B Physics Overview: Rare Hadronic and Radiative Decays

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ABSTRACT: I discuss three B physics results with potential for exhibiting new flavour-changing interactions: (1) the time-dependent CP asymmetry in $B \to J/\psi K_S$ decay; (2) branching fractions and CP asymmetries in $B \to \pi\pi, \pi K$ decays; (3) inclusive and exclusive semi-hadronic $b \to s$ transitions.

1. Introduction

The primary motivation for B physics is to complete the determination of the CKM matrix including its CP violating phase and to perhaps discover flavour-changing interactions beyond the standard weak interaction. With the large number of $B$ mesons now being produced at SLAC and KEK, these efforts have entered a new phase in which the standard theory will finally be seriously challenged. For this talk I have chosen to discuss the theory and interpretation of three sets of key observables, for which experimental results are already available, and are likely to be much improved in the near future.

2. Time-dependent CP asymmetry in $B \to J/\psi K_S$ decay

The time-dependent asymmetry in $B \to J/\psi K_S$ decay is predicted to be

$$A_{\text{mix}}(t) \equiv \frac{\Gamma(B^0(t) \to J/\psi K_S) - \Gamma(B^0(t) \to J/\psi K_S)}{\Gamma(B^0(t) \to J/\psi K_S) + \Gamma(B^0(t) \to J/\psi K_S)} = \sin \Phi_d \sin(\Delta M_{B_d} t),$$

i.e. it measures the phase of the $B\bar{B}$ mixing amplitude, $\Phi_d$, assuming the standard parameterization of the CKM matrix in which $V_{cb}V_{cs}^*$ has (almost) no phase. In the Standard

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Model \( \Phi_d = 2\beta \) (with \( \beta \) one of the angles of the unitarity triangle) is related to the phase of \( V_{td} \).

Eq. (2.1) is based on several assumptions, all of which are verified experimentally or expected to be valid to the percent level or better:

- The lifetime difference of the \( B_d \) mass eigenstates and CP violation in the \( B\bar{B} \) mixing amplitude is negligible.

- CP violation in \( K\bar{K} \) mixing is negligible.

- The decay proceeds only through amplitudes with a common weak phase, the phase of \( V_{cb}V^*_cs \).

The last assumption can be checked by fitting a \( \cos(\Delta M_{B_d}t) \) term to the time-dependence of \( A_{\text{mix}}(t) \), although the absence of such a term does not rigorously exclude a decay amplitude with a different weak phase. The above assumptions continue to hold in reasonable extensions of the standard model, so that the CP asymmetry \( A_{\text{mix}}(t) \) always measures the \( B\bar{B} \) mixing phase, but it may or may not equal \( 2\beta \).

The current results reported by BaBar \( (0.59 \pm 0.15 \ [3]) \) and Belle \( (0.99 \pm 0.15 \ [4]) \) together with an earlier CDF measurement \( [5] \) average to

\[
\sin(2\beta) = 0.79 \pm 0.10. \tag{2.2}
\]

This is in good agreement with an indirect determination of the \( B\bar{B} \) mixing phase that follows from the mass difference of the \( B_d \) and \( B_s \) eigenstates, CP violation in \( K\bar{K} \) mixing and \( |V_{ub}/V_{cb}| \): \( \sin(2\beta) = 0.68 \pm 0.21 \ [1], 0.68 \pm 0.15 \ [7] \) at 95% CL.

A large difference between the direct and indirect determination of the \( B\bar{B} \) mixing phase would probably have implied a new source of flavour violation beyond the CKM matrix. Because of this the good agreement is perhaps not very surprising, since models with non-CKM flavour violation are generally disfavoured by the smallness of CP violation in \( K\bar{K} \) mixing and the suppression of flavour-changing neutral currents. Eq. (2.2) therefore implies two important facts about Nature:

1) The Kobayashi-Maskawa mechanism \( [8] \) is most likely the dominant source of CP violation at the electroweak scale.

2) CP violating phases are large, i.e. CP is not an approximate symmetry.\(^2\)

\(^1\)This has been updated to account for new data published after the conference. At the time of the conference the result was \( 0.47 \pm 0.16 \).

\(^2\)In extensions of the Standard Model the observed amount of CP violation in the kaon system can be explained with small CP violating phases avoiding the requirement of a large CKM phase. This possibility might have been attractive from a conceptual point of view, since most extensions of the SM cause too much CP violation with generic CP violating phases. (However, the problem of large flavour-changing neutral currents would then have had to be solved by other means.)
3. Branching fractions and CP asymmetries in $B \to \pi\pi$, $\pi K$ decays

The decay $B \to J/\psi K$ with its unambiguous theoretical interpretation is unique. The generic situation is that $B$ decays which are interesting for constraining the unitarity triangle have significant decay amplitudes with different CP violating phases. This leads to new CP violating observables (“direct CP violation”) but to determine the CP phases one needs knowledge of the relative magnitudes of the different amplitudes – a strong interaction problem. There exists now a continuously increasing number of measured charmless $B$ decay modes, which are presumed to have two (nearly) comparable amplitudes, and which are sensitive to relative phases due to interference of the amplitudes. In the following I discuss methods to extract information on the phase of $V_{ub}$ from decays to pions and kaons. These final states have attracted interest recently, since the branching fractions for some decays into pions are observed to be somewhat lower than expected, while those into a pion and a kaon are larger than expected.

Two complementary strategies have been adopted to approach the strong interaction problem. The first approach begins with a general parameterisation of the strong interaction amplitudes based on SU(2) isospin symmetry. One then eliminates as many as possible strong interaction parameters through measurements. Additional, but less accurate constraints may follow from assuming SU(3) flavour symmetry. This approach is limited by the need of many accurate measurements and by SU(3) breaking effects. The second approach is based on a calculation of the strong interaction amplitudes. The feasibility of such calculations has only recently been understood and relies on the heavy quark limit. The limitation of this approach is that the $b$ quark mass is not that large.

3.1 Flavour symmetry

As is well-known the time-dependent CP asymmetry in $B \to \pi^+\pi^-$ decay determines $\gamma$, if the penguin amplitude is neglected. However, this is not a good approximation.

Neglecting only electroweak penguin amplitudes the isospin amplitude system for the three $\pi\pi$ final states contains five real strong interaction parameters, just as many as there are independent branching fractions under the same assumption. Solving this system allows one to determine $\gamma$ up to discrete ambiguities. Since this method requires a measurement of the small $B \to \pi^0\pi^0$ branching fractions, it has practical difficulties. Already bounds on the CP averaged $\pi^0\pi^0$ branching fraction can be useful to constrain the amplitude system. If $\text{Br}(\pi^0\pi^0)$ is small, the strong phase of the penguin-to-tree ratio cannot be large. In fact $\text{Br}(\pi^0\pi^0) = 0$ implies $\text{Br}(\pi^+\pi^-) = 2\text{Br}(\pi^\pm\pi^0)$. Conversely, a deviation from the last relation implies that $\text{Br}(\pi^0\pi^0)$ cannot be too small. At this moment the $\pi^\pm\pi^0$ measurement is still too uncertain to draw meaningful conclusions. Further constraints on the $\pi\pi$ modes can be obtained only by assuming also SU(3) or U-spin symmetry. This relates, for example, $B_d \to \pi^+\pi^-$ to $B_s \to K^+K^-$. The inverted CKM hierarchy of penguin and tree amplitude in the second decay can in principle be

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3I assume here that the $BB$ mixing phase is taken from experiment. Then $B \to \pi^+\pi^-$ determines the phase of $V_{ub}^*$, i.e. $\gamma$.

4The electroweak penguin amplitude is of order 1% and can be corrected for analytically.
used to determine $\gamma$ from a combined measurement of the time-dependent and direct CP asymmetries in both decays [15].

The $B \to \pi K$ decays are penguin-dominated, because the tree amplitudes are CKM suppressed. The final states $\pi^0 K^\pm$ and $\pi^0 K^0$ have significant electroweak penguin contributions. The amplitude system contains 11 real strong interaction parameters, too many to determine them all by measurements. Flavour symmetry is useful to constrain some of the amplitude parameters:

- Isospin symmetry implies [16]

$$\text{Br}(\pi^0 K^0) = \frac{\text{Br}(\pi^+ K^-) \text{Br}(\pi^- K^0)}{4 \text{Br}(\pi^0 K^-)} \times \{1 + O(\epsilon^2)\},$$

(3.1)

where $\epsilon \sim 0.3$ is related to the tree-to-penguin ratio. Unless the correction term is unexpectedly large this relation suggests a $\pi^0 K^0$ branching fraction of order $6 \times 10^{-6}$, about a factor $1.5 - 2$ smaller than the current measurements. Further isospin relations of this type have been derived [17].

- SU(3) or U-spin symmetry imply:

  a. The dominant electroweak penguin amplitude is determined [18].
  b. The magnitude of the tree amplitude for $I = 3/2$ final states is related to $\text{Br}(\pi^+ \pi^0)$.
  c. Rescattering and annihilation contributions to the (otherwise) pure penguin decay $B^+ \to \pi^+ K^0$ are constrained by $\text{Br}(K^+ K^0)$, where they are CKM enhanced relative to the penguin amplitude [13].

SU(3) flavour symmetry together with a few further dynamical assumptions (detailed below) suffice to derive bounds on $\gamma$ from CP averaged branching fractions alone. The inequality [20]

$$\sin^2 \gamma \leq \frac{\tau(B^+)}{\tau(B_d)} \frac{\text{Br}(\pi^+ K^-)}{\text{Br}(\pi^- K^0)} \equiv R$$

(3.2)

excludes $\gamma$ near 90° if $R < 1$ and is derived upon assuming that the rescattering contribution mentioned above and a colour-suppressed electroweak penguin amplitude are negligible. Current data give $R = 1.06 \pm 0.18$. The ratio of charged decay modes satisfies [18] (neglecting again the rescattering contribution to $B^+ \to \pi^+ K^0$)

$$2 \cdot \frac{\text{Br}(\pi^0 K^-)}{\text{Br}(\pi^- K^0)} \equiv R^{-1} \leq \left(1 + \tilde{\epsilon}_{3/2} |q - \cos \gamma| \right)^2 + \tilde{\epsilon}_{3/2}^2 \sin^2 \gamma,$$

(3.3)

where $q$ and $\tilde{\epsilon}_{3/2}$ are determined according to a. and b. above, respectively. This bound is particularly interesting, since, if $R^{-1} > 1$, it excludes a region in $\gamma$ around 55°, which is favoured by the indirect unitarity triangle constraints. Current data give $R^{-1} = 1.40 \pm 0.23$. This prefers $\gamma > 90°$, but the error is still too large to speculate about the implications.

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[5] Here and in the remainder of the text “Br” always refers to CP averaged branching fractions.
of this statement. Eq. (3.3) can be turned into a determination of $\gamma$ if one assumes that the strong phase of the tree amplitude relative to the penguin amplitude is not too large $^{21}$. This assumption is justified by theoretical calculations as discussed next, but will eventually be verified experimentally by the observation of small direct CP asymmetries.

### 3.2 QCD factorisation

QCD calculations of hadronic two-body decays deal with matrix elements

$$\langle M_1 M_2 | O_i | \bar{B} \rangle$$

of operators $O_i$, which appear in the weak effective Hamiltonian. A widely used approach has been to approximate the matrix element by the factorisation ansatz $^{22}$, but the validity of the approximation has never been clear, apart from its inadequacy to describe strong phases. We have shown recently $^{23, 24}$ that the arguments that lead to perturbative factorisation theorems for many strong interaction processes at large momentum transfer also imply that the matrix elements (3.4) factor into short-distance and long-distance parts. The short-distance scale is provided by the large mass of the decaying meson and the correspondingly large energy of the final state mesons. The long-distance parts are sufficiently simple to make this approach predictive. Schematically, the matrix element is computed as

$$F_{B \to M_1} \cdot T_{I}^{I} \ast \Phi_{M_2} + \Phi_{B} \ast T_{II}^{I} \ast \Phi_{M_1} \ast \Phi_{M_2}$$

in the heavy quark limit, i.e. neglecting corrections that scale as $1/m_b$. In this equation $F_{B \to M_1}$ denotes a form factor and $\Phi_{X}$ denote the meson’s leading-twist light cone distribution amplitudes, which are taken as non-perturbative inputs. The $T_{I,II}^{I}$ represent perturbative functions, which can be computed as series in $\alpha_s$, the strong coupling. A dependence on the distribution amplitudes arises only at order $\alpha_s$; at leading order the factorisation formula reproduces the earlier factorisation ansatz, but without phenomenological parameters. Strong interaction phases are generated by perturbative rescattering in the heavy quark limit and appear through imaginary parts of the hard-scattering functions $T_{I,II}^{I}$. For decays to pions and kaons these functions have been computed to next-to-leading order $^{21, 23, 25, 26}$. Since the factorisation formula holds only in the heavy quark limit one may expect non-negligible $1/m_b$ corrections. This is true in particular for strong interaction phases, which are either of order $\alpha_s$ or $\Lambda_{QCD}/m_b$. Some potentially large corrections have been estimated and have been included in a theoretical error estimate, but a general parametrisation of $1/m_b$ effects has neither been given nor is it likely to be useful in practice. It is therefore important to devise tests of the theory while simultaneously extracting information on CP violation.

The main results of our analysis $^{21}$ of $\pi\pi$ and $\pi K$ final states are briefly summarised as follows:

- The branching fractions for the modes $B^+ \to \pi^+\pi^0$ and $B^+ \to \pi^+\bar{K}^0$, which depend only on a single weak phase to very good approximation, are well described by the theory. This demonstrates that the magnitude of the tree and penguin amplitude
Figure 1: 95% (solid), 90% (dashed) and 68% (short-dashed) confidence level contours in the $(\bar{\rho}, \bar{\eta})$ plane obtained from a global fit to the CP averaged $B \to \pi K, \pi\pi$ branching fractions, using the scanning method [6]. The darker dot shows the overall best fit, whereas the light dot indicates the best fit for the default parameter set.

is obtained correctly. There is, however, a relatively large normalisation uncertainty for the $\pi K$ final states, which are sensitive to weak annihilation and the strange quark mass through the scalar penguin amplitude. This uncertainty can be partially eliminated by taking ratios of branching fractions.

- Branching fractions with a significant interference of tree and penguin amplitudes deviate slightly from their measured values [27, 28, 29], if $\gamma < 90^\circ$. As a consequence charmless $B$ decays appear to favour larger values of $\gamma$ than the standard unitarity triangle fit, although again the errors are too large to reach a definite conclusion.

- With some exceptions (such as the final state $\pi^0\pi^0$) strong phase differences are not large, so that QCD factorisation predicts small direct CP asymmetries, up to $(10 - 15)\%$ (with a preference around 5%) for the $\pi^0 K^\pm$, $\pi^\pm K^0$ final states. The current data seem to favour small CP asymmetries [28, 30, 31, 32], but they are not yet accurate enough to decide upon whether the QCD factorisation approach allows one to predict strong phases quantitatively, or whether it is limited to the qualitative statement that strong phases are small. Establishing direct CP violation in $B$ decays will not only constitute an important physics result but also provide a crucial test of our theoretical framework.

The current interpretation of charmless non-leptonic $B$ decays can be summarised by a global fit of the Wolfenstein parameters $\bar{\rho}, \bar{\eta}$ to the six CP averaged $\pi\pi$, $\pi K$ branching fractions shown in Figure 1. The constraint on $V_{ub}$ from semi-leptonic charmless $B$ decays is not used in this fit. The result is consistent with the conventional unitarity triangle fit [3] (shown as the shaded (yellow) region in Figure 1), but prefers somewhat larger $\gamma$ or smaller $|V_{ub}/V_{cb}|$. The quality of the fit is very good, independent of the phenomenological parameter $\rho_A$ which quantifies the size of weak annihilation. The darker (green) dot shows the value of $(\bar{\rho}, \bar{\eta})$ at which the $\chi^2$ function is minimized within the allowed ranges of theoretical input parameters. At this minimum $\chi^2$/d.o.f. $\approx 0.5$, $\gamma = 85^\circ$, $|V_{ub}/V_{cb}| = 0.071$. 


Table 1: Input and output of the global $(\bar{\rho}, \bar{\eta})$ fit. CP averaged branching fractions in units of $10^{-6}$.

| Decay Mode | Fit | Exp. Average |
|------------|-----|--------------|
| $B^0 \rightarrow \pi^+ \pi^-$ | 4.6 | 4.4 ± 0.9 |
| $B^{\pm} \rightarrow \pi^\mp \pi^0$ | 5.3 | 5.6 ± 1.5 |
| $B^0 \rightarrow \pi^+ K^\mp$ | 17.9 | 17.2 ± 1.5 |
| $B^{\pm} \rightarrow \pi^0 K^\mp$ | 11.3 | 12.1 ± 1.7 |
| $B^{\pm} \rightarrow \pi^\pm K^0$ | 17.7 | 17.2 ± 2.5 |
| $B^0 \rightarrow \pi^0 K^0$ | 7.1 | 10.3 ± 2.5 |

The corresponding theory input parameters and output branching fractions are given in Table 1.

4. Inclusive and exclusive semi-hadronic $b \rightarrow s$ transitions

The flavour-changing neutral current transition $b \rightarrow s\gamma$ has been constraining extensions of the standard model severely, since it was first observed by CLEO [34]. The decay amplitude is not only loop-suppressed (a property in common with $B^+ \rightarrow \pi^+ K^0$ discussed above), but also suppressed due to the chiral weak interactions. In the effective low-energy theory the transition occurs through the operator $g\bar{s}\sigma^{\mu\nu}(1 + \gamma_5)bF_{\mu\nu}$. The left-handedness of weak interactions implies that the coefficient scales with $m_b/M_W^2$, so that the operator is effectively dimension six. New, non-chiral, flavour interactions can replace $m_b$ by a weak scale quantity, elevating the operator to dimension five.

Because the leading contribution is loop-induced, the computation of the first correction in renormalisation group improved perturbation theory is a difficult task, which however has been completed [35, 36, 37]. Recently, it has been suggested to parametrise

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The results discussed here disagree with another analysis [33] that has prematurely declared a failure of QCD factorisation. This overly pessimistic assessment of the current situation is based on an incomplete use of the QCD factorisation input (neglecting in particular radiative corrections to the scalar penguin contribution) together with neglecting errors in some of the important theory input parameters. A phenomenological parameter (similar to $\rho_A$ above) for the QCD penguin amplitude is then introduced as a further fit parameter. Despite this, due to the rigid theory input parameters, the quality of the fit is in fact worse than the fit based on QCD factorisation.
the dominant correction in terms of the charm quark $\overline{MS}$ mass rather than the pole mass \[38\]. This increases the theoretical prediction by about 10% to

$$\text{Br}(B \to X_s \gamma |_{E_\gamma > m_b/20} = (3.73 \pm 0.30) \cdot 10^{-4}$$

and suggests that the theory is perhaps not yet as precise as commonly assumed. In view of this the prediction seems to be consistent with the experimental average \[39, 40, 41\]

$$(2.96 \pm 0.35) \cdot 10^{-4}$$. The experimental error includes a theoretical error from extrapolating the photon energy spectrum to small photon energies, which are not measured.

The exclusive transitions $B \to K^* \gamma$ are easier to measure, but more dependent on hadronic parameters. One obstacle to a precise prediction of the exclusive decay has recently been removed \[42, 43\] by demonstrating that non-factorisable strong interaction effects can be computed in perturbation theory in the heavy quark limit. The exclusive mode can then be treated by renormalisation group methods in a manner analogous to the inclusive mode and is now also known to next-to-leading-logarithmic order. The theoretical prediction is proportional to the $B \to K^*$ tensor form factor $T_1(0)$, and is given by

$$\text{Br}(\overline{B} \to \overline{K}^* \gamma) = (7.3 \pm 1.4) \cdot 10^{-5} \times \left(\frac{\tau_B}{1.6 \text{ps}}\right) \left(\frac{m_b(m_b)}{4.2 \text{GeV}}\right)^2 \left(\frac{T_1(0)}{0.38}\right)^2 .$$

The next-to-leading order correction is about 30% on the amplitude level, enhancing the decay rate significantly. This enhancement is not unexpected since a large part of the correction is identical to the equally large NLO correction to the inclusive rate. The theoretical prediction is now significantly higher than the experimental results $\text{Br}(\overline{B}^0 \to \overline{K}^{*0} \gamma)_{\text{exp}} = (4.54 \pm 0.37) \cdot 10^{-5}$ and $\text{Br}(B^- \to \overline{K}^{*-\gamma})_{\text{exp}} = (3.81 \pm 0.68) \cdot 10^{-5}$ \[44, 45, 40\]. Given the agreement in the inclusive sector a non-Standard Model explanation of this difference appears unlikely since it would have to be connected with spectator specific interactions. The difference must then be blamed on a sizeable $1/m_b$ correction to the amplitude or a smaller tensor form factor. Isospin breaking effects are absent at leading order in $1/m_b$ in the decay amplitude, so that the branching fractions for charged and neutral $B$ mesons may differ only through the different $B$ meson lifetimes in this approximation. If the trend indicated by the experimental result is real, it may signal a gross failure of the theory or a new isospin-violating interaction. In both cases this effect should be seen more clearly in $B \to \rho \gamma$ decay.

Requiring the inclusive branching fraction to be consistent with data in the Standard Model extended by a second Higgs doublet implies that the charged Higgs boson cannot be light \[46, 47\]. The same is true in the supersymmetric standard model unless chargino exchange interferes destructively with the charged Higgs contribution. This constrains the model parameter space, requiring in particular that the product $A_t \mu$ of the soft-supersymmetry breaking parameter $A_t$ and the $\mu$-parameter be negative. For models with large $\tan \beta$ the theoretical prediction becomes unreliable, unless corrections enhanced by powers of $\tan \beta$ are resummed. This resummation strengthens the constraints on the parameter space \[48, 49\]. The discussion of charged Higgs boson and chargino effects in supersymmetry usually proceeds under the assumption that flavour-changing gluino
vertices are negligible. Once such couplings are allowed, implying non-diagonal squark mass matrices in a standard quark mass eigenstate basis, they constitute generically the dominant supersymmetric effect, so that the data severely limits the magnitude of these additional flavour-violating interactions \[50\]. Including all supersymmetric effects simultaneously opens the possibility for cancellations \[51\], but the general conclusion remains true that the allowed squark mass matrix pattern is rather restricted.

The transition \(b \rightarrow s\ell^+\ell^-\) opens new possibilities to probe new flavour interactions, since it is sensitive to flavour-changing \(Z\) boson couplings \(\bar{s}\Gamma bZ\) \[52\]. At low energies the corresponding dimension-six operators involve two Wilson coefficients, \(C_9\) and \(C_{10}\). Both can be determined from the lepton invariant mass spectrum and the forward-backward asymmetry. The forward-backward asymmetry is a particularly interesting quantity \[53\], since the value \(q^2_0\) at which the asymmetry vanishes for the exclusive decay \(\bar{B} \rightarrow \bar{K}^*\ell^+\ell^-\) provides a relation between the coefficient of the magnetic penguin operator \(C_7\) and \(C_9\) which is nearly free of hadronic uncertainties \[54\]:

\[
C_9 + \text{Re}(Y(q^2_0)) = -\frac{2M_B m_b}{q^2_0} C^\text{eff}_7, \tag{4.3}
\]

where \(Y\) is a calculable function. The next-to-leading order corrections to the exclusive decay have recently been calculated \[42\] (using the new result for the two-loop virtual correction to the inclusive decay \[55\]) and have been shown to correct Eq. (4.3) significantly, as seen from the shift of the asymmetry zero in Figure 2. However, the conclusion that \(C_9\) can be determined from the asymmetry zero with little uncertainty remains valid after this correction is taken into account.

5. Conclusion

The three sets of observables discussed here could easily have come out different from their Standard Model expectation, but apparently did not. The challenge is therefore to understand what distinguishes the Kobayashi-Maskawa mechanism of CP violation and the
GIM mechanism of suppressing flavour-changing neutral currents from the many sources of CP violation and flavour-changing neutral currents that arise in generic extensions of the Standard Model.

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