CP violation in $\Lambda \to p\pi^-$ beyond the standard model

Xiao-Gang He
Institute of Theoretical Science, University of Oregon, Eugene, Oregon 97403

G. Valencia
Department of Physics, Iowa State University, Ames, Iowa 50011

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The CP-violating asymmetry $A(\Lambda^0_-)$ has been estimated to occur at the level of a few times of 10^{-5} within the minimal standard model. The experiment E871 expects to reach a sensitivity of 10^{-4} to the asymmetry $A(\Lambda^0_-) + A(\Xi^-_-)$. In this paper we study some of the implications of such a measurement for CP violation beyond the minimal standard model. We find that it is possible to have $A(\Lambda^0_-)$ at the few times 10^{-4} level while satisfying the constraints imposed by the measurements of CP violation in kaon decays.

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I. INTRODUCTION

The origin of CP violation remains one of the outstanding problems in particle physics. In the attempt to understand this problem many experimental and theoretical efforts have been launched [1]. One of the systems where it is possible to search for CP violation is the nonleptonic decay of hyperons. Although this has been known for many years [2], it is only recently that it has become conceivable to carry out an experimental program to look for CP-violating signals in the decays of Ξ and Λ hyperons [3,4].

Of particular interest is the upcoming experiment E871 that expects to reach a sensitivity of 10^{-4} for the sum of asymmetries $A(\Lambda^0_-) + A(\Xi^-_-)$ [4]. Unfortunately, the calculation of these asymmetries is plagued by theoretical uncertainties in the estimate of the hadronic matrix elements involved. Nevertheless, a conservative study of these asymmetries within the minimal standard model indicated that $A(\Lambda^0_-)$ is likely to occur at the level of a few times 10^{-5}. In view of this, the potential results of E871 are very exciting.

One of the questions we would like to answer is whether the phase in the Cabibbo-Kobayashi-Maskawa (CKM) matrix of the three-generation minimal standard model is the sole source of CP violation. The experimental information that we have so far is the following.

A nonzero value of the parameter ϵ in kaon decays [5]:

$$|\epsilon| = 2.26 \times 10^{-3} \ . \tag{1}$$

A measurement of the parameter ϵ' [6]:

$$\frac{\epsilon'}{\epsilon} = \begin{cases} (2.3 \pm 0.65) \times 10^{-3} \ \ \text{NA31} \ , \\ (0.74 \pm 0.52 \pm 0.29) \times 10^{-3} \ \text{E731} \ . \end{cases} \tag{2}$$

The first result indicates that there is CP violation in nature, but it does not pinpoint its origin. The best one can say is that it is possible for the minimal standard model to accommodate this number. If the second number turns out to be nonzero, it would establish the existence of direct $|\Delta S| = 1$ CP violation, ruling out some superweak models. The current experimental numbers are consistent with the minimal standard model, although the theoretical calculations are also plagued with uncertainty from the evaluation of hadronic matrix elements.

The present situation is, therefore, that there is no need for CP violation beyond the phase in the three-generation CKM matrix,[1] but that other sources of CP violation have not been ruled out.

The question we want to address in this paper is whether it is possible for E871 to find a nonzero asymmetry given its expected sensitivity and the current values of ϵ and ϵ'/ϵ. To this end and in keeping with the results of all the precise experiments conducted to date, we will assume that the minimal three-generation standard model is a very good low energy approximation to the electroweak interactions. We will, therefore, discuss any possible new physics in terms of an effective Lagrangian consistent with the symmetries of the standard model and will only look at operators of dimension six.

Our paper is organized as follows. In Sec. II we review the notation for CP-violating observables in hyperon decays as well as the standard model estimate of $A(\Lambda^0_-)$. In Sec. III we compute the contributions of CP-violating four-quark operators to $A(\Lambda^0_-)$ and the constraints that result from the measurements of CP violation in $K \to \pi\pi$. In Sec. IV we repeat this analysis for the two-quark operators of dimension six (so called penguin operators). Finally, we present our conclusions.

II. CP VIOLATION IN $\Lambda^0 \to p\pi^-$

In this section we review the basic features of CP violation in the reaction $\Lambda^0 \to p\pi^-$, denoted by (Λ^0_-). In

[1] Except perhaps in the origin of the baryon asymmetry of the universe. We will not discuss that issue in this paper.

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the Λ^0 rest frame, $\vec{\omega}_{i,f}$ will denote unit vectors in the directions of the Λ and p polarizations, and $\vec{q}$ will denote the proton momentum. The isospin of the final state is $I = 1/2$ or $3/2$, and each of these two states can be reached via a $\Delta I = 1/2$ or $3/2$ weak transition, respectively. There are also two possibilities for the parity of the final state. They are the s-wave, $l = 0$, parity-odd state (thus reached via a parity-violating amplitude), and the p-wave, $l = 1$, parity-even state reached via a parity-conserving amplitude.

A model-independent analysis of the decay can be done by writing the most general matrix element consistent with Lorentz invariance [2]:

$$\mathcal{M} = G_F m_\pi^2 \bar{u}_P (A - B\gamma_5) u_\Lambda \; . \tag{3}$$

It is customary to introduce the quantities

$$s = A \; ,$$

$$p = \left(\frac{|\vec{q}|}{E_P + M_P}\right) B \tag{4}$$

to write the total decay

$$\Gamma = \frac{|\vec{q}|(E_P + M_P)}{4\pi M_\Lambda} G_F^2 m_\pi^4 (|s|^2 + |p|^2) \; . \tag{5}$$

The angular distribution is proportional to

$$\frac{d\Gamma}{d\Omega} \sim 1 + \gamma \vec{\omega}_i \cdot \vec{\omega}_f + (1-\gamma)\hat{q}\cdot\vec{\omega}_i \hat{q}\cdot\vec{\omega}_f + \alpha \hat{q}\cdot(\vec{\omega}_i + \vec{\omega}_f) + \beta \hat{q}\cdot(\vec{\omega}_f \times \vec{\omega}_i) \; , \tag{6}$$

where we have used the standard notation [2]:

$$\alpha \equiv \frac{2\mathrm{Re}s^* p}{|s|^2 + |p|^2}, \quad \beta \equiv \frac{2\mathrm{Im}s^* p}{|s|^2 + |p|^2}, \quad \gamma \equiv \frac{|s|^2 - |p|^2}{|s|^2 + |p|^2} \; . \tag{7}$$

If the proton polarization is not observed, α is the parameter that governs the angular distribution:

$$\frac{d\Gamma}{d\Omega} = \frac{\Gamma}{4\pi}(1 + \alpha \hat{q}\cdot\vec{\omega}_i) \; . \tag{8}$$

Similarly, if the initial Λ is unpolarized, α determines the polarization of the proton:

$$\vec{\mathcal{P}}_p = \alpha_\Lambda \hat{q} \; . \tag{9}$$

E871 will not measure the correlation governed by the parameter β, and so we will not deal with it in this paper.

The CP-odd observable $A(\Lambda^0_-)$ is constructed by comparing the parameter α in the reaction $\Lambda^0 \to p\pi^-$ with the corresponding parameter $\bar{\alpha}$ in the reaction $\bar{\Lambda}^0 \to \bar{p}\pi^+$. One can show that CP symmetry predicts that

$$\bar{\alpha} = -\alpha \; , \tag{10}$$

so that a CP-odd observable is [7]

$$A \equiv \frac{\alpha\Gamma + \bar{\alpha}\bar{\Gamma}}{\alpha\Gamma - \bar{\alpha}\bar{\Gamma}} \approx \frac{\alpha + \bar{\alpha}}{\alpha - \bar{\alpha}} \; . \tag{11}$$

Other possible CP-odd observables have been discussed in the literature: a rate asymmetry that is significantly smaller than A [7] and an asymmetry based on the parameter β that will not be accessible to E871. For this reasons we concern ourselves with the observable[2] $A(\Lambda^0_-)$.

It is convenient to decompose the amplitudes according to isospin and to introduce the following notation for the phases:

$$\begin{aligned} s(\Lambda^0_-) &= -\sqrt{2/3} s_1 e^{i(\delta_1^s + \phi_1^s)} + \sqrt{1/3} s_3 e^{i(\delta_3^s + \phi_3^s)} \; , \\ p(\Lambda^0_-) &= -\sqrt{2/3} p_1 e^{i(\delta_1^p + \phi_1^p)} + \sqrt{1/3} p_3 e^{i(\delta_3^p + \phi_3^p)} \; , \end{aligned} \tag{12}$$

where δ_J^I is the strong rescattering phase for the pion nucleon system and ϕ_J^I is the CP-violating phase.

In terms of these quantities, one finds [7]

$$\begin{aligned} A(\Lambda^0_-) = -\tan(\delta_1^p - \delta_1^s)\sin(\phi_1^p - \phi_1^s)\Bigg[1 &+ \frac{1}{\sqrt{2}}\frac{s_3}{s_1}\left(\frac{\cos(\delta_1^p - \delta_3^s)}{\cos(\delta_1^p - \delta_1^s)} - \frac{\sin(\delta_1^p - \delta_3^s)\sin(\phi_1^p - \phi_3^s)}{\sin(\delta_1^p - \delta_1^s)\sin(\phi_1^p - \phi_1^s)}\right) \\ &+ \frac{1}{\sqrt{2}}\frac{p_3}{p_1}\left(\frac{\cos(\delta_3^p - \delta_1^s)}{\cos(\delta_1^p - \delta_1^s)} - \frac{\sin(\delta_3^p - \delta_1^s)\sin(\phi_3^p - \phi_1^s)}{\sin(\delta_1^p - \delta_1^s)\sin(\phi_1^p - \phi_1^s)}\right)\Bigg] \; . \end{aligned} \tag{13}$$

Experimentally we know the values of the strong rescattering phases [9],

$$\delta_1^s \approx 6.0^\circ, \quad \delta_3^s \approx -3.8^\circ, \quad \delta_1^p \approx -1.1^\circ, \quad \delta_3^p \approx -0.7^\circ \; , \tag{14}$$

with all the errors on the order of 1°, the $\Delta I = 3/2$ amplitudes are much smaller than the $\Delta I = 1/2$ amplitudes [10],

$$s_3/s_1 = 0.027 \pm 0.008, \quad p_3/p_1 = 0.03 \pm 0.037 \; , \tag{15}$$

and the s and p amplitudes (assuming they are dominated by the CP-conserving, $\Delta I = 1/2$, transitions)

[2] In fact, E871 will be sensitive to the sum $A(\Lambda^0_-) + A(\Xi^-_-)$. An analysis of $A(\Xi^-_-)$ parallels the one we will carry out, but does not really affect our conclusions given the inherent uncertainties in the computation of matrix elements. It has also been argued that $A(\Xi^-_-)$ is probably smaller than $A(\Lambda^0_-)$ due to much smaller strong rescattering phases [8].

$$s \approx -\sqrt{\frac{2}{3}} s_1 = 1.47 \pm 0.01 \; ,$$
$$p \approx -\sqrt{\frac{2}{3}} p_1 = \left(\frac{|\vec{q}|}{E_P + M_P} \right) (9.98 \pm 0.24) \; . \tag{16}$$

Substituting the experimental numbers for the amplitudes and strong rescattering phases, one gets

$$A(\Lambda^0_-) \approx 0.13 \sin(\phi_1^p - \phi_1^s) + 0.001 \sin(\phi_1^p - \phi_3^s) \\ -0.0024 \sin(\phi_3^p - \phi_1^s) \; . \tag{17}$$

A. Standard model calculation

In the case of the minimal standard model, the CP-violating phase resides in the CKM matrix. For low energy transitions, this phase shows up as the imaginary part of the Wilson coefficients in the effective weak Hamiltonian. In the notation of Buras and Harlander [11],

$$H_W^{\rm SM} = \frac{G_F}{\sqrt{2}} V_{ud}^* V_{us} \sum_i c_i(\mu) Q_i(\mu) + \text{H.c.} \tag{18}$$

$Q_i(\mu)$ are four-quark operators, and $c_i(\mu)$ are the Wilson coefficients that are usually written as

$$c_i(\mu) = z_i(\mu) + \tau y_i(\mu) \; ,$$
$$\tau = -\frac{V_{td}^* V_{ts}}{V_{ud}^* V_{us}} \; , \tag{19}$$

with the CP-violating phase being the phase of τ. Numerical values for these coefficients can be found, for example, in Buchalla *et al.* [11].

The calculation of the weak phases would proceed by evaluating the hadronic matrix elements of the four-quark operators in Eq. (18) to obtain real and imaginary parts for the amplitudes, schematically,

$$\langle p\pi | H_w^{\rm eff} | \Lambda^0 \rangle |_l^I = \text{Re} M_l^I + i \text{Im} M_l^I \; , \tag{20}$$

and to the extent that the CP-violating phases are small, they can be approximated by

$$\phi_l^I \approx \frac{\text{Im} M_l^I}{\text{Re} M_l^I} \; . \tag{21}$$

At present, however, we do not know how to compute the matrix elements, and so we cannot actually implement this calculation.

For a simple estimate, we can take the real part of the matrix elements from experiment (assuming that the measured amplitudes are real, that is, that CP violation is small) and compute the imaginary parts in vacuum saturation. This approach provides a conservative estimate for the weak phases because the model calculation of the real part of the amplitudes is smaller than the experimental value. Nevertheless, the numbers should be viewed with great caution.

The approximate weak phases estimated in vacuum saturation are [12]

$$\phi_s^1 \approx -3 y_6 \text{Im}\tau \; ,$$
$$\phi_p^1 \approx -0.3 y_6 \text{Im}\tau \; ,$$
$$\phi_s^3 \approx \left[3.6(y_1 + y_2) + 2.7(y_7 + 3y_8) \frac{B_0^2}{m_K^2} \right] \text{Im}\tau \; , \tag{22}$$
$$\phi_p^3 \approx \left[0.5(y_1 + y_2) - 0.4(y_7 + 3y_8) \frac{B_0^2}{m_K^2} \right] \text{Im}\tau \; .$$

These provide numerical estimates using the values for the Wilson coefficients[3] of Buchalla *et al.* [11], $y_6 \approx -0.08$, and the value of B_0 given in the Appendix. For the quantity $\text{Im}\tau$ (we use the Wolfenstein parametrization of the CKM matrix), we take

$$\text{Im}\tau = A^2 \lambda^4 \eta \leq 0.001 \; . \tag{23}$$

Putting all the numbers together and using the upper limit in Eq. (23) yields[4]

$$A(\Lambda^0_-) \approx 3 \times 10^{-5} \; . \tag{24}$$

Other models of CP violation contain additional short-distance operators with CP-violating phases [13–15] and predict different values for $A(\Lambda^0_-)$ [7,16]. A summary of results can be found, for example, in Ref. [17].

III. FOUR-QUARK OPERATORS

We now study, in a model-independent manner, the contributions to $A(\Lambda^0_-)$ that occur due to physics beyond the minimal standard model. In this section we look at the effect of all the four-quark operators and in the next section we discuss the two-quark operators. We assume that the physics that lies beyond the minimal standard model is characterized by a scale $\Lambda \gg M_W$ and, therefore, that its most important low energy effects can be parametrized by the lowest-dimension operators of the most general effective Lagrangian consistent with the symmetries of the standard model. Such a Lagrangian has been written down by Buchmüller and Wyler [18]. In the Appendix we list all the operators that occur at dimension six with $|\Delta S| = 1$.

The calculation then proceeds as in the previous section, but with the effective Hamiltonian

$$H_{\rm eff} = H_W^{\rm SM} + \frac{g^2}{\Lambda^2} \left(\sum_i \lambda_i \mathcal{O}_i^{\rm new} + \text{H.c.} \right) \; . \tag{25}$$

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[3]For $\mu = 1$ GeV, $\Lambda_{\rm QCD} = 200$ MeV.

[4]See Ref. [12] for additional discussions of this calculation.

To the usual, QCD-corrected, standard weak Hamiltonian of the previous section, we add all the four-fermion operators with $|\Delta S| = 1$ that come from the new physics sector. We will sidestep the issue of the possible origin of the effective CP-violating operators. We use the notation of Ref. [18] as detailed in the Appendix. These operators violate CP if the coupling λ_i has an imaginary part. The normalization has been chosen for convenience.

A. $K_L \to \pi\pi$ and ϵ'/ϵ

The standard notation for the $K \to \pi\pi$ amplitudes is

$$A(K^0 \to \pi^+\pi^-) = A_0 e^{i\delta_0} + \frac{A_2}{\sqrt{2}} e^{i\delta_2} ,$$
$$A(K^0 \to \pi^0\pi^0) = A_0 e^{i\delta_0} - \sqrt{2} A_2 e^{i\delta_2} , \tag{26}$$

where $\delta_{0,2}$ are the strong $\pi\pi$ scattering phases in the $I = 0, 2$ channel. The amplitudes A_0 and A_2 are real unless there is CP violation. Experimentally it is known that the $\Delta I = 3/2$ amplitude A_2 is much smaller than the $\Delta I = 1/2$ amplitude A_0:

$$\omega \equiv \frac{\mathrm{Re}A_2}{\mathrm{Re}A_0} \approx \frac{1}{22} . \tag{27}$$

The contribution of the dominant penguin operator ($\mathcal{O}_6$ in the notation of Ref. [11]) to ϵ'/ϵ is given by

$$\left(\frac{\epsilon'}{\epsilon}\right)_6 = -\frac{\omega}{\sqrt{2}|\epsilon|} \frac{\mathrm{Im}(A_0)_6}{|A_0|} = \frac{\omega}{2|\epsilon|} \frac{G_F}{|A_0|} y_6 \lambda \mathrm{Im}\tau \langle \pi^+\pi^- | \mathcal{O}_6 | K^0 \rangle . \tag{28}$$

The hadronic uncertainty enters the calculation through the matrix element of the four-quark operator. We will use the estimate of Ref. [11] for the matrix element of $\mathcal{O}_6$:

$$\langle \pi^+\pi^- | \mathcal{O}_6 | K^0 \rangle \bigg|_{I=0} = -4\sqrt{2} f_\pi \frac{m_K^2 - m_\pi^2}{\Lambda_\chi^2} \left(\frac{m_K^2}{m_s + m_d}\right)^2 \approx -0.26 \ \mathrm{GeV}^3 . \tag{29}$$

Using the values $A = 0.9$, $\lambda = 0.22$, and $\eta = 0.5$, one finds that

$$\left(\frac{\epsilon'}{\epsilon}\right)_6 \approx 1.5 \times 10^{-3} . \tag{30}$$

The usual standard model analysis of ϵ'/ϵ consists of computing this contribution of the "penguin" operator and of normalizing all other contributions to it in terms of a parameter Ω:

$$\frac{\epsilon'}{\epsilon} = \left(\frac{\epsilon'}{\epsilon}\right)_6 (1 - \Omega_{\mathrm{SM}} - \Omega_{\mathrm{new}}) . \tag{31}$$

Ω_{SM} is given, for example, in Ref. [11], and we have introduced an analogous term Ω_{new} for the contributions of the new four-quark operators. Given the experimental result in Eq. (2), we will place bounds on the new physics by requiring, conservatively, that $\Omega_{\mathrm{new}} \leq 1$. We find

$$\Omega_{\mathrm{new}} = 8 \left(\frac{M_W}{\Lambda}\right)^2 \sum_i \left(\frac{\mathrm{Im}\lambda_i}{A^2\lambda^5\eta}\right) \left[\frac{\langle \pi^+\pi^- | \mathcal{O}_i | K^0 \rangle_{I=0}}{y_6 \langle \pi^+\pi^- | \mathcal{O}_6 | K^0 \rangle_{I=0}} - \frac{\sqrt{2}}{\omega} \frac{\langle \pi^+\pi^- | \mathcal{O}_i | K^0 \rangle_{I=2}}{y_6 \langle \pi^+\pi^- | \mathcal{O}_6 | K^0 \rangle_{I=0}}\right] . \tag{32}$$

Because there is no way at present to compute the matrix elements of four-quark operators reliably, we will simply use vacuum saturation. The new contributions to ϵ' can thus be computed with the aid of the matrix elements listed in Table VIII, below. We use, as before, $A^2\lambda^5\eta \approx 2 \times 10^{-4}$, and we explicitly separate the contributions from the different isospin components of each operator for later convenience. We thus write

$$\Omega_{\mathrm{new}} = 4 \times 10^4 \left(\frac{M_W}{\Lambda}\right)^2 \sum_i \mathrm{Im}\lambda_i (\omega_{0i} + \omega_{2i}) , \tag{33}$$

where $\omega_{0,2i}$ refers to the $\Delta I = 1/2, 3/2$ component of $\mathcal{O}_i$. We present numerical results for $\omega_{0,2i}$ in Table I.

Requiring that $\Omega_{\mathrm{new}} < 1$, we can constrain the size of the CP-violating couplings $\mathrm{Im}\lambda_i/\Lambda^2$. By assuming that there is no accidental cancellation between the contributions of different operators to Ω_{new}, we may constrain each operator separately. The isospin decomposition is useful because it is possible to construct combinations of operators with definite isospin transformation properties. The constraints that apply to operators that are purely $\Delta I = 1/2$ are different from those that apply to operators that are purely $\Delta I = 3/2$.

TABLE I. Numerical coefficients for Eq. (33).

Operator	ω_0	ω_2
$\mathcal{O}_{qq}^{(1,1)}$	-0.06	0
$\mathcal{O}_{qq}^{(8,1)}$	-0.3	0
$\mathcal{O}_{qq}^{(1,3)}$	-0.3	0
$\mathcal{O}_{qq}^{(8,3)}$	0.32	0
$\mathcal{O}_{dd}^{(1)}$	0.08	2.5
$\mathcal{O}_{ud}^{(1)}$	-0.02	-2.5
$\mathcal{O}_{dd}^{(8)}$	0.1	3.3
$\mathcal{O}_{ud}^{(8)}$	0.2	-3.3
$\mathcal{O}_{qu}^{(1)}$	2.4	-36.8
$\mathcal{O}_{qu}^{(8)}$	0.1	3.3
$\mathcal{O}_{qd}^{(1)}$	1.5	0
$\mathcal{O}_{qsd}^{(1)}$	-3.9	36.8
$\mathcal{O}_{qsd}^{(1)}$	-0.1	-3.3
$\mathcal{O}_{qsq}^{(1)}$	1.8	-31.2
$\mathcal{O}_{qsq}^{(8)}$	-3.5	33
$\mathcal{O}_{qqs}^{(1)}$	-3.5	31
$\mathcal{O}_{qqs}^{(8)}$	2.1	-33
$\mathcal{O}_{qq}^{(1s)}$	1.8	0
$\mathcal{O}_{qq}^{(8s)}$	1.3	0

B. K^0-$\bar{K}^0$ mixing and ϵ

In general, ϵ' provides tighter constraints on new CP-violating interactions than does ϵ. Nevertheless, it is necessary to consider constraints from ϵ because the ones that arise from ϵ' do not apply to parity-conserving operators that do not contribute to the decay $K^0 \to \pi\pi$. In the operator basis that we are using, all the operators have parity-conserving and -violating components. However, it is possible to construct parity-conserving combinations of operators just as it is possible to construct combinations of operators with definite isospin.

All of the $|\Delta S| = 1$ four-quark operators that we consider contribute to ϵ when combined with a second $|\Delta S| = 1$ vertex from the usual weak Hamiltonian. In terms of the K^0-$\bar{K}^0$ mixing matrix, each operator gives a contribution to ϵ of the form

$$|\epsilon|_i \approx \frac{1}{\sqrt{2}} \frac{|\mathrm{Im} M_{12}|_i}{\Delta m_k} . \quad (34)$$

We estimate the long-distance contributions to $\mathrm{Im}M_{12}$ due to intermediate pion and η poles [21]. Using the matrix elements of Table IX, below, we find that there is a cancellation between the contributions of the pion and octet-η poles at leading order in SU(3) breaking. This situation is unfortunate because it makes the estimates very unreliable. For our purposes we will use the model of Ref. [19] to deal with this problem.

The contribution of each operator to ϵ is given by

$$|\epsilon|_{\mathcal{O}_i} = \sqrt{2} g^2 \frac{g_8}{M_W^2} \left(\frac{m_K}{\Delta m_K}\right) |\mathrm{Im}\lambda_i| \left(\frac{M_W}{\Lambda}\right)^2 \frac{m_K^2}{m_K^2 - m_\pi^2} |\xi_i| \approx 9.3 |\mathrm{Im}\lambda_i| \left(\frac{M_W}{\Lambda}\right)^2 |\xi_i| , \quad (35)$$

where g_8 is defined in Eq. (53) and ξ_i is given according to the model of Ref. [19] by

$$\xi_i = \frac{1}{\sqrt{2} f_\pi^2 m_K^2} \frac{f_K}{f_\pi} \left(\langle \pi^0 | \mathcal{O}_i | \bar{K}^0 \rangle + \frac{\langle \eta_8 | \mathcal{O}_i | \bar{K}^0 \rangle}{\sqrt{3}} \left\{ \left(\frac{m_K^2 - m_\pi^2}{m_K^2 - m_\eta^2} \right) [(1+\xi)\cos\theta + 2\sqrt{2}\rho\sin\theta] \left[\frac{f_{\eta_8}}{f_\pi} \cos\theta - \sqrt{2} \frac{f_{\eta_0}}{f_\pi} \sin\theta \right] \right.\right.$$
$$\left.\left. + \left(\frac{m_K^2 - m_\pi^2}{m_K^2 - m_\eta'^2} \right) \left[(1+\xi)\sin\theta - 2\sqrt{2}\rho\cos\theta \right] \left[\frac{f_{\eta_8}}{f_\pi} \sin\theta + \sqrt{2} \frac{f_{\eta_0}}{f_\pi} \cos\theta \right] \right\} \right) . \quad (36)$$

We choose the parameters that Ref. [19] considers more physical: $\theta = -20°$, $\xi = 0.17$, $f_{\eta_8} = 1.25 f_\pi$, and $f_{\eta 0} = 1.04 f_\pi$.

Once more we present separate results for the $\Delta I = 1/2, 3/2$ components of each operator in Table II. We emphasize again that we present our results in this form because it is possible to construct combinations of operators that have definite isospin transformation properties.

TABLE II. Factors ξ_i for Eq. (36).

Operator	$\xi_{i,1/2}(\rho = 0.8)$	$\xi_{i,1/2}(\rho = 1.2)$	$\xi_{i,1/2}(\pi \text{ only})$	$\xi_{i,3/2}$
$\mathcal{O}_{qq}^{(1,1)}$	−0.24	0.14	−0.04	0
$\mathcal{O}_{qq}^{(8,1)}$	−0.41	−0.12	−0.22	0
$\mathcal{O}_{qq}^{(1,3)}$	−0.33	−0.17	−0.21	0
$\mathcal{O}_{qq}^{(8,3)}$	−0.16	0.71	0.22	0
$\mathcal{O}_{dd}^{(1)}$	−0.18	0.04	−0.06	−0.11
$\mathcal{O}_{ud}^{(1)}$	−0.06	0.1	0.01	0.11
$\mathcal{O}_{dd}^{(8)}$	−0.23	0.06	−0.07	−0.15
$\mathcal{O}_{ud}^{(8)}$	−0.18	−0.18	−0.15	0.15
$\mathcal{O}_{qu}^{(1)}$	−1.8	−1.7	−1.5	1.5
$\mathcal{O}_{qu}^{(8)}$	−0.05	0.24	0.07	0.15
$\mathcal{O}_{qd}^{(1)}$	−4.2	−1.2	−2.2	0
$\mathcal{O}_{qsd}^{(1)}$	−2.4	0.57	−0.75	−1.5
$\mathcal{O}_{qsd}^{(8)}$	−0.23	0.06	−0.07	−0.15
$\mathcal{O}_{qsq}^{(1)}$	−1.7	−2.2	−1.6	1.2
$\mathcal{O}_{qsq}^{(8)}$	0.46	−2.1	−0.65	−1.3
$\mathcal{O}_{qqs}^{(1)}$	0.22	−2.6	−1.0	−1.2
$\mathcal{O}_{qqs}^{(8)}$	−1.6	−1.6	−1.3	1.3
$\mathcal{O}_{qq}^{(1s)}$	−1.5	−4.8	−2.6	0
$\mathcal{O}_{qq}^{(8s)}$	−1.1	−3.7	−2.0	0

For the $\Delta I = 1/2$ component, there is sensitivity to the parameters in the model of Ref. [19]. We illustrate this by presenting results for $\rho = 0.8$, $\rho = 1.2$, and for just the pion pole. For the $\Delta I = 3/2$ component there is only a pion pole.

C. $\Lambda \to p\pi^-$ and $A(\Lambda^0_-)$

The starting point of the calculation is Eq. (17). We study the effect of the new physics one operator at a time and always assume that the CP-violating amplitudes are small, so that the experimental value of the amplitudes is approximately equal to the CP-conserving amplitude. All the CP-violating phases are then small, and we can write

$$A(\Lambda^0_-) \approx 3\times 10^{-5} + \sum_i [0.13(\phi_1^p - \phi_1^s) + 0.001(\phi_1^p - \phi_3^s) - 0.0024(\phi_3^p - \phi_1^s)]_i \ , \quad (37)$$

where the sum runs over all the operators in Eq. (25).

We carry out the calculation in the same manner as the standard model analysis of the previous section [12]. That is, we compute the imaginary part of the amplitudes by taking matrix elements of each new four-quark operator in vacuum saturation. Further, we will not compute perturbative QCD corrections to the effective Hamiltonian of the new physics sector. We will also assume that the new physics does not significantly alter the CP-conserving amplitudes, but we will comment on this later on. As discussed in Ref. [12], this vacuum saturation calculation is not reliable at all; nevertheless, we will use it for lack of anything better.

Calculating the imaginary part of the amplitudes taking the real part from experiment as in the previous section, we find that each operator $\mathcal{O}_i$ induces the phases

$$\begin{aligned}
(\phi_1^p)_i &= -8\frac{G_F}{\sqrt{2}}\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im}\lambda_i \frac{\langle p\pi^-|\mathcal{O}_i|\Lambda\rangle_1^P}{9.98G_F m_\pi^2}\ , \\
(\phi_1^s)_i &= 8\frac{G_F}{\sqrt{2}}\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im}\lambda_i \frac{\langle p\pi^-|\mathcal{O}_i|\Lambda\rangle_1^S}{1.47G_F m_\pi^2}\ , \\
(\phi_3^p)_i &= 8\frac{G_F}{\sqrt{2}}\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im}\lambda_i \frac{\langle p\pi^-|\mathcal{O}_i|\Lambda\rangle_3^P}{0.21G_F m_\pi^2}\ , \\
(\phi_3^s)_i &= -8\frac{G_F}{\sqrt{2}}\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im}\lambda_i \frac{\langle p\pi^-|\mathcal{O}_i|\Lambda\rangle_3^S}{0.03G_F m_\pi^2}\ .
\end{aligned} \quad (38)$$

The matrix elements are estimated in vacuum saturation and listed in Table VII, below, in the Appendix. Numerically we find

$$A(\Lambda^0_-) \approx 3\times 10^{-5} + \left(\frac{M_W}{\Lambda}\right)^2 \sum_i \mathrm{Im}\lambda_i a_i\ , \quad (39)$$

where the coefficients a_i are listed in Table III. We present two different values: In the first column we include only the $\Delta I = 1/2$ component of each operator, whereas in the second column we include both isospin components. We can see from Table III that the CP-violating asymmetry $A(\Lambda^0_-)$ is dominated by the interference of the s and p waves in the $\Delta I = 1/2$ amplitude, as can be anticipated from Eq. (17).

TABLE III. Factors a_i for $A(\Lambda^0_-)_{\rm new}$ in Eq. (39).

Operator	$(a_i)_{1/2}$	a_i
$\mathcal{O}_{qq}^{(1,1)}$	0.03	0.03
$\mathcal{O}_{qq}^{(8,1)}$	0.15	0.14
$\mathcal{O}_{qq}^{(1,3)}$	0.14	0.14
$\mathcal{O}_{qq}^{(8,3)}$	−0.15	−0.15
$\mathcal{O}_{dd}^{(1)}$	−0.05	−0.06
$\mathcal{O}_{ud}^{(1)}$	0.01	0.02
$\mathcal{O}_{dd}^{(8)}$	−0.06	−0.08
$\mathcal{O}_{ud}^{(8)}$	−0.1	−0.1
$\mathcal{O}_{qu}^{(1)}$	−1.3	−1.1
$\mathcal{O}_{qu}^{(8)}$	−0.05	−0.08
$\mathcal{O}_{qd}^{(1)}$	1.5	1.5
$\mathcal{O}_{qsd}^{(1)}$	−0.6	−0.8
$\mathcal{O}_{qsd}^{(8)}$	0.05	0.08
$\mathcal{O}_{qsq}^{(1)}$	−1.4	−1.2
$\mathcal{O}_{qsq}^{(8)}$	0.6	0.7
$\mathcal{O}_{qqs}^{(1)}$	−0.9	−1.0
$\mathcal{O}_{qqs}^{(8)}$	−1.1	−1.0
$\mathcal{O}_{qq}^{(1s)}$	1.8	1.8
$\mathcal{O}_{qq}^{(8s)}$	1.4	1.3

We also see from Table III that a_i is of order one in some cases. Equation (39) then tells us that a measurement of $A(\Lambda^0_-)$ at the 10^{-4} level is sensitive, in principle, to new CP-violating interactions generated at a scale $\Lambda \leq 8$ TeV and is thus potentially interesting.

We can use the constraints from CP violation in kaon decays to place bounds on the magnitude of $A(\Lambda^0_-)$ that each of the four-quark operators can induce In general, the bounds coming from direct CP violation in ϵ' are stronger than those coming from ϵ. However, it is necessary to consider both because it is possible to construct parity-conserving combinations of operators that do not contribute to $K \to \pi\pi$ amplitudes and, thus, evade the bounds from ϵ'. Similarly, ϵ' places stronger constraints on $\Delta I = 3/2$ operators than on $\Delta I = 1/2$ operators due to the enhancement factor of $1/\omega$ in Eq. (32). To take into account these distinctions, we list in Table IV the bounds on each of the weak phases separately. The blank entries indicate that there is no bound because the particular operator does not contribute to that amplitude.

The bounds on the p-wave phases arise from the contributions of the operator to ϵ and are weaker than the bounds on the s-wave phases that arise from the contributions to ϵ'. For the operator basis that we have been using, the bounds on the different components are not independent. This, however, is not an important point because there is nothing special about this operator basis. We prefer to illustrate separately the bounds on each parity and isospin amplitude because it is possible to construct operators with definite parity and isospin.

TABLE IV. Bounds on the phases that enter $A(\Lambda^0_-)$.

Operator	$\phi_1^p \times 10^5$	$\phi_1^s \times 10^5$	$\phi_3^p \times 10^5$	$\phi_3^s \times 10^5$
$\mathcal{O}_{qq}^{(1,1)}$	2.9	-10		
$\mathcal{O}_{qq}^{(8,1)}$	9.0	-10		
$\mathcal{O}_{qq}^{(1,3)}$	10	-10		
$\mathcal{O}_{qq}^{(8,3)}$	-24	-10		
$\mathcal{O}_{dd}^{(1)}$	5.3	-10	400	-16
$\mathcal{O}_{ud}^{(1)}$	-3.7	-10	400	-16
$\mathcal{O}_{dd}^{(8)}$	5.3	-10	400	-16
$\mathcal{O}_{ud}^{(8)}$	14	-10	400	-16
$\mathcal{O}_{qu}^{(1)}$	14	-9.3	400	-15
$\mathcal{O}_{qu}^{(8)}$	-24	-10	400	-16
$\mathcal{O}_{qd}^{(1)}$	9	22		
$\mathcal{O}_{qsd}^{(1)}$	5.4	2.8	-400	-15
$\mathcal{O}_{qsd}^{(8)}$	5.3	-10	400	-16
$\mathcal{O}_{qsq}^{(1)}$	16	-14	400	-15
$\mathcal{O}_{qsq}^{(8)}$	24	-2.8	-400	15
$\mathcal{O}_{qqs}^{(1)}$	-74	4.2	400	-15
$\mathcal{O}_{qqs}^{(8)}$	14	-9.3	400	-15
$\mathcal{O}_{qq}^{(1s)}$	29	22		
$\mathcal{O}_{qq}^{(8s)}$	29	22		

IV. TWO-QUARK OPERATORS OF DIMENSION SIX

In addition to the four-quark operators of dimension six considered in the previous section, there are also two-quark operators of dimension six that can contribute to the processes under consideration [18]. These operators are the SU(3)×SU(2)×U(1)-invariant versions of "penguin" operators that naively appear to be dimension five [20]. There are two types of operators that contribute to CP violation in $|\Delta S| = 1$ processes. The first one in the notation of Ref. [18] is

$$\mathcal{O}_{dG} = (\bar{q}\sigma_{\mu\nu}\lambda^a d)\phi G^a_{\mu\nu} \; . \tag{40}$$

The operator of interest to us is obtained when the scalar doublet ϕ takes its vacuum expectation value. This leads to the effective Lagrangian (with the same overall normalization that we used before and $v \approx 246$ GeV)

$$\begin{aligned}\mathcal{L}_p = \frac{g^2}{\Lambda^2}\frac{v}{\sqrt{2}}\Bigg[&\lambda_{ds}\bar{d}\sigma_{\mu\nu}\lambda^a\left(\frac{1+\gamma_5}{2}\right)sG^a_{\mu\nu} \\ &+\lambda^*_{sd}\bar{d}\sigma_{\mu\nu}\lambda^a\left(\frac{1-\gamma_5}{2}\right)sG^a_{\mu\nu}\Bigg] + \text{H.c.} \\ \equiv \frac{g^2}{\Lambda^2}\frac{v}{\sqrt{2}}&\bar{d}\sigma_{\mu\nu}t^a(f_{\rm PC}+\gamma_5 f_{\rm PV})sG^a_{\mu\nu} + \text{H.c.}\end{aligned} \tag{41}$$

There are also analogous operators where the gluon field strength tensor is replaced by field strength tensors for electroweak gauge bosons. The matrix elements of these operators are suppressed by a power of $\alpha = 1/137$ with respect to the gluon operator and we will, therefore, neglect them.

A. Constraint on the parity-conserving coupling

The parity-conserving coupling $f_{\rm PC}$ is constrained by the contribution of Eq. (41) to the parameter ϵ. Unlike the four-quark operators of the previous section, we cannot use vacuum saturation to compute the matrix elements of this operator. However, this is the same operator that arises in the Weinberg model of CP violation, and the analysis has been carried out by Donoghue and Holstein [21] using MIT bag model matrix elements. We can simply take over their results to find

$$|\epsilon|_p \approx 1.5\times 10^5|\xi|\left(\frac{M_W}{\Lambda}\right)^2|\text{Im}f_{\rm PC}| \; . \tag{42}$$

This contribution to ϵ is due to long-distance effects as those discussed in the previous section. In complete analogy we have introduced the parameter ξ which takes the values $\xi = 0.12$ for $\rho = 0.8$ and $\xi = -0.48$ for $\rho = 1.2$. We find that the sensitivity of the result to the SU(3)-breaking parameters of the pole model is larger in this case than it was for the four-quark operators.

B. Constraint on the parity-violating coupling

The constraint on the parity-violating coupling $f_{\rm PV}$ comes from an analysis of ϵ'. Just as we did for $f_{\rm PC}$, we simply take over the results of Ref. [21] with a suitable identification of the coupling. We find

$$\left|\frac{\epsilon'}{\epsilon}\right|_p \approx 2.2\times 10^5\left(\frac{M_W}{\Lambda}\right)^2|\text{Im}f_{\rm PV}| \; . \tag{43}$$

C. Contribution to $A(\Lambda^0_-)$

Once again, we use the fact that up to coupling constants this operator is the same one appearing in the Weinberg model of CP violation. Its matrix elements using the MIT bag model can thus be taken from Ref. [7]. We find

$$\phi_1^s \approx 7\times 10^4 \left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im} f_{\mathrm{PV}} \,,$$
$$\phi_1^p \approx -8\times 10^4 \left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im} f_{\mathrm{PC}} \,. \quad (44)$$

From these it follows that

$$A(\Lambda^0_-)_p \approx -10^4 \left(\frac{M_W}{\Lambda}\right)^2 (\mathrm{Im} f_{\mathrm{PC}} + 0.9 \mathrm{Im} f_{\mathrm{PV}}) \,. \quad (45)$$

In the Weinberg model this operator appears with $f_{\mathrm{PV}} = -f_{\mathrm{PC}}$ and there is a large cancellation between the two phases, leading to a smaller value for $A(\Lambda^0_-)$ than would have been obtained from each phase individually [7]. In our general operator analysis, the bounds from Eqs. (42) and (43) can be combined to obtain (with $\xi = -0.5$)

$$\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im} f_{\mathrm{PC}} < 2.7\times 10^{-8} \,,$$
$$\left(\frac{M_W}{\Lambda}\right)^2 \mathrm{Im} f_{\mathrm{PV}} < 6.8\times 10^{-9} \,, \quad (46)$$

or, in terms of the hyperon decay observable,

$$A(\Lambda^0_-) \leq \begin{cases} 3\times 10^{-4} & \text{parity-conserving operator}\,, \\ 6\times 10^{-5} & \text{parity-violating operator}\,. \end{cases} \quad (47)$$

Before ending this section we should comment on one class of two-quark operators that we have not discussed. In the notation of Ref. [18] it is

$$\mathcal{O}_{qG} = i(\bar{q}\lambda^a\gamma_\mu D_\nu d)\phi G^a_{\mu\nu} \quad (48)$$

and related operators with field strength tensors for electroweak gauge bosons instead of the gluon. These latter ones will have matrix elements suppressed by α compared to Eq. (48). We have not found a simple way to estimate the matrix elements of these operators, and for this reason we do not discuss them in detail. We do not expect the behavior of this type of operator to be significantly different from the others that we have discussed.

V. SUMMARY AND CONCLUSIONS

The minimal standard model of electroweak interactions is in extraordinary agreement with all experiments conducted so far, and there is no evidence for any new particles below 100 GeV or so. In view of this, it is reasonable to assume that any new physics beyond the minimal standard model is associated with a scale $\Lambda \geq M_W$, and it is, therefore, possible to represent the low energy effects of any such new physics with an effective Lagrangian that respects the symmetries of the standard model.

In this paper we have studied all the $|\Delta S| = 1$, CP-violating, operators that occur at dimension six. We have investigated the constraints that exist on the couplings of these operators from the measurements of ϵ and ϵ', and estimated what their largest contribution to CP violation in $A(\Lambda^0_-)$ could be.

The operators that we have discussed also contribute to CP-conserving and flavor changing amplitudes. We might thus worry that the constraints on the real part of the couplings are such that it is not natural for the imaginary (CP-violating) part of the couplings to attain the upper bounds allowed by the values of ϵ and ϵ'. We briefly address this issue in this section.

Consider the contributions to K^0-$\bar{K}^0$ mixing. If we fix $\mathrm{Im}\lambda_i$ to its maximum allowed value, we find that the constrain $2\mathrm{Re}M_{12,i} \leq \Delta m_K$ is also satisfied if

$$\mathrm{Re}\lambda_i \leq \frac{\mathrm{Im}\lambda_i}{2\sqrt{2}\epsilon} \,. \quad (49)$$

Therefore, the CP-conserving constraint is also satisfied if both real and imaginary parts of the couplings are of the same size or if the imaginary part is smaller than the real part by a factor of ϵ. The strongest constraints on flavor-changing operators in the CP-conserving case are known to come from K^0-$\bar{K}^0$ mixing. If we set the couplings to be of order one, we obtain a lower bound on the scale of new physics Λ requiring that $2\mathrm{Re}M_{12,i} \leq \Delta m_K$. It is easy to check that with couplings and scales satisfying this bound the new operators do not make any significant contributions to the real part of the amplitudes in $K \to \pi\pi$ or $\Lambda \to p\pi$. Therefore, we conclude that fixing the imaginary part of the couplings to their maximum allowed value is not in conflict with CP-conserving constraints.

In the minimal standard model, we have estimated previously [12] that $A(\Lambda^0_-)$ is of the order of a few times 10^{-5}. For the new physics considered in this paper, we find that most of the operators would naturally induce contributions to $A(\Lambda^0_-)$ at the 10^{-5} level, making them indistinguishable from the minimal standard model (as long as precise calculations of the matrix elements are not available) and inaccessible to the search to be conducted by E871. However, we have also found that for certain operators $\mathcal{O}^{(1)}_{qqs}$ and $\mathcal{O}_{dG}$, $A(\Lambda^0_-)$ could be as large as a few times 10^{-4}.

Given our crude estimate of the hadronic matrix elements involved, all our numerical results should be viewed with caution. Nevertheless, our results suggest that the search for CP violation in $A(\Lambda^0_-)$ at the 10^{-4} level of sensitivity that is expected for E871 is potentially very interesting. Our results also suggest that this measurement is complementary to the measurement of ϵ'/ϵ, in that it probes potential sources of CP violation at a level that has not been probed by the kaon experiments. This is particularly true for parity-conserving

TABLE V. Division six $|\Delta S| = 1$ four-quark operators. We list in the second column the gauge-invariant version of the operator in the notation of Ref. [18] and in the third column the $|\Delta S| = 1$ components (in some cases there is more than one).

Operator	Ref. [18]	$\|\Delta S\| = 1$
$\mathcal{O}_{qq}^{(1,1)}$	$\mathcal{O}^{(1,1)}$	$\frac{1}{2}d_L\gamma_\mu s_L(\bar{u}_L\gamma_\mu u_L + \bar{d}_L\gamma_\mu d_L)$
$\mathcal{O}_{qq}^{(8,1)}$	$\mathcal{O}^{(8,1)}$	$\frac{1}{2}\bar{d}_L\lambda^a\gamma_\mu s_L(\bar{u}_L\lambda^a\gamma_\mu u_L + \bar{d}_L\lambda^a\gamma_\mu d_L)$
$\mathcal{O}_{qq}^{(1,3)}$	$\mathcal{O}^{(1,3)}$	$\frac{1}{2}(2\bar{u}_L\gamma_\mu s_L\bar{d}_L\gamma_\mu u_L - \bar{u}_L\gamma_\mu u_L\bar{d}_L\gamma_\mu s_L + \bar{d}_L\gamma_\mu d_L\bar{d}_L\gamma_\mu s_L)$
$\mathcal{O}_{qq}^{(8,3)}$	$\mathcal{O}^{(8,3)}$	$\frac{1}{2}(2\bar{u}_L\lambda^a\gamma_\mu s_L\bar{d}_L\lambda^a\gamma_\mu u_L - \bar{u}_L\lambda^a\gamma_\mu u_L\bar{d}_L\lambda^a\gamma_\mu s_L + \bar{d}_L\lambda^a\gamma_\mu d_L\bar{d}_L\lambda^a\gamma_\mu s_L)$
$\mathcal{O}_{dd}^{(1)}$	$\mathcal{O}_{dd}^{(1)}$	$\frac{1}{2}\bar{d}_R\gamma_\mu s_R\bar{d}_R\gamma_\mu d_R$
$\mathcal{O}_{ud}^{(1)}$	$\mathcal{O}_{ud}^{(1)}$	$\frac{1}{2}\bar{u}_R\gamma_\mu u_R\bar{d}_R\gamma_\mu s_R$
$\mathcal{O}_{dd}^{(8)}$	$\mathcal{O}_{dd}^{(8)}$	$\frac{1}{2}\bar{d}_R\lambda^a\gamma_\mu s_R\bar{d}_R\lambda^a\gamma_\mu d_R$
$\mathcal{O}_{ud}^{(8)}$	$\mathcal{O}_{ud}^{(8)}$	$\frac{1}{2}\bar{u}_R\lambda^a\gamma_u u_R\bar{d}_R\lambda^a\gamma_\mu s_R$
$\mathcal{O}_{qu}^{(1)}$	$\mathcal{O}_{qu}^{(1)}$	$\bar{d}_L u_R\bar{u}_R s_L$
$\mathcal{O}_{qu}^{(8)}$	$\mathcal{O}_{qu}^{(8)}$	$\bar{d}_L\lambda^a u_R\bar{u}_R\lambda^a s_L$
$\mathcal{O}_{qd}^{(1)}$	$\mathcal{O}_{qd}^{(1)}$	$\bar{u}_L s_R\bar{d}_R u_L + \bar{d}_L s_R\bar{d}_R d_L$
$\mathcal{O}_{qsd}^{(1)}$	$\mathcal{O}_{qd}^{(1)}$	$\bar{d}_L d_R\bar{d}_R s_L$
$\mathcal{O}_{qsd}^{(8)}$	$\mathcal{O}_{qd}^{(8)}$	$\bar{d}_L\lambda^a d_R\bar{d}_R\lambda^a s_L$
$\mathcal{O}_{qsq}^{(1)}$	$\mathcal{O}_{qq}^{(1)}$	$-\bar{u}_R s_L\bar{d}_R u_L$
$\mathcal{O}_{qsq}^{(8)}$	$\mathcal{O}_{qq}^{(8)}$	$-\bar{u}_R\lambda^a s_L\bar{d}_R\lambda^a u_L$
$\mathcal{O}_{qqs}^{(1)}$	$\mathcal{O}_{qq}^{(1)}$	$\bar{d}_R s_L\bar{u}_R u_L$
$\mathcal{O}_{qqs}^{(8)}$	$\mathcal{O}_{qq}^{(8)}$	$\bar{d}_R\lambda^a s_L\bar{u}_R\lambda^a u_L$
$\mathcal{O}_{qq}^{(1s)}$	$\mathcal{O}_{qq}^{(1)}$	$\bar{u}_L u_R\bar{d}_L s_R - \bar{d}_L u_R\bar{u}_L s_R$
$\mathcal{O}_{qq}^{(8s)}$	$\mathcal{O}_{qq}^{(8)}$	$\bar{u}_L\lambda^a u_R\bar{d}_L\lambda^a s_R - \bar{d}_L\lambda^a u_R\bar{u}_L\lambda^a s_R$

interactions that do not contribute to ϵ' and are only constrained by ϵ.

We conclude that it is possible for E871 to observe a CP-violating signal at the 10^{-4} level. Our study indicates that if such a signal is observed, it would probably be evidence for physics beyond the minimal standard model. However, a reliable determination of hadronic matrix elements is necessary to reach any definite conclusion.

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APPENDIX: OPERATOR ANALYSIS

1. Dimension six $|\Delta S| = 1$ four-quark operators

In Table V we list all the four-quark operators of dimension six that change strangeness by one unit. We use the explicitly SU(3) $\times$ SU(2) $\times$ U(1) gauge-invariant notation of Ref. [18]. For each class of gauge-invariant operator, we give the components needed for this paper.

2. Isospin decomposition

For convenience we provide the isospin decomposition of the four-quark operators in Table VI.

3. Matrix elements for $\Lambda \to p\pi^-$ in vacuum saturation

We use the normalization in which $f_\pi = 93$ MeV and neglect $m_{u,d}/m_s$. In terms of

$$B_0 \equiv \frac{m_\pi^2}{m_u + m_d} = \frac{m_K^2}{m_s + m_u} \approx 11 m_\pi \ , \qquad \text{(A1)}$$

we find:

$$\langle p\pi^-|\bar{d}\gamma_\mu\gamma_5 u\bar{u}\gamma^\mu s|\Lambda\rangle \equiv M_V$$

$$= i\sqrt{2}f_\pi(M_\Lambda - M_P)\sqrt{3/2}\,\bar{\Psi}_p\Psi_\Lambda,$$

TABLE VI. Isospin decomposition of four-quark operators.

Operator	$\Delta I = 1/2$	$\Delta I = 3/2$
$3\bar{u}s\bar{d}u$	$2\bar{u}s\bar{d}u - \bar{d}s\bar{u}u + \bar{d}s\bar{d}d$	$\bar{u}s\bar{d}u + \bar{d}s\bar{u}u - \bar{d}s\bar{d}d$
$3\bar{d}s\bar{u}u$	$2\bar{d}s\bar{u}u - \bar{u}s\bar{d}u + \bar{d}s\bar{d}d$	$\bar{u}s\bar{d}u + \bar{d}s\bar{u}u - \bar{d}s\bar{d}d$
$3\bar{d}s\bar{d}d$	$\bar{u}s\bar{d}u + \bar{d}s\bar{u}u + 2\bar{d}s\bar{d}d$	$-\bar{u}s\bar{d}u - \bar{d}s\bar{u}u + \bar{d}s\bar{d}d$

TABLE VII. Matrix elements in $\Lambda \to p\pi^-$.

Operator	$\Delta I = 1/2$	$\Delta I = 3/2$
$\mathcal{O}_{qq}^{(1,1)}$	$\frac{1}{8N}(M_A - M_V)$	0
$\mathcal{O}_{qq}^{(8,1)}$	$\frac{1}{4}\left(1 - \frac{1}{N^2}\right)(M_A - M_V)$	0
$\mathcal{O}_{qq}^{(1,3)}$	$\frac{1}{4}\left(1 - \frac{1}{2N}\right)(M_A - M_V)$	0
$\mathcal{O}_{qq}^{(8,3)}$	$\frac{1}{4}\left(\frac{1}{N^2} - 1\right)(M_A - M_V)$	0
$\mathcal{O}_{dd}^{(1)}$	$\frac{1}{24}\left(1 + \frac{1}{N}\right)(M_A + M_V)$	$-\frac{1}{24}\left(1 + \frac{1}{N}\right)(M_A + M_V)$
$\mathcal{O}_{ud}^{(1)}$	$\frac{1}{24}\left(\frac{2}{N} - 1\right)(M_A + M_V)$	$\frac{1}{24}\left(1 + \frac{1}{N}\right)(M_A + M_V)$
$\mathcal{O}_{dd}^{(8)}$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(M_A + M_V)$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(M_A + M_V)$
$\mathcal{O}_{ud}^{(8)}$	$\frac{1}{6}\left(1 - \frac{1}{N^2}\right)(M_A + M_V)$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(M_A + M_V)$
$\mathcal{O}_{qu}^{(1)}$	$\frac{1}{12}\left(2\frac{B_0^2}{m_K^2}(M_V + M_A) + \frac{1}{2N}(M_V - M_A)\right)$	$\frac{1}{12}\left(\frac{B_0^2}{m_K^2}(M_V + M_A) - \frac{1}{2N}(M_V - M_A)\right)$
$\mathcal{O}_{qu}^{(8)}$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(M_V - M_A)$	$\frac{1}{12}\left(\frac{1}{N^2} - 1\right)(M_V - M_A)$
$\mathcal{O}_{qd}^{(1)}$	$-\frac{B_0^2}{4m_K^2}(M_V - M_A)$	0
$\mathcal{O}_{qsd}^{(1)}$	$\frac{1}{12}\left[\frac{B_0^2}{m_K^2}(M_V + M_A) - \frac{1}{2N}(M_V - M_A)\right]$	$\frac{1}{12}\left[-\frac{B_0^2}{m_K^2}(M_V + M_A) + \frac{1}{2N}(M_V - M_A)\right]$
$\mathcal{O}_{qsd}^{(8)}$	$\frac{1}{12}\left(\frac{1}{N^2} - 1\right)(M_V - M_A)$	$-\frac{1}{12}\left(\frac{1}{N^2} - 1\right)(M_V - M_A)$
$\mathcal{O}_{qsq}^{(1)}$	$\frac{1}{12}\left(2 + \frac{1}{2N}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$	$\frac{1}{12}\left(1 - \frac{1}{2N}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$
$\mathcal{O}_{qsq}^{(8)}$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$
$\mathcal{O}_{qqs}^{(1)}$	$\frac{1}{12}(1 + \frac{1}{N})\frac{B_0^2}{m_K^2}(M_V + M_A)$	$-\frac{1}{12}(1 - \frac{1}{2N})\frac{B_0^2}{m_K^2}(M_V + M_A)$
$\mathcal{O}_{qqs}^{(8)}$	$\frac{1}{6}\left(1 - \frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}(M_V + M_A)$
$\mathcal{O}_{qq}^{(1s)}$	$-\frac{1}{4}\left(1 + \frac{1}{2N}\right)\frac{B_0^2}{m_K^2}(M_V - M_A)$	0
$\mathcal{O}_{qq}^{(8s)}$	$-\frac{1}{4}\left(1 - \frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}(M_V - M_A)$	0

TABLE VIII. Matrix elements in $\bar{K}^0 \to \pi^+\pi^-$.

Operator	A_0	$A_2/\sqrt{2}$
$\mathcal{O}_{qq}^{(1,1)}$	$\frac{1}{8N}(V_2 - V_1)$	0
$\mathcal{O}_{qq}^{(8,1)}$	$\frac{1}{4}\left(1 - \frac{1}{N^2}\right)(V_2 - V_1)$	0
$\mathcal{O}_{qq}^{(1,3)}$	$\frac{1}{8}\left(2 - \frac{1}{N}\right)(V_2 - V_1)$	0
$\mathcal{O}_{qq}^{(8,3)}$	$-\frac{1}{4}\left(1 - \frac{1}{N^2}\right)(V_2 - V_1)$	0
$\mathcal{O}_{dd}^{(1)}$	$\frac{1}{24}\left(1 + \frac{1}{N}\right)(V_1 - V_2)$	$-\frac{1}{24}\left(1 + \frac{1}{N}\right)(V_1 + 2V_2)$
$\mathcal{O}_{ud}^{(1)}$	$\frac{1}{24}\left(\frac{2}{N} - 1\right)(V_1 - V_2)$	$\frac{1}{24}\left(1 + \frac{1}{N}\right)(V_1 + 2V_2)$
$\mathcal{O}_{dd}^{(8)}$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 - V_2)$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 + 2V_2)$
$\mathcal{O}_{ud}^{(8)}$	$\frac{1}{6}\left(1 - \frac{1}{N^2}\right)(V_1 - V_2)$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 + 2V_2)$
$\mathcal{O}_{qu}^{(1)}$	$\frac{1}{12}\left(2S_1 + \frac{1}{2N}(V_1 + V_2)\right)$	$\frac{1}{12}\left(S_1 - \frac{1}{2N}(V_1 - 2V_2)\right)$
$\mathcal{O}_{qu}^{(8)}$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 + V_2)$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 - 2V_2)$
$\mathcal{O}_{qd}^{(1)}$	$\frac{1}{4}(S_2 - S_1)$	0
$\mathcal{O}_{qsd}^{(1)}$	$\frac{1}{12}\left((S_1 - 3S_2) - \frac{1}{2N}(V_1 + V_2)\right)$	$\frac{1}{12}\left(-S_1 + \frac{1}{2N}(V_1 - 2V_2)\right)$
$\mathcal{O}_{qsd}^{(8)}$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 + V_2)$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(V_1 - 2V_2)$
$\mathcal{O}_{qsq}^{(1)}$	$\frac{1}{12}\left(2S_1 + \frac{1}{2N}(S_1 - 3S_2)\right)$	$\frac{1}{12}\left(1 - \frac{1}{2N}\right)S_1$
$\mathcal{O}_{qsq}^{(8)}$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)(S_1 - 3S_2)$	$-\frac{1}{12}\left(1 - \frac{1}{N^2}\right)S_1$
$\mathcal{O}_{qqs}^{(1)}$	$\frac{1}{12}\left(\left(1 + \frac{1}{N}\right)S_1 - 3S_2\right)$	$-\frac{1}{12}\left(1 - \frac{1}{2N}\right)S_1$
$\mathcal{O}_{qqs}^{(8)}$	$\frac{1}{6}\left(1 - \frac{1}{N^2}\right)S_1$	$\frac{1}{12}\left(1 - \frac{1}{N^2}\right)S_1$
$\mathcal{O}_{qq}^{(1s)}$	$-\frac{1}{4}\left(1 + \frac{1}{2N}\right)(S_1 - S_2)$	0
$\mathcal{O}_{qq}^{(8s)}$	$-\frac{1}{4}\left(1 - \frac{1}{N^2}\right)(S_1 - S_2)$	0

$$\langle p\pi^- | \bar{d}\gamma_\mu\gamma_5 u \bar{u}\gamma^\mu\gamma_5 s|\Lambda\rangle \equiv M_A$$

$$= -i\frac{2f_\pi f_K m_\pi^2}{m_K^2 - m_\pi^2}(-13.3)\bar{\Psi}_p\gamma_5\Psi_\Lambda, \tag{A2}$$

$$\langle p\pi^- | \bar{d}\gamma_5 u\bar{u}s|\Lambda\rangle = \frac{B_0^2}{m_K^2} M_V \ ,$$

$$\langle p\pi^- | \bar{d}\gamma_5 u\bar{u}\gamma_5 s|\Lambda\rangle = -\frac{B_0^2}{m_K^2} M_A \ .$$

We list the matrix element for each operator using vacuum saturation in Table VII.

4. Matrix elements for $\bar{K}^0 \to \pi^+\pi^-$ in vacuum saturation

In vacuum saturation we find

$$\begin{aligned}
V_1 &\equiv \langle\pi^-|\bar{d}\gamma_\mu\gamma_5 u|0\rangle\langle\pi^+|\bar{u}\gamma_\mu s|\bar{K}^0\rangle \\
&= -i\sqrt{2}f_\pi(m_K^2 - m_\pi^2) \ , \\
V_2 &\equiv \langle\pi^+\pi^-|\bar{u}\gamma_\mu u|0\rangle\langle 0|\bar{d}\gamma_\mu\gamma_5 s|\bar{K}^0\rangle \\
&= -\langle\pi^+\pi^-|\bar{d}\gamma_\mu d|0\rangle\langle 0|\bar{d}\gamma_\mu\gamma_5 s|\bar{K}^0\rangle \\
&= 0 \ , \\
S_1 &\equiv \langle\pi^-|\bar{d}\gamma_5 u|0\rangle\langle\pi^+|\bar{u}s|\bar{K}^0\rangle \\
&= -i\sqrt{2}f_\pi B_0^2\left(1 + 2\frac{m_\pi^2}{\Lambda_\chi^2}\right) \ , \\
S_2 &\equiv \langle\pi^+\pi^-|\bar{u}u|0\rangle\langle 0|\bar{d}\gamma_5 s|\bar{K}^0\rangle \\
&= \langle\pi^+\pi^-|\bar{d}d|0\rangle\langle 0|\bar{d}\gamma_5 s|\bar{K}^0\rangle \\
&= -i\sqrt{2}f_\pi B_0^2\left(1 + 2\frac{m_K^2}{\Lambda_\chi^2}\right) \ .
\end{aligned} \tag{A3}$$

We have to introduce momentum-dependent terms in the last two expressions because the leading terms cancel in

TABLE IX. Matrix elements in $\bar{K}^0 \to \pi^0, \eta_8, \eta_0$. An overall $\sqrt{2}f_\pi^2 m_K^2$ has been factored from all the entries in the table.

Operator	$\langle\pi^0\vert\mathcal{O}_1\vert\bar{K}^0\rangle$	$\sqrt{3}\langle\eta_8\vert\mathcal{O}_1\vert\bar{K}^0\rangle$	$\langle\pi^0\vert\mathcal{O}_3\vert\bar{K}^0\rangle$
$\mathcal{O}_{qq}^{(1,1)}$	$-\frac{1}{8N}$	$\frac{1}{8}\left(2+\frac{1}{N}\right)$	0
$\mathcal{O}_{qq}^{(8,1)}$	$-\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	$\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	0
$\mathcal{O}_{qq}^{(1,3)}$	$\frac{1}{8}\left(\frac{1}{N}-2\right)$	$\frac{3}{8N}$	0
$\mathcal{O}_{qq}^{(8,3)}$	$\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	$\frac{3}{4}\left(1-\frac{1}{N^2}\right)$	0
$\mathcal{O}_{dd}^{(1)}$	$-\frac{1}{24}\left(1+\frac{1}{N}\right)$	$\frac{1}{8}\left(1+\frac{1}{N}\right)$	$-\frac{1}{12}\left(1+\frac{1}{N}\right)$
$\mathcal{O}_{ud}^{(1)}$	$\frac{1}{24}\left(1-\frac{2}{N}\right)$	$\frac{1}{8}$	$\frac{1}{12}\left(1+\frac{1}{N}\right)$
$\mathcal{O}_{dd}^{(8)}$	$-\frac{1}{12}\left(1-\frac{1}{N^2}\right)$	$\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	$-\frac{1}{6}\left(1-\frac{1}{N^2}\right)$
$\mathcal{O}_{ud}^{(8)}$	$-\frac{1}{6}\left(1-\frac{1}{N^2}\right)$	0	$\frac{1}{6}\left(1-\frac{1}{N^2}\right)$
$\mathcal{O}_{qu}^{(1)}$	$-\frac{1}{12}\left(2\frac{B_0^2}{m_K^2}-\frac{1}{2N}\right)$	$\frac{1}{8N}$	$\frac{1}{12}\left(2\frac{B_0^2}{m_K^2}+\frac{1}{N}\right)$
$\mathcal{O}_{qu}^{(8)}$	$\frac{1}{12}\left(1-\frac{1}{N^2}\right)$	$\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	$\frac{1}{6}\left(1-\frac{1}{N^2}\right)$
$\mathcal{O}_{qd}^{(1)}$	$-\frac{1}{4}\frac{B_0^2}{m_K^2}$	$\frac{1}{4}\left(\frac{1}{N}+\frac{B_0^2}{m_K^2}\right)$	0
$\mathcal{O}_{qsd}^{(1)}$	$\frac{1}{12}\left(-\frac{B_0^2}{m_K^2}-\frac{1}{2N}\right)$	$-\frac{1}{4}\left(-\frac{B_0^2}{m_K^2}-\frac{1}{2N}\right)$	$\frac{1}{6}\left(-\frac{B_0^2}{m_K^2}-\frac{1}{2N}\right)$
$\mathcal{O}_{qsd}^{(8)}$	$-\frac{1}{12}\left(1-\frac{1}{N^2}\right)$	$\frac{1}{4}\left(1-\frac{1}{N^2}\right)$	$-\frac{1}{6}\left(1-\frac{1}{N^2}\right)$
$\mathcal{O}_{qsq}^{(1)}$	$-\frac{1}{12}\left(2+\frac{1}{2N}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{8N}\frac{B_0^2}{m_K^2}$	$\frac{1}{6}\left(1-\frac{1}{2N}\right)\frac{B_0^2}{m_K^2}$
$\mathcal{O}_{qsq}^{(8)}$	$-\frac{1}{12}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{4}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{6}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$
$\mathcal{O}_{qqs}^{(1)}$	$-\frac{1}{12}\left(1+\frac{1}{N}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{4}\frac{B_0^2}{m_K^2}$	$-\frac{1}{6}\left(1-\frac{1}{2N}\right)\frac{B_0^2}{m_K^2}$
$\mathcal{O}_{qqs}^{(8)}$	$-\frac{1}{6}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$	0	$\frac{1}{6}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$
$\mathcal{O}_{qq}^{(1s)}$	$-\frac{1}{4}\left(1+\frac{1}{2N}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{4}\left(1+\frac{1}{2N}\right)\frac{B_0^2}{m_K^2}$	0
$\mathcal{O}_{qq}^{(8s)}$	$-\frac{1}{4}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$	$-\frac{1}{4}\left(1-\frac{1}{N^2}\right)\frac{B_0^2}{m_K^2}$	0

the difference $S_1 - S_2$. The scale of chiral symmetry breaking, $\Lambda_\chi \approx 1$ GeV, can be related to the ratio f_K/f_π [22]. We list the matrix elements for each operator in Table VIII.

5. Matrix elements for $\bar{K}^0 \to \pi^0, \eta_8, \eta_0$ in vacuum saturation

The matrix elements for the $\bar{K}^0 \to \pi^0, \eta_8, \eta_0$ transition in the standard model are, to lowest order in chiral perturbation theory,

$$
\begin{aligned}
-i\langle \pi^0 | H_W | \bar{K}^0(q)\rangle &= -i2\frac{g_8}{f_\pi^2}\frac{q^2}{\sqrt{2}}\,, \\
-i\langle \eta_8 | H_W | \bar{K}^0(q)\rangle &= -i2\frac{g_8}{f_\pi^2}\frac{q^2}{\sqrt{6}}\,, \\
-i\langle \eta_0 | H_W | \bar{K}^0(q)\rangle &= i2\frac{g_8}{f_\pi^2}\frac{2q^2}{\sqrt{3}}\,,
\end{aligned}
\tag{A4}
$$

where $g_8 \approx 7.8 \times 10^{-8} f_\pi^2 \approx 10^{-13} M_W^2$.

For the matrix elements of the four-quark operators, we use vacuum saturation and U(3) symmetry to include the η singlet. The results are listed in Table IX where we have factored out a common $\sqrt{2} f_\pi^2 m_K^2$. For all the operators in Table IX, we have $\langle \eta_0 | \mathcal{O} | \bar{K}^0\rangle = \sqrt{2}\langle \eta_8 | \mathcal{O} | \bar{K}^0\rangle$.

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