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New bounds on the Cabibbo-Kobayashi-Maskawa matrix from $B \to K\pi\pi$ Dalitz plot analyses

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We present a new technique to extract information on the unitarity triangle from the study of $B \to K\pi\pi$ Dalitz plots. Using the sensitivity of Dalitz analyses to the absolute values and the phases of decay amplitudes and isospin symmetry, we obtain a new constraint on the elements of the CKM matrix. We discuss in detail the role of electroweak penguin contributions and outline future prospects.

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I. INTRODUCTION

The study of flavor and CP violation in B decays allows us to test the flavor structure of the standard model (SM) and to look for new physics (NP). In the last few years, B factories have provided us with a large amount of new data in this field, in particular, the first measurements of the angles α and γ of the unitarity triangle (UT). These measurements are in agreement with the indirect determinations of the UT angles [1,2] and therefore provide an important test of the standard model and stringent constraints on NP [3].

$B \to K\pi$ decays could in principle constitute another source of information on γ [4,5]. However, the fact that the tree-level $b \to u$ transition, carrying the phase γ, is doubly Cabibbo suppressed in these channels, together with a large dynamical enhancement of the Cabibbo allowed penguin contribution [6], make a model-independent extraction of γ from $B \to K\pi$ impossible. Model-dependent studies of these channels have tackled this issue with no success [7,8].

In this letter, we consider the possibility of obtaining a tree-level determination of γ from $B \to K^*\pi$ decays, cancelling the effect of penguins through the rich set of information available from Dalitz plot analyses. In particular, the B factories can provide the magnitude and phase of decay amplitudes separately for B and $\bar{B}$ decays. One can think of exploiting this information, together with isospin symmetry, to build combinations of amplitudes that are proportional to a single weak phase.

This very simple idea, however, in the case of $B \to K\pi\pi$, has to face the presence of Electroweak Penguins (EWP), doubly Cabibbo enhanced with respect to current-current operators in $b \to s$ transitions. This enhancement can largely compensate the $\alpha_{\rm em}$ suppression, leading to an $\mathcal{O}(1)$ correction to the decay amplitude to the $I = 3/2$ final state. This fact, however, does not spoil the possibility of extracting information on the UT with small hadronic uncertainty, as explained below.

Dalitz plot analyses combined with isospin have already shown their effectiveness in the extraction of α from $B \to \pi\pi\pi$ [9]. A proposal relying on isospin for extracting γ using a global analysis of $B \to K_S(\pi\pi)_{I=0,2}$, including time-dependent CP asymmetries at fixed values of the Mandelstam variables, can be found in Ref. [10]. We focus instead on $K^*\pi$ final states which allow us to perform a simpler analysis with no need of time-dependent measurements.

II. EXTRACTING CKM MATRIX ELEMENTS FROM $B \to K\pi\pi$ DALITZ PLOTS

Let us first illustrate our idea for the simplified case in which we neglect EWP contributions. To this aim, we write the amplitudes using isospin symmetry, in terms of Renormalization Group Invariant (RGI) complex parameters [11], obtaining

$$\begin{aligned} A(K^{*+}\pi^-) &= V^*_{tb}V_{ts}P_1 - V^*_{ub}V_{us}(E_1 - P_1^{\rm GIM}) \\ \sqrt{2}A(K^{*0}\pi^0) &= -V^*_{tb}V_{ts}P_1 - V^*_{ub}V_{us}(E_2 + P_1^{\rm GIM}) \\ \sqrt{2}A(K^{*+}\pi^0) &= V^*_{tb}V_{ts}P_1 - V^*_{ub}V_{us}(E_1 + E_2 + A_1 - P_1^{\rm GIM}) \\ A(K^{*0}\pi^+) &= -V^*_{tb}V_{ts}P_1 + V^*_{ub}V_{us}(A_1 - P_1^{\rm GIM}), \end{aligned} \qquad (1)$$

where $P_1^{\rm (GIM)}$ represent (GIM-suppressed) penguin contributions, A_1 the disconnected annihilation and E_1 (E_2) the connected (disconnected) emission topologies. $\bar{B}$ decay amplitudes are simply obtained by conjugating the CKM factors V_{ij}. Similar expressions hold for higher K^* resonances.

Considering the two combinations of amplitudes

$$A^0 = A(K^{*+}\pi^-) + \sqrt{2}A(K^{*0}\pi^0) = -V^*_{ub}V_{us}(E_1 + E_2), \qquad (2)$$

$$\bar{A}^0 = A(K^{*-}\pi^+) + \sqrt{2}A(\bar{K}^{*0}\pi^0) = -V_{ub}V^*_{us}(E_1 + E_2), \qquad (3)$$

the ratio

$$R^0 = \frac{\bar{A}^0}{A^0} = \frac{V_{ub}V^*_{us}}{V^*_{ub}V_{us}} = e^{-2i\gamma} \qquad (4)$$

provides a clean determination of the weak phase γ. We now discuss how to extract A_0 and $\bar{A}_0$. Looking at the decay chains $B^0 \to K^{*+}(\to K^+\pi^0)\pi^-$ and $B^0 \to K^{*0}(\to K^+\pi^-)\pi^0$, one can obtain A^0 from the $K^+\pi^-\pi^0$ Dalitz plot, including the phase in a given convention, for example $\mathrm{Im}A(K^{*+}\pi^-) = 0$. Similarly, $\bar{A}^0$ can be extracted from the $K^-\pi^+\pi^0$ Dalitz plot using the same procedure, choosing $\mathrm{Im}A(K^{*-}\pi^+) = 0$. However, in general, this choice does not reproduce the physical phase difference between $A(K^{*+}\pi^-)$ and $A(K^{*-}\pi^+)$, which has to be fixed using additional information.

This information can be provided by the $K_S\pi^+\pi^-$ Dalitz plot, considering the decay chain $B^0 \to K^{*+}(\to K^0\pi^+)\pi^-$ and the CP conjugate $\bar{B}^0 \to K^{*-}(\to \bar{K}^0\pi^-)\pi^+$. These two decay channels do not interfere directly on the Dalitz plot, but they both interfere with the decays $B, \bar{B} \to \rho^0(\to \pi^+\pi^-)K_S$ and with other resonances contributing to the same Dalitz plot. Therefore the Dalitz analysis of B, $\bar{B} \to K_S\pi^+\pi^-$ should include the $\rho^0 K_S$ final state. In a time-integrated analysis, the $\rho^0 K_S$ final state comes from a mixture of B and $\bar{B}$ mesons, while the $K^{*+(-)}\pi^{-(+)}$ final state only originates from B ($\bar{B}$) decay. Looking at the phases of $A(K^{*+(-)}\pi^{-(+)})$ relative to $A(\rho^0 K_S)$, we can extract the phase difference between $A(K^{*+}\pi^-)$ and $A(K^{*-}\pi^+)$ with no need of resolving the flavor of the B in $A(\rho^0 K_S)$ [12].

A similar isospin relation involves charged B decays. We have

$$A^+ = A(K^{*0}\pi^+) + \sqrt{2}A(K^{*+}\pi^0) = -V_{ub}^* V_{us}(E_1 + E_2), \tag{5}$$

$$A^- = A(\bar{K}^{*0}\pi^-) + \sqrt{2}A(K^{*-}\pi^0) = -V_{ub} V_{us}^*(E_1 + E_2), \tag{6}$$

and the ratio

$$R^{\mp} = \frac{A^-}{A^+} = e^{-2i\gamma}. \tag{7}$$

As before, $A^\pm$ can be extracted from the decay chains $B^\pm \to K^{*\pm}(\to K^0\pi^\pm)\pi^0$ and $B^\pm \to K^{*0}(\to K^0\pi^0)\pi^\pm$ entering the $K_S\pi^\pm\pi^0$ Dalitz plot. Electric charge forbids the extraction of the relative phase of the two Dalitz plots along the way discussed above, so that a strategy based on theoretical arguments has to be adopted. In particular, one can follow two possible paths.

The first one is to use isospin symmetry to relate charged and neutral B decays:

$$A(K^{*+}\pi^-) + \sqrt{2}A(K^{*0}\pi^0) - A(K^{*0}\pi^+) - \sqrt{2}A(K^{*+}\pi^0) = 0 \tag{8}$$

$$A(K^{*-}\pi^+) + \sqrt{2}A(\bar{K}^{*0}\pi^0) - A(\bar{K}^{*0}\pi^-) - \sqrt{2}A(K^{*-}\pi^0) = 0 \tag{9}$$

so that the relative phases can be fixed. In this way, no additional information on γ can be extracted from $R^{\mp}$. However, the full information coming from charged and neutral B's can be combined to improve the accuracy of the determination of γ, thanks to the increase of available statistics. The second possibility is to use the penguin-dominated channel $K^{*0}\pi^+$ to fix the phase difference between the amplitudes in the two Dalitz plots. In this way an independent, albeit more uncertain, determination of γ can be obtained from $R^{\mp}$. To illustrate this point, let us write down the phase for the $K^{*0}\pi^\pm$ final state

$$\arg(A(K^{*0}\pi^+)) = \beta_s + \arg\left(1 + \frac{V_{ub}^* V_{us}}{V_{tb}^* V_{ts}} \Delta^+ e^{i\delta_{\Delta^+}}\right)$$

$$\arg(A(\bar{K}^{*0}\pi^-)) = -\beta_s + \arg\left(1 + \frac{V_{ub} V_{us}^*}{V_{tb} V_{ts}^*} \Delta^+ e^{i\delta_{\Delta^+}}\right)$$

where $\beta_s = \arg(-V_{ts}V_{tb}^*/(V_{cs}V_{cb}^*))$ and

$$\Delta^+ e^{i\delta_{\Delta^+}} = \frac{A_1 - P_1^{\mathrm{GIM}}}{P_1}. \tag{10}$$

We now take advantage of the fact that $|V_{ub}^* V_{us}|/|V_{tb}^* V_{ts}| \ll 1$ to simplify the above equations and we obtain

$$\arg(A(\bar{K}^{*0}\pi^-)) = \arg(A(K^{*0}\pi^+)) - 2\beta_s + 2\Delta^+ \mathrm{Im}\frac{V_{ub}V_{us}^*}{V_{tb}V_{ts}^*}\cos\delta_{\Delta^+}. \tag{11}$$

On general grounds, we expect $\Delta^+ \sim \mathcal{O}(1)$. The error induced by the last term in Eq. (11) can be estimated at the level of $|V_{ub}V_{us}^*|/|V_{tb}V_{ts}^*| \sim \lambda^2$. The determination of γ from $R^{\mp}$ is not as theoretically clean as the one obtained from R^0. Nevertheless, the uncertainty induced by our dynamical assumption, being of $\mathcal{O}(\lambda^2)$, is much smaller than the expected experimental error (at least in the near future).

III. INCLUSION OF ELECTROWEAK PENGUINS

The inclusion of the effect of EWP's completely changes Eqs. (4) and (7). In fact, even though EPW's give a subdominant contribution to branching ratios (because of the $\mathcal{O}(\alpha_{\mathrm{em}})$ suppression with respect to the strong penguin contribution), they provide an $\mathcal{O}(1)$ contribution to R^0 and $R^{\mp}$ (more generally, they provide an $\mathcal{O}(1)$ correction to CP violating effects in charmless $b \to s$ decays). Fortunately, as we shall discuss in the following, the dominant EWP's (i.e. left-handed EWP operators) can be eliminated at the operator level, so that no additional hadronic matrix elements are introduced. The net effect of EWP's is that R^0 and $R^{\mp}$ depend not only on γ but also on other CKM parameters.

Let us consider the effective Hamiltonian for $b \to s$ transitions given for instance in Eq. (5) of Ref. [11]. There is a hierarchy in the values of the Wilson coefficients for EWP operators: $|C_{9,10}| \gg |C_{7,8}|$. Let us therefore neglect the effect of $Q_{7,8}$ and focus on $Q_{9,10}$, barring a very large dynamical enhancement of $\langle Q_{7,8}\rangle$ in the case of B decays. There is an exact operator relation that allows to eliminate $Q_{9,10}$ [13] in favor of current-current and penguin operators [14]:

$$Q_9 = \frac{3}{2}(Q_2^{\rm suu} - Q_2^{\rm scc}) + 3Q_2^{\rm scc} - \frac{1}{2}Q_3^s$$
$$Q_{10} = \frac{3}{2}(Q_1^{\rm suu} - Q_1^{\rm scc}) + 3Q_1^{\rm scc} - \frac{1}{2}Q_4^s \qquad (12)$$

so that the effective Hamiltonian becomes

$$H_{\rm eff} = \frac{G_F}{\sqrt{2}}\bigg[\bigg(V_{ub}^*V_{us}C_+ - \frac{3}{2}V_{tb}^*V_{ts}C_+^{\rm EW}\bigg)(Q_+^{\rm suu} - Q_+^{\rm scc}) + \bigg(V_{ub}^*V_{us}C_- + \frac{3}{2}V_{tb}^*V_{ts}C_-^{\rm EW}\bigg)(Q_-^{\rm suu} - Q_-^{\rm scc}) - V_{tb}^*V_{ts}H^{\Delta I=0}\bigg]. \qquad (13)$$

where $Q_\pm = (Q_1 \pm Q_2)/2$, $C_\pm = C_1 \pm C_2$ and $C_\pm^{\rm EW} = C_9 \pm C_{10}$. We observe that the relation

$$\frac{C_+^{\rm EW}}{C_+} = \frac{C_-^{\rm EW}}{C_-} \qquad (14)$$

is exact at the LO and is broken by $\mathcal{O}(\alpha_s\alpha_{\rm em}\log)$ corrections. Numerically, Eq. (14) holds with high accuracy, being violated at the percent level. Using this relation, we can write

$$H_{\rm eff} \simeq \frac{G_F}{\sqrt{2}}\bigg\{V_{ub}^*V_{us}(1+\kappa_{\rm EW})\bigg[C_+(Q_+^{\rm suu} - Q_+^{\rm scc}) + \frac{1-\kappa_{\rm EW}}{1+\kappa_{\rm EW}}C_-(Q_-^{\rm suu} - Q_-^{\rm scc})\bigg] - V_{tb}^*V_{ts}H^{\Delta I=0}\bigg\}, \qquad (15)$$

where

$$\kappa_{\rm EW} \equiv -\frac{3}{2}\frac{C_+^{\rm EW}}{C_+}\frac{V_{tb}^*V_{ts}}{V_{ub}^*V_{us}} = \frac{3}{2}\frac{C_+^{\rm EW}}{C_+}\bigg(1 + \frac{1-\lambda^2}{\lambda^2(\bar\rho + i\bar\eta)} + \mathcal{O}(\lambda^2)\bigg). \qquad (16)$$

Therefore R^0 and $R^\mp$ can be rewritten as

$$R^{0,\mp} = e^{-2i(\gamma + \arg(1+\kappa_{\rm EW}))}\frac{1 + \frac{1-\kappa_{\rm EW}^*}{1+\kappa_{\rm EW}^*}\frac{C_-}{C_+}re^{i\theta_r}}{1 + \frac{1-\kappa_{\rm EW}}{1+\kappa_{\rm EW}}\frac{C_-}{C_+}re^{i\theta_r}}, \qquad (17)$$

where

$$re^{i\theta_r} = \frac{\langle K^*\pi(I=3/2)|Q_-|B\rangle}{\langle K^*\pi(I=3/2)|Q_+|B\rangle}. \qquad (18)$$

While for $B \to K\pi$ decays the $SU(3)$ flavor symmetry guarantees that $\langle K\pi(I=3/2)|Q_-|B\rangle$ vanishes [5,15], the same argument, based on the symmetry property of the final state wave function, does not apply to $K^*\pi$ final states [16]. However, using factorized amplitudes and form factors as given in Ref. [7], one obtains

$$r = \bigg|\frac{f_{K^*}F_0^{B\to\pi} - f_\pi A_0^{B\to K^*}}{f_{K^*}F_0^{B\to\pi} + f_\pi A_0^{B\to K^*}}\bigg| \lesssim 0.05. \qquad (19)$$

While this numerical result depends on the estimate of the form factors, the good agreement between QCD sum rules and lattice QCD calculations makes it rather robust. In our analysis, however, we do not assume a specific model for computing the amplitudes, rather we let r to vary in the conservative range 0–0.3.

Concerning $\kappa_{\rm EW}$, using the values $C_+(m_b) = 0.877$, $C_+^{\rm EW}(m_b) = -1.017\alpha_{\rm em}$ and $\bar\rho = 0.216$, $\bar\eta = 0.342$ [1], we obtain $\kappa_{\rm EW} = -0.35 + 0.53i$. One can thus verify that $\kappa_{\rm EW}$ is an $\mathcal{O}(1)$ correction to the decay amplitude to the $I = 3/2$ final state.

The above equations allow to translate the experimental results for R^0 and $R^\mp$ into allowed regions in the $\bar\rho$-$\bar\eta$ plane. Neglecting terms of $\mathcal{O}(r^2)$, for a given value of R^0 one obtains the linear relation $\bar\eta = -\tan(\frac{1}{2}\arg R^0)\times(\bar\rho - \bar\rho_0)$, where

$$\bar\rho_0 = -\bigg[\frac{3C_+^{\rm EW}}{2C_+ + 3C_+^{\rm EW}} - \frac{12C_+^{\rm EW}C_- r\cos\theta_r}{(2C_+ + 3C_+^{\rm EW})^2}\bigg]\cdot\frac{1-\lambda^2}{\lambda^2} + \mathcal{O}(\lambda^2) \qquad (20)$$

Including terms of $\mathcal{O}(r^2)$ or higher, the relation between $\bar\eta$ and $\bar\rho$ becomes quadratic, but the deviation from the result in Eq. (20) is small and mainly amounts to a depletion of the region $\bar\eta \sim 0$. Furthermore, the effect of r is constrained by the measurement of $|R^{0,\mp}|$, since $|R^{0,\mp}| - 1 \propto r\sin\theta_r$.

We can test this new idea using the experimental result of the *BABAR* Collaboration on $B^0 \to K^+\pi^-\pi^0$ decays [17]. Among the other measurements, this analysis provides the decay amplitudes for B^0 and $\bar B^0$ decays to $K^*\pi$ and to $K^*(1430)\pi$ final states. We implement this information in our analysis using directly the shape of the multidimensional likelihood from *BABAR* (including all correlations). In this way, we obtain an error of 38° on $\arg R^0$. Unfortunately, the $K_S\pi^+\pi^-$ Dalitz plot is not yet available so that we cannot fix the relative phase of B and $\bar B$ decays at present. For the sake of illustration, we assume a central value for this relative phase such that the constraint on $\bar\rho$ and $\bar\eta$ from Eq. (17) is compatible with the SM UT fit result [1]. For the experimental uncertainty on the relative phase, we consider two cases, corresponding to ±20° or ±40°, leading to the constraints exhibited in Fig. 1, obtained without expanding in r.

The situation can be improved by fitting $R^0_{K^*\pi}$ and $R^0_{K^*(1430)\pi}$ directly from data, cancelling out part of the systematic error. In addition, the determination of the UT parameters can be further improved with the experimental measurement of $R^\mp_{K^*\pi}$ and $R^\mp_{K^*(1430)\pi}$ and adding Belle data.

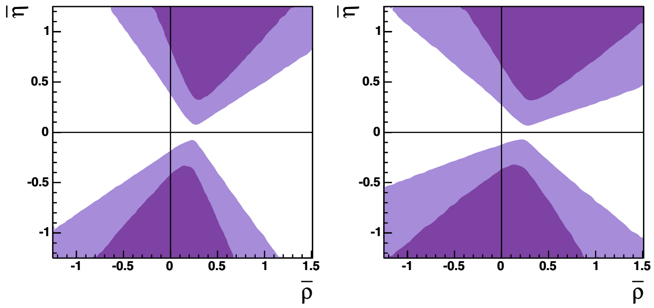

FIG. 1 (color online). Bounds in the $\bar{\rho}$-$\bar{\eta}$ plane from the analysis of B^0 decays, assuming a measurement of the relative phase with an error of 20° (left) or 40° (right). The output of the present UT fit analysis [1] is shown as a reference. Dark (light) regions correspond to 68% (95%) probability.

Let us finally comment on the sensitivity to NP of our analysis. Making the very reasonable assumption that NP effects only enter at the loop level, we can envisage three possibilities. First of all, NP could affect the coefficients of QCD penguin operators. In this case, the analysis of R^0 is completely unaffected, while the phase of NP contributions would modify Eq. (11). This could produce a discrepancy between the constraints on the UT obtained from R^0 and $R^{\mp}$ using Eq. (11). A second possibility is that NP modifies EWP coefficients, respecting however the hierarchy $C_{9,10} \gg C_{7,8}$. In this case, the only effect would be a modification of $\kappa_{\rm EW}$, so that the constraint on the UT obtained using the SM value for $\kappa_{\rm EW}$ could be inconsistent with the SM UT fit result. Finally, NP could produce contributions to EWP operators such that $C_{9,10} \sim C_{7,8}$, or give rise to new $\Delta I = 3/2$ operators that cannot be eliminated. In this case, one would observe $|R^{0,\mp}| \neq 1$. Present data give $|R^0| = 0.96 \pm 0.17$. A small $|R^{0,\mp}| - 1$ could also be generated by $\langle Q_- \rangle$ or by a large dynamical enhancement of $\langle Q_{7,8} \rangle$ within the SM.

IV. CONCLUSIONS AND OUTLOOK

We have presented a new method to constrain the unitarity triangle using $B \to K\pi\pi$ decays. This can be achieved with both neutral and charged B decays, using the amplitude ratios R^0 and $R^{\mp}$ defined in Eqs. (4) and (7). The theoretical uncertainty is negligible with respect to the foreseen experimental error. We have discussed in detail how to take into account electroweak penguins. Our exploratory study shows that this new constraint can have a sizable impact on the unitarity triangle analysis in the near future. We have discussed possible improvements and sensitivity to New Physics.

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