MORE PRECISE DETERMINATION of $V_{cb}$ & $V_{ub}$ and
DIRECT CP VIOLATION IN CHARMLESS B DECAYS

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More precise extractions on the two important CKM matrix elements $|V_{cb}|$ and $|V_{ub}|$ are presented
up to $1/m_Q^2$ corrections based on the complete heavy quark effective field theory (HQEFT) of QCD.
The HQEFT as a large component QCD provides us a complete theoretical framework for correctly
and systematically evaluating the subleading and higher order contributions in $1/m_Q^2$ expansion of
heavy quarks. A global analysis on charmless B decays is made based on the isospin and SU(3) flavor
symmetry. An isospin relation is found to be very useful for studying SU(3) symmetry breaking
effects of strong phases and exploring new type of electroweak penguin effects. The direct CP
violation in charmless B decays is predicted and its precise measurement is helpful either for testing
SU(3) symmetry breaking effects of strong phases or for probing new physics beyond standard model.

Introduction. The standard model (SM) has been well tested in the gauge sector. The unsolved and unclear
problems in the SM are mainly concerned in the Yukawa sector which is strongly related to flavor physics, such as:
origin and mechanism of CP violation, origins of the quark and lepton masses as well as their mixing. It involves
thirteen parameters whose origins are all unknown. Therefore, precisely extracting those parameters and testing CP
violation mechanism as well as probing new physics become hot topics in flavor physics. In fact, flavor physics has
already indicated the existence of new physics.\[1\]. Much efforts have been made by experimentalists in flavor physics.\[2\].
For instance, after forty years of discovery for indirect CP violation, direct CP violation in kaon decays has recently
been established by two experimental groups\[3, 4\], which is also consistent with the theoretical predictions\[7\].
CP violation has also been observed in B decays at the two B-factories. More precise experimental results in flavor physics
will become available in the recent years. It is known that exclusive semileptonic and inclusive B decays play a crucial
role for extracting two important parameters $V_{cb}$ and $V_{ub}$ in the CKM matrix elements. Rare B decays and direct
CP violations are also of great importance in determining weak phase angles of the unitarity triangle and testing the
Kobayashi-Maskawa (KM) mechanism \[6\] in SM as well as probing new physics. In this talk, we will pay attention
to the more precise extraction of the CKM matrix elements $V_{cb}$ and $V_{ub}$ and the direct CP violation in charmless B
decays.

Complete HQEFT of QCD and incompleteness of the usual HQET. Heavy quark effective field theory
(HQET) of QCD, which was first explored in \[7\] and recently developed in detail by a series of papers \[8, 10, 11, 12, 13, 14, 15, 16, 17\],
provides a promising and systematic tool in correctly evaluating the hadronic matrix elements of heavy quarks and precisely extracting the CKM matrix elements $V_{cb}$ and $V_{ub}$ from B decays via heavy quark expansion (HQE). As the HQEFT is a theoretical framework derived directly from QCD, it explicitly displays
the heavy quark symmetry (HQS) \[11\] in the infinite mass limit $m_Q \rightarrow \infty$ \[21\] and symmetry breaking corrections
for finite mass case in the real world. In fact, the HQEFT has been shown to be as a large component QCD \[7, 17\].
At the leading order, it coincides with the usual heavy quark effective theory (HQET) \[21\] which is constructed based on the heavy quark symmetry in the infinite mass limit. The differences between HQEFT of QCD and the usual
HQET arise from the subleading terms in the $1/m_Q^2$ expansion. This is because in the construction of HQET the
class and antiparticle components were separately treated based on the assumption that the particle number and
antiparticle number are conserved separately in the effective Lagrangian. However, such an assumption is only valid in
the infinite mass limit. Note that the particle number and antiparticle number are always conserved in the transition
matrix though the particle and antiparticle number is not conserved in the Lagrangian due to the pair creation and
annihilation interaction terms, which is independent of heavy quark limit and is in fact the basic principle of quantum
field theory. Obviously, the quark-antiquark coupled terms that correspond to the pair creation and annihilation interaction terms in full QCD were inappropriately dropped away in the usual HQET. Those terms have been shown in HQEFT of QCD to be suppressed by $1/m_Q$ and they truly become vanishing in infinite mass limit. Thus the usual HQET based on the assumption of particle and antiparticle number conservation in the effective Lagrangian cannot be regarded as a complete effective theory for evaluating the subleading and higher order corrections, it must be an incomplete effective theory. Unlike the derivation of the usual HQET from QCD by making the assumption of particle and antiparticle number conservation in the effective Lagrangian [22], we have derived the HQEFT [7] from full QCD by carefully treating all contributions of the field components, i.e., large and small, ‘particle’ and ‘antiparticle’ in the effective Lagrangian, so that the resulting effective Lagrangian should form the right basis for a complete effective field theory of heavy quarks [15]. Though the quark-antiquark mixing terms are suppressed by $1/m_Q$, their physics effects at subleading and higher orders have been found in some cases to be significant and crucial for obtaining consistent results [8, 9, 10, 11, 12, 13, 14, 15, 16]. Therefore, to correctly and precisely consider the finite quark mass corrections, it is necessary to include the contributions from the components of the antiquark fields. Namely, consistent results [8, 9, 10, 11, 12, 13, 14, 15, 16]. Therefore, to correctly and precisely consider the finite quark mass corrections, it is necessary to include the contributions from the components of the antiquark fields. Namely, consistent results [8, 9, 10, 11, 12, 13, 14, 15, 16].

More Precise Extraction of $|V_{cb}|$. Exclusive semileptonic decays $B \to D^+(D)\nu$ provide one of the main approaches to extract $|V_{cb}|$. The differential decay rates are:

$$\frac{d\Gamma(B \to D^+\nu)}{d\omega} = \frac{G_F^2}{48\pi^3} (m_B - m_{D^+})^3 m_{D^+}^3 \omega^2 - 1(\omega + 1)^2 \times \left[ (1 + \frac{4\omega}{\omega + 1}) \frac{m_B^2 - 2\omega m_B m_{D^+} + m_{D^+}^2}{(m_B - m_{D^+})^2} \right] |V_{cb}|^2 \mathcal{F}(\omega),$$

$$\frac{d\Gamma(B \to D\nu)}{d\omega} = \frac{G_F^2}{48\pi^3} (m_B + m_D)^3 m_D^3 (\omega^2 - 1)^{3/2} |V_{cb}|^2 \mathcal{G}(\omega)$$

with

$$\mathcal{F}(1) = \eta_A h_A(1) = \eta_A(1 + \delta^*), \quad \mathcal{G}(1) = \eta_V [h_+(1) - \frac{m_B - m_D}{m_B + m_D} h_-(1)] = \eta_V (1 + \delta),$$

where the QCD radiative corrections to two loops give the short distance coefficients $\eta_A = 0.960 \pm 0.007$ and $\eta_V = 1.022 \pm 0.004$ [24]. $\omega$ is the product of the four-velocities of the $B$ and $D^+(D)$ mesons, $\omega = v \cdot v'$. The weak transition form factors $h_A(\omega)$, $h_+(\omega)$ and $h_-(\omega)$ can be expanded in powers of $1/m_Q$ and represented by the heavy quark spin-flavor independent wave functions. Up to order $1/m_Q^2$ and neglecting the contributions of operators containing two gluon field strength tensors, one has at the zero recoil point $\omega = 1$

$$h_{A_1} = 1 + \frac{1}{8\Lambda} \left[ \frac{1}{m_b} (\kappa_1 + 3\kappa_2) - \frac{1}{m_c} (\kappa_1 - \kappa_2) \right]^2 - \frac{1}{8m_b^2 \Lambda} (F_1 + 3F_2 - 2\Lambda_1 + 6\Lambda_2)$$

$$+ \frac{1}{8m_c^2 \Lambda} (F_1 - F_2 - 2\Lambda_1 + 2\Lambda_2) + \frac{1}{4m_b m_c \Lambda} (F_1 + F_2 - 2\Lambda_1 - 2\Lambda_2),$$

$$h_+ = 1 + \frac{1}{8\Lambda} \left[ \frac{1}{m_b} - \frac{1}{m_c} \right]^2 [(\kappa_1 + 3\kappa_2)^2 - (F_1 + 3F_2) + 2(\Lambda_1 + 3\Lambda_2)], \quad h_- = 0,$$

where the parameters in the rhs. of Eqs. (4)-(5) are defined in HQEFT [8]. The binding energy of a heavy meson $M$ is defined as $\Lambda \equiv \lim_{m_q \to \infty} \Lambda_M = \lim_{m_q \to \infty} (m_M - m_Q)$. It is seen from Eqs. (3) and (4) that both $h_{A_1}$ and $h_+$ in HQEFT of QCD do not receive $1/m_Q$ order correction at the zero recoil point. More interestingly, the vanishing value of $h_-$ in HQEFT arises from the fact of partial cancellation between the $1/m_Q$ correction in the current expansion.
and the $1/m_Q$ correction coming from the insertion of the effective Lagrangian into the transition matrix elements. Such a cancellation is not observed in the incomplete HQET because in the latter framework the quark-antiquark couplings are not taken into account explicitly. Though the HQET protects the weak transition matrix elements from $1/m_Q$ order correction at zero recoil point, it does not protect $h_-$ from such correction. As a result, in the incomplete HQET framework, only the $B \to D^*\ell\nu$ decay rate at zero recoil is strictly protected against $1/m_Q$ order correction when $B \to D\ell\nu$ decay rate is not. This is the main reason besides the experimental considerations for the conclusion that the $B \to D\ell\nu$ decay is not as favorable as $B \to D^*\ell\nu$ decay for $|V_{cb}|$ extraction in HQET. While the complete HQEFT of QCD is more reliable to extract $|V_{cb}|$ from both $B \to D^*\ell\nu$ and $B \to D\ell\nu$ decays, because both rates of these decays do not receive $1/m_Q$ order correction, as can be seen from Eqs. \textbf{11-15}.

Furthermore, the HQEFT of QCD gives interesting relations between meson masses and wave functions, which enable one to get the zero recoil values of HQEFT wave functions straightforwardly from the heavy meson mass spectrum. On the other hand, the wave functions can also be calculated via other nonperturbative approaches such as sum rules. As these parameters are estimated consistently, $|V_{cb}|$ can be determined to a good accuracy.

**Inclusive semileptonic $B$ decays** are the other alternative to determine $|V_{cb}|$. In the usual HQET the light quark in a hadron is generally treated as a spectator, which does not affect the heavy hadron properties to a large extent. This treatment may be one possible reason for the failure of HQET in some applications. For example, the world average value for bottom hadron lifetime $\tau(B) = \tau(B^0)$ can not be explained well in the usual framework of HQET. Instead of simply applying the equation of motion for infinitely heavy free quark, $i\nu \cdot DQ^+_V = 0$, we treat the heavy quark in a hadron as a dressed particle, which means that the residual momentum $k$ of the heavy quark within a hadron is considered to comprise contributions from the light degrees of freedom. This simple picture is adopted to conveniently take into account the effects of light degrees of freedom and the binding effects of heavy and light components of the hadron but not deal with the complex dynamics of hadronization directly. Explicitly, one may use the relation $(i\nu \cdot D) \equiv \langle B |i\nu \cdot DQ^+_V |H_s\rangle/2\mathcal{M}_H \approx \mathcal{K} \neq 0$. To acquire a good convergence of HQE, we perform the expansion in terms of $k - v(i\nu \cdot D)$ (or say, equivalently, in terms of $1/(m_Q + \mathcal{K})$). Then the $B \to X_c\ell\nu$ decay rate is found to be \textbf{19, 11}:

$$\Gamma(B \to X_c\ell\nu) = \frac{G_F^2 \tilde{m}_b^2 V_{cb}^2}{192\pi^3} \eta_{cl}(\rho, \rho, \mu) \left[ I_0(\rho, \rho, \hat{\rho}) + I_1(\rho, \rho, \hat{\rho}) \frac{\kappa_1}{3\tilde{m}_0^2} - I_2(\rho, \rho, \hat{\rho}) \frac{\kappa_2}{\tilde{m}_0^2} \right],$$

(6)

where $\tilde{m}_b = m_b + \mathcal{K}$, $I_0$, $I_1$ and $I_2$ are functions of the mass square ratios $\rho = m_0^2/\tilde{m}_b^2$, $\rho^2 = \hat{m}_0^2/\tilde{m}_b^2$ and $\rho^2 = m_0^2/\hat{m}_0^2$. Here the calculation is performed up to nonperturbative order $1/\tilde{m}_0^2$ and perturbative order $\alpha_s^2$. The function $\eta_{cl}$ characterizes QCD radiative corrections. $\kappa_2$ is often extracted from the known $B - B^*$ mass splitting $\kappa_2 \approx \frac{1}{8}(m_{B^*}^2 - m_B^2) \approx 0.06 \text{GeV}^2$ which is consistent with the sum rule result. There are several points to be mentioned for Eq. \textbf{6}. Firstly, in deriving Eq. \textbf{6}, the effects of light degrees of freedom are explicitly accounted for in the picture of a dressed heavy quark in a hadron. Secondly, it is seen that the next leading order contributions vanish in our HQE in terms of the inverse dressed heavy quark mass, $1/\tilde{m}_b$. Furthermore, our HQE in terms of $k - v(i\nu \cdot D)$ (or $1/\tilde{m}_b$) has a good convergence. It is found that the $1/\tilde{m}_0^2$ order contributions induce only $-0.7 \approx 5\%$ corrections to the total width $\Gamma(B \to X_c\ell\nu)$. Therefore we conclude that the higher order nonperturbative corrections can be safely neglected. Finally, now one needs only to treat the dressed quark mass $\tilde{m}_b = m_b + \mathcal{K}$ instead of considering the uncertainties arising from the two quantities $m_b$ and $\mathcal{K}$ separately. Note that these features can not be observed in the HQE in the usual HQET, where one assumes $(i\nu \cdot D)$ to be zero or of higher order of $1/m_b$. In HQET the next leading order corrections can be absent only when the HQE is performed in terms of $1/m_b$, and the heavy quark mass $m_b$ and the binding energy $\mathcal{K}$ have to be treated separately. Thus in HQET the theoretical prediction of the total decay width strongly depends on the value of bottom quark mass $m_b$.

**Numerical Values of $V_{cb}$**. Using the current world average $|V_{cb}|F(1) = 0.0383 \pm 0.0005 \pm 0.0009$, $|V_{cb}|G(1) = 0.0413 \pm 0.0029 \pm 0.0027$, the lifetime $\tau(B^0) = 1.540 \pm 0.014 \text{ps}$ and the branching ratio from CLEO $\text{Br}(B \to X_c\ell\nu) = (10.49 \pm 0.17 \pm 0.43)\%$, we have \textbf{18}:

$$|V_{cb}| = 0.0395 \pm 0.0011_{\text{exp}} \pm 0.0019_{\text{th}} \quad \text{from} \quad B \to D^*\ell\nu, \quad O(1/m_Q^2)$$
$$|V_{cb}| = 0.0434 \pm 0.0041_{\text{exp}} \pm 0.0020_{\text{th}} \quad \text{from} \quad B \to D\ell\nu, \quad O(1/m_Q^2)$$
$$|V_{cb}| = 0.0394 \pm 0.0010_{\text{exp}} \pm 0.0014_{\text{th}} \quad \text{from} \quad B \to X_c\ell\nu, \quad O(1/m_Q^2)$$

(7)

where the result extracted from $B \to D\ell\nu$ decay receives a larger experimental uncertainty than that from $B \to D^*\ell\nu$ decay but a similar theoretical uncertainty as the latter. The result obtained from $B \to D^*\ell\nu$ decay agrees quite well with that from inclusive $B \to X_c\ell\nu$ decay. These results then give the average \textbf{18}:

$$|V_{cb}| = 0.0402 \pm 0.0014_{\text{exp}} \pm 0.0017_{\text{th}}$$.

(8)
Alternatively it can be represented as $A = 0.83 \pm 0.07$ in the Wolfenstein parameterization [21] $|V_{cb}| = A \lambda^2$ with $\lambda = |V_{ud}| = 0.22$.

More precise extraction of $|V_{ub}|$. It is similar to that of $|V_{cb}|$. The main difference is now one generally has to deal with heavy-to-light decays. For exclusive decays $B \to \pi(p)\ell\nu$, one can parameterize the leading order transition matrix elements in HQEFT as [32, 33]

$$
\langle \pi(p) | \Gamma Q^+ | B \rangle = -Tr[\pi(v,p)\Gamma M(v)], \quad \langle \rho(p, e^*) | \Gamma Q^+ | B \rangle = -iTr[\Omega(v,p)\Gamma M_v],
$$

where

$$
\pi(v, p) = \gamma^\mu_A \langle v | A(p, v, \mu) | B \rangle,
\Omega(v, p) = L_1(v \cdot p)\hat{\phi} + L_2(v \cdot p)(v \cdot e^*) + [L_3(v \cdot p)\hat{\phi}^* + L_4(v \cdot p)(v \cdot e^*)] \hat{\phi}
$$

with $\hat{\phi} = \frac{p^\mu}{v^\mu}$. A, B and $L_i (i = 1, 2, 3, 4)$ are the leading order wave functions in HQEFT. HQS and relevant effective theories are useful for studying heavy-to-light decays as they give us relations between different channels. For example, since $L_i$ are heavy quark mass independent, $B \to \rho \ell\nu$ and $D \to \rho \ell\nu$ are characterized by the same set of wave functions $L_i$. In this sense, HQS and effective theories simplify the heavy-to-light decays, though for a single channel the number of independent functions is not reduced. We calculated these wave functions from light cone sum rules, considering the $\pi$ distribution functions up to twist 4 and the $\rho$ distribution functions up to twist 2 for $B \to \pi \ell\nu$ decay [12] and for $B \to \rho \ell\nu$ decay [13]. Recently, the calculation on $B \to \pi \ell\nu$ up to $1/m_Q$ order has also been performed and the finite mass correction in HQEFT is found to be small. Inclusive decays $B \to ul\nu$ can be investigated in a similar way as for $B \to el\nu$ decays by using the HQE in HQEFT.

Numerical Values of $|V_{ub}|$. Our numerical results are summarized as follows

$$
|V_{ub}|_{LO} = (3.4 \pm 0.5\exp \pm 0.5_{th}) \times 10^{-3} \quad \text{from } B \to \pi\ell\nu,
|V_{ub}|_{LO} = (3.7 \pm 0.6\exp \pm 0.7_{th}) \times 10^{-3} \quad \text{from } B \to \rho\ell\nu,
|V_{ub}|_{NLO} = (3.2 \pm 0.5\exp \pm 0.2_{th}) \times 10^{-3} \quad \text{from } B \to \pi\ell\nu,
|V_{ub}| = (3.48 \pm 0.62\exp \pm 0.11_{th}) \times 10^{-3} \quad \text{from } B \to X_u\ell\nu \quad O(1/m_Q^2)
$$

All these results can be compared with the average of CLEO [30]: $|V_{ub}| = (3.25^{+0.25}_{-0.32} \pm 0.55) \times 10^{-3}$.

Summary I. Within the framework of complete HQEFT of QCD, various processes and approaches all lead to consistent and more precise values for $|V_{cb}|$ and $|V_{ub}|$.

Charmless B decays and direct CP violation. With the successful running of B factories, high precision data on the rare hadronic B decay modes such as $B \to \pi\pi, \pi K$ [34, 35, 36, 37, 38, 39] have been obtained, which provide us good opportunities to extract the weak phase angle $\gamma$, to test the theoretical approaches for evaluating the hadronic transition matrix elements, and to explore new physics beyond the SM.

The recently proposed methods such as QCD Factorization [40, 41] and pQCD approach [42, 43] have been extensively discussed. From those methods, useful information of weak phase angles such as $\gamma$ can be extracted [44, 45].

On the other hand, methods based on flavor isospin and SU(3) symmetries are still helpful and important [46, 47, 48, 49, 50]. The advantage of this kind of approaches is obvious that they are model independent and more convenient in studying the interference between weak and strong phases. Recently the flavor isospin and SU(3) symmetries in charmless B decays are studied by using global fits to the experiment data [51, 52, 53]. In a general isospin decomposition, there exist a lot of independent free parameters. But using the flavor isospin and SU(3) symmetries, the number of parameters can be greatly reduced and the method of global fit becomes applicable. Through direct fit, the isospin or SU(3) invariant amplitudes as well as the corresponding strong phases can be extracted with a reasonable precision. Our early results [51] have already indicated some unexpected large isospin amplitudes and strong phases. The fitted amplitudes and strong phases can also provide useful information for the weak phase $\gamma$ [52]. However unlike isospin symmetry, the flavor SU(3) symmetry is known to be broken down sizably [53, 54]. The ways of introducing SU(3) breaking may have significant influence on the final results. In the usual considerations, the main effects of SU(3) breaking are often taken into accounted only in the amplitudes. To be more general, the study of SU(3) breaking including strong phases is necessary, which can be significant and bring a consistent explanation to the present data [55]. Of particular, their effects can lead to different predictions on direct CP violation [56].

Isospin Relation. Let us take decays $B \to \pi\pi$ as an example, the final states of $\pi\pi$ have isospin of 2 and 0. The isospin amplitudes $A_2$ and $A_0$ are defined as follows

$$
A_2 \equiv \langle \pi\pi, I = 2 | H_{eff}^{1/2} | B \rangle = \lambda_0 a_2^0 e^{i\delta_0} + \lambda_0 a_2^0 e^{i\delta_0},
A_0 \equiv \langle \pi\pi, I = 0 | H_{eff}^{1/2} | B \rangle = \lambda_0 a_0^0 e^{i\delta_0} + \lambda_0 a_0^0 e^{i\delta_0},
$$

(11)
where \( a_i^q, (q = u, c \text{ and } I = 2, 0) \) are the amplitudes associated with \( \lambda_q = V_{q3}V^*_{q4} \). The isospin structure of effective Hamiltonian leads to the following relations between two isospin amplitudes

\[
\frac{a_2^u}{a_2^c} = R_{EW} = \frac{3}{2} \cdot \frac{C_9 + C_{10}}{C_1 + C_2 + C_9 + C_{10}}.
\]

(12)

Taking the Wilson coefficients at \( \mu = m_b \), one has \( C_1 = 1.144, C_2 = -0.308, C_9 = -1.28\alpha \) and \( C_{10} = 0.328\alpha \). Thus

\[
R_{EW} = -1.25 \times 10^{-2}, \quad \text{and} \quad \delta_2^u = \delta_2^c.
\]

(13)

Note that the relation is obtained without the knowledge of the matrix element \( \langle I = 2 | O(3/2) | 0 \rangle \). It can not be affected by the final state inelastic rescattering processes with lower isospin as it is only related to the highest isospin component \( \Delta I = 3/2 \). Furthermore, the value of \( R_{EW} \) is the ratio between the electroweak penguin and tree diagrams. It is then sensitive to new physics effects beyond the SM in electroweak penguin sector. The new physics effects on \( R_{EW} \) have been discussed in Refs. \[57, 58, 59, 60\]. It seems quite sensitive to several new physics models. A precise determination of \( R_{EW} \) from experiments may be helpful to single out possible new physics or study flavor symmetry breaking in charmless \( B \) decays. To describe the possibility that the value of \( R_{EW} \) extracted from experiments could be different from the SM calculations, we introduce a factor \( \kappa \) as follows

\[
R_{EW}^{\text{exp}} = \kappa \cdot R_{EW} \simeq -0.0125 \cdot \kappa, \quad k = 1 \text{ in SM}
\]

(14)

where \( R_{EW}^{\text{exp}} \) stands for its value extracted from experiments and obviously \( \kappa = 1 \) in SM.

**SU(3) Analysis.** In flavor SU(3) limit, the decay amplitudes for \( B \to \pi \pi \) and \( B \to \pi K \) is directly connected

\[
a_0^u e^{i\delta_0^u} = a_{1/2}^u e^{i\delta_{1/2}^u}, \quad a_0^c e^{i\delta_0^c} = a_{1/2}^c e^{i\delta_{1/2}^c}, \quad a_2^u e^{i\delta_2^u} = a_{3/2}^u e^{i\delta_{3/2}^u}, \quad a_2^c e^{i\delta_2^c} = a_{3/2}^c e^{i\delta_{3/2}^c}.
\]

(15)

If these relations are adopted, the number of free parameters is reduced to be nine. From Eq. (12) and the above relation, one finds that \( \frac{a_{3/2}^u}{a_{3/2}^c} = \frac{a_2^u}{a_2^c} = R_{EW} \). Thus the highest isospin amplitudes for the \( B \to \pi K \) decays satisfy the same relation as the one in the \( B \to \pi \pi \) decay. When SU(3) breaking effects are considered, the above relations have to be modified. At present stage, it is not very clear how to describe the SU(3) breaking effects. A widely used approach is introducing a breaking factor \( \xi \) which characterizes the ratio between \( B \to \pi K \) and \( \pi \pi \) decay amplitudes, i.e., \( a_{1/2}^u = \xi a_0^u, \quad a_{3/2}^u = \xi a_2^u, \quad \text{but strong phases are assumed to remain satisfying the SU(3) relations}

\[
\delta_{1/2}^u = \delta_0^u, \quad \delta_{3/2}^u = \delta_2^u.
\]

(16)

Typically \( \xi = f_K/f_\pi \simeq 1.23 \) with \( f_\pi \) and \( f_K \) being the pion and kaon meson decay constants, which comes from the naive factorization calculations. It is easy to see that this pattern of SU(3) breaking is a quite special one. The value of \( \xi \) is highly model dependent. It can only serve as an order of magnitude estimation and it is even not clear whether a single factor can be applied to all the isospin amplitudes. The equal strong phase assumption implies that the SU(3) breaking effects on strong phase are all ignored, which may be far away from the reality. In a more general case, all the strong phases could be different when SU(3) is broken down. The breaking effects on strong phases may have significant effects on the prediction for the direct CP violations in those decay modes.

**SU(3) Symmetry Breaking Effects of Strong Phases.** To describe the possible violations of relations in eq. (10) or the SU(3) breaking effects on strong phases, we may introduce the following phase differences \( \Delta_0^q(q = u, c \text{ and } I = 3/2, 1/2) \):

\[
\delta_0^q = \delta_{1/2}^q + \Delta_{1/2}^q, \quad \delta_2^q = \delta_{3/2}^q + \Delta_{3/2}^q \quad (q = u, c).
\]

(17)

On the other hand, the SU(3) breaking effects in amplitudes may also be given in a more general way

\[
a_{1/2}^q = \xi^q a_0^q, \quad a_{3/2}^q = \xi^q a_2^q \quad (q = u, c)
\]

(18)

The SU(3) limit corresponds to the case that all \( \Delta_0^q \) vanish and \( \xi^q = 1 \).

However, as it can be seen in table 3, the global fit to the latest data show that for \( \xi = 1 \) or 1.23 the best fitted value of \( \kappa \) is quite large, around 10. The fits with different values of \( \gamma \) show that the such a large \( \kappa \) insensitive to \( \gamma \). This result is closely related to the observed large branching ration of \( B \to \pi^0 K^0 \) and \( B \to \pi^0 \pi^0 \). The value of \( \kappa \) is sensitive to the contributions from electroweak penguin diagrams. Since many new physics models can give significant corrections to this sector, it may be helpful to study new physics effects on \( \kappa \). However, to explore any new physics effects and arrive at a definitive conclusion for the existence of new physics from the hadronic decays, it is necessary
to check all the theoretical assumptions and make the most general considerations. It is noted that the above results are obtained by assuming SU(3) symmetry with its breaking only in amplitudes. Therefore, we need first examine the above results to a more general case of SU(3) symmetry breaking before claiming any possible new physics signals. Within the framework of the SM, we found that the breaking effects of flavor SU(3) symmetry could be considerable.

We then explore the parameter space of those SU(3) breaking factors and found that for some typical non-zero values of the phase shift $\Delta_1$ the best fit of $\kappa$ does restore to it’s values in SM, i.e. close to unity. For example, in cases of $\Delta_{1/2}^u = +\pi/6$, $\Delta_2^c = +\pi/6, +\pi/3$ and $\Delta_{3/2}^c = +\pi/6, +\pi/3$, the best fitted values of $\kappa$ are around 1.5 with the minimal $\chi^2_{min}$ $\leq$ 4. The results is similar for other values of $\gamma$ as long as $\gamma < 120^\circ$.

**Direct CP Violation.** From the fit result, the corresponding direct CP violation can also be obtained. The best fitted direct CP violation for example, in case (c) is given by

$$A_{CP}(\pi^+ \pi^-) \simeq 0.3, \quad A_{CP}(\pi^0 \pi^0) \simeq 0.4$$

$$A_{CP}(\pi^+ K^-) \simeq -0.1, \quad A_{CP}(\pi^0 \bar{K}^0) \simeq -0.1$$

$$A_{CP}(\pi^0 K^-) \simeq -0.0, \quad A_{CP}(\pi^- \bar{K}^0) \simeq 0.1$$

In cases of $\Delta_{1/2}^u = +\pi/6$, $\Delta_{2/3}^c = +\pi/6, +\pi/3$ and $\Delta_{3/2}^c = +\pi/6, +\pi/3$, the best fitted values of $\kappa$ are around 1.5 with the minimal $\chi^2_{min}$ $\leq$ 4. The direct CP violation for $\Delta_{1/2}^u = +\pi/6(\Delta_{2/3}^c = +\pi/6)$ is as follows,

$$A_{CP}(\pi^+ \pi^-) \simeq 0.1(0.5), \quad A_{CP}(\pi^0 \pi^0) \simeq 0.5(0.2),$$

$$A_{CP}(\pi^+ K^-) \simeq -0.1(-0.1), \quad A_{CP}(\pi^0 \bar{K}^0) \simeq -0.2(-0.1),$$

$$A_{CP}(\pi^0 K^-) \simeq -0.1(-0.0), \quad A_{CP}(\pi^- \bar{K}^0) \simeq 0.1(0.1).$$

Compared with the ones with SU(3) symmetry case, the predicted values of direct CP violation can be quite different.

**Summary II.** In the case of SU(3) limits and also the case with SU(3) breaking only in amplitudes, the fitting results lead to an unexpected large ratio between two isospin amplitudes $a_{3/2}^c/a_{3/2}^u$, which is about an order of magnitude larger than the SM prediction. The results are found to be insensitive to the weak phase $\gamma$. By including SU(3) breaking effects on the strong phases, one is able to obtain a consistent fit to the current data within the SM, which implies that the SU(3) breaking effect on strong phases may play an important role in understanding the observed charmless hadronic B decay modes $B \to \pi\pi$ and $\pi K$. It is possible to test those breaking effects in the near future from more precise measurements of direct CP violation in B factories.

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[1] H. Fritzsch, theory summary talk in this proceedings.
