) Absorptive Parts $N = 2$ NLL QCD Coefficients, BSW model, $(\rho, eta) = (-0.12, 0.34)$						
		Decay Parameters				
Channel	Amp	$a_{\epsilon'}$	$a_{\epsilon+\epsilon'}$	$< BR >$	Re A$\times 10^3$	Im A$\times 10^3$
	t + p + e	-0.0834	-0.938	6.11×10^{-7}	1.17	1.79
	t + p	-0.0615	-0.877	7.76×10^{-7}	1.47	1.88
$\pi^0\pi^0$	t + p'	0.0	-0.881	7.69×10^{-7}	1.56	1.70
	t	0.0	-0.951	3.41×10^{-7}	0.511	1.45
	t + p + e	0.0620	-0.781	8.75×10^{-6}	-1.37	-7.40
	t + p	0.0609	-0.793	8.83×10^{-6}	-1.46	-7.43
$\pi^+\pi^-$	t + p'	0.0	-0.794	8.82×10^{-6}	-1.32	-7.69
	t	0.0	-0.951	1.06×10^{-5}	-2.84	-8.05
	t + p + e	0.000399	–	7.83×10^{-6}	2.14	7.03
	t + p	0.0	–	8.31×10^{-6}	2.52	7.14
$\pi^0\pi^-$	t + p'	0.0	–	8.31×10^{-6}	2.52	7.14
	t	0.0	–	8.31×10^{-6}	2.52	7.14
	t + p + e	0.0264	0.697	2.70×10^{-6}	-4.26	-0.444
	t + p	0.0145	0.670	4.58×10^{-6}	-5.59	-0.424
$\bar{K}^0\pi^0$	t + p'	0.0	0.672	4.48×10^{-6}	-5.56	0.446
	t	0.0	-0.951	2.21×10^{-8}	0.131	0.370
	t + p + e	-0.0659	–	1.30×10^{-5}	-9.22	-3.35
	t + p	-0.0709	–	1.20×10^{-5}	-8.83	-3.35
$K^-\pi^+$	t + p'	0.0	–	1.18×10^{-5}	-8.79	-2.13
	t	0.0	–	8.09×10^{-7}	-0.789	-2.24
	t + p + e	-0.0560	–	9.39×10^{-6}	7.85	2.72
	t + p	-0.0809	–	6.40×10^{-6}	6.38	2.74
$K^-\pi^0$	t + p'	0.0	–	6.29×10^{-6}	6.35	1.88
	t	0.0	–	6.16×10^{-7}	0.689	1.95
Pure Penguin Modes						
	p + e	0.0886	0.0959	7.57×10^{-7}	2.00	0.907
	p	0.0862	0.0940	7.95×10^{-7}	2.06	0.925
$K^0\bar{K}^0$	p'	0.0	0.110	7.68×10^{-7}	2.26	0.536
	p + e	0.0883	–	7.82×10^{-7}	2.04	0.922
	p	0.0860	–	8.21×10^{-7}	2.09	0.939
K^0K^-	p'	0.0	–	7.94×10^{-7}	2.30	0.546
	p + e	-0.00584	–	9.09×10^{-6}	-7.90	-1.12
	p	-0.00569	–	9.54×10^{-6}	-8.09	-1.12
$\bar{K}^0\pi^-$	p'	0.0	–	9.32×10^{-6}	-8.05	0.108

Tab. 4

Sum Rules for Reduced $B \to \pi\pi, \pi K, K\bar{K}$ Amplitudes Tree, Strong and EW Penguins t,p (p'), e with (without) Absorptive Parts $N = \infty$ NLL QCD Coefficients, BSW model, $(\rho, \eta) = (-0.12, 0.34)$			
Channel	Amp	ReA$\times 10^3$	ImA$\times 10^3$
Sum Rules: A=B=C=D and E=F			
$A = \sqrt{2}\pi^0\pi^-$	t + p + e	2.06	6.63
	t + p	2.37	6.73
	t + p'	2.37	6.73
	t	2.37	6.73
$B = \sqrt{2}\pi^0\pi^0 - \pi^+\pi^-$	t + p + e	2.01	6.61
	t + p	2.34	6.70
	t+ p'	2.34	6.72
	t	2.37	6.74
$C = (\sqrt{2}\bar{K}^0\pi^0 - K^-\pi^+)/r_u$	t + p + e	8.78	8.40
	t + p	3.20	8.49
	t + p'	2.70	8.50
	t	2.96	8.53
$D = (\sqrt{2}K^-\pi^0 + \bar{K}^0\pi^-)/r_u$	t + p + e	8.95	8.40
	t + p	2.80	8.54
	t + p'	2.60	8.54
	t	3.00	8.44
$E = \pi^+\pi^- - K^-K^0$	t + p + e	-4.27	-9.79
	t + p	-4.17	-9.76
	t + p'	-4.26	-9.59
	t	-3.31	-9.37
$F = (K^-\pi^+ - \bar{K}^0\pi^-)/r_u$	t + p + e	-1.70	-11.45
	t + p	-3.70	-11.44
	t + p'	-3.40	-11.47
	t	-4.01	-11.5

Tab. 5

Sum Rules for Reduced $B \to \pi\pi, \pi K, K\bar{K}$ Amplitudes Tree, Strong and EW Penguins t,p (p'), e with (without) Absorptive Parts $N = 2$ NLL QCD Coefficients, BSW model, $(\rho, \eta) = (-0.12, 0.34)$			
Channel	Amp	ReA$\times 10^3$	ImA$\times 10^3$
Sum Rules: A=B=C=D and E=F			
$A = \sqrt{2}\pi^0\pi^-$	t + p + e	3.03	9.94
	t + p	3.56	10.10
	t + p'	3.56	10.10
	t	3.56	10.10
$B = \sqrt{2}\pi^0\pi^0 - \pi^+\pi^-$	t + p + e	3.02	9.93
	t + p	3.54	10.09
	t+ p'	3.53	10.09
	t	3.57	10.10
$C = (\sqrt{2}\bar{K}^0\pi^0 - K^-\pi^+)/r_u$	t + p + e	14.10	12.01
	t + p	4.08	12.13
	t + p'	4.09	12.18
	t	4.30	12.20
$D = (\sqrt{2}K^-\pi^0 + \bar{K}^0\pi^-)/r_u$	t + p + e	14.13	12.03
	t + p	4.12	12.16
	t + p'	4.10	12.20
	t	4.30	12.17
$E = \pi^+\pi^- - K^-K^0$	t + p + e	-3.41	-8.32
	t + p	-3.55	-8.37
	t + p'	-3.62	-8.24
	t	-2.84	-8.05
$F = (K^-\pi^+ - \bar{K}^0\pi^-)/r_u$	t + p + e	-5.83	-9.84
	t + p	-3.27	-9.84
	t + p'	-3.26	-9.88
	t	-3.48	-9.89