Analysis of *CP* violation through phase angles in weak charged currents in a relativistic quark model*

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We test models of *CP* violation through phase angles in the weak currents of the type of Glashow for $K \to 2\pi$ decays. The analysis, based on a relativistic quark model for mesons, shows that the phases are bound to be very small ($\sim 10^{-4}$).

I. INTRODUCTION

A year after the discovery[1] of the decay $K_L \to 2\pi$ it was proposed that the *CP* violation could be due to a phase difference between the vector and the axial-vector Cabibbo currents.[2] Ten years later, the validity of these proposals is still uncertain, owing to the difficulty of computing their predictions. In fact, only estimates based on symmetry considerations have been done.[3] Here we present a detailed numerical analysis based on explicit Bethe-Salpeter (BS) amplitudes for the mesons,[4] which have been shown to reproduce accurately their decay rates,[5] as well as to imply the symmetry results of *CP* violation in the "milliweak" theory with charged currents.[6] In Sec. II we present the *CP*-violation model and the main phenomenology. In Sec. III we compute the $K \to 2\pi$ amplitudes and in Sec. IV we discuss the determination of the *CP*-violating phases and the results of the model.

II. *CP*-VIOLATION MODELS AND PHENOMENOLOGY

There have been proposed several models[2] in which the *CP* violation is attributed to phase angles between the weak vector and axial-vector charged currents. In the Glashow[2] model the phases are added in such a way as to have a weak charged current of $V-A$ type

$$J = \cos\theta_C (V + e^{i\phi}A)_{1+i2} + \sin\theta_C (V + e^{i\omega}A)_{4+i5} \tag{2.1}$$

instead of the standard Cabibbo current ($\phi = \omega = 0$).

In a model proposed by two of the authors,[2] the universal weak current was obtained through a generalization of the Cabibbo rotation of the isospin axial-vector current so as to include also a rotation around the other neutral F_6 direction in the unitary space. For the vector part the usual Cabibbo rotation is kept, but for the axial-vector part of the current, a rotation $\exp^{(2i\theta_A F_7 + 2i\phi_A F_6)}$ is performed. Then the *CP*-violation effects one gets for the nonleptonic weak Hamiltonian are due only to a "small" presence of the F_6 rotation in the axial-vector current. In this model, the weak charged current of $V-A$ type is

$$J = \cos\theta_C (V + A)_{1+i2} + \sin\theta_C (V + e^{i\omega}A)_{4+i5}, \tag{2.2}$$

where

$$e^{i\omega} = \frac{\theta_A + i\phi_A}{\theta_C}. \tag{2.3}$$

The second model, which leads to the current (2.2), can of course be formally considered as a particular case of (2.1), where no *CP* violation is allowed in the $\Delta S = 0$ current as suggested by β decay. However, in this model the phase angle ω has a physical meaning in terms of the Cabibbo-type rotations.

As usual the weak nonleptonic Hamiltonian is of the current × current form

$$\mathcal{H}_{NL} = \frac{G}{2\sqrt{2}} \{J, J^\dagger\}. \tag{2.4}$$

In order to connect the experimental *CP*-violating parameters with the phases ϕ and ω let us state a few remarks about the usual *CP* phenomenological analysis, and fix our phases and conventions.

The neutral kaons of strangeness $S = \pm 1$, K^0 and $\bar{K}^0$, are defined as

$$|\bar{K}^0\rangle = CP|K^0\rangle. \tag{2.5}$$

The neutral-kaon states with a definite lifetime $K_{S,L}$ are given in terms of the strangeness ±1 basis, using CPT invariance, as

$$|K_{S,L}\rangle=\frac{1}{[2(1+|\epsilon|^2)]^{1/2}}[(|K^0\rangle\pm|\overline{K}^0\rangle)+\epsilon(|K^0\rangle\mp|\overline{K}^0\rangle)]. \qquad (2.6)$$

The CP impurity in the corresponding $|K_S\rangle$ or $|K_L\rangle$ state is given by the complex parameter ϵ, whose real part measures its lack of orthogonality,

$$\langle K_S|K_L\rangle=\frac{2\,\mathrm{Re}\epsilon}{(1+|\epsilon|^2)}, \qquad (2.7)$$

and is expected to be of the order of magnitude of 10^{-3}. In fact $\mathrm{Re}\epsilon$ can be experimentally measured, because the charge asymmetry parameter δ_L for the semileptonic decays of the neutral kaon can be written as

$$\delta_L=\frac{N_+-N_-}{N_++N_-}=2\,\mathrm{Re}\epsilon \qquad (2.8)$$

if one assumes that the $\Delta S=-\Delta Q$ amplitudes are small compared to the $\Delta Q=\Delta S$ amplitudes (its ratio, x, is consistent with zero; $\mathrm{Re}x=-0.003\pm0.027$, $\mathrm{Im}x=-0.005\pm0.038$). We shall use therefore $\mathrm{Re}\epsilon$ as an experimental input, its value[7] being

$$2\,\mathrm{Re}\epsilon=(3.34\pm0.12)\times10^{-3}. \qquad (2.9)$$

The experimental parameters we shall use are

$$\eta_{ij}=\frac{A(K_L\to\pi^i\pi^j)}{A(K_S\to\pi^i\pi^j)}=\frac{A^L_{ij}}{A^S_{ij}}=\frac{(A_{ij}-\overline{A}_{ij})+\epsilon(A_{ij}+\overline{A}_{ij})}{(A_{ij}+\overline{A}_{ij})+\epsilon(A_{ij}-\overline{A}_{ij})}, \qquad (2.10)$$

where

$$A_{ij}\equiv A(K^0\to\pi^i\pi^j),\quad \overline{A}_{ij}\equiv A(\overline{K}^0\to\pi^i\pi^j)$$

and (i,j) are either $(+,-)$ or $(0,0)$.

The current experimental values for η_{+-} and η_{00} as given in Ref. 7 are

$$\begin{aligned}&|\eta_{+-}|=(2.30\pm0.035)\times10^{-3},\\&|\eta_{00}/\eta_{+-}|=1.013\pm0.046,\\&\phi_{+-}=49.4^\circ\pm1.7^\circ,\\&\phi_{00}-\phi_{+-}=4^\circ\pm13^\circ.\end{aligned} \qquad (2.11)$$

The strong interaction of the final $\pi\pi$ system is taken into account, as usual, by the phase shifts in the isospin channels. The phase shifts can be determined from the experimental decay rates of $K_S\to2\pi$ and $K^+\to\pi^+\pi^0$ and can also be extracted from the $\pi\pi$ phase-shift analysis. We shall use our result of Ref. 8, $\delta_2-\delta_0=53^\circ$, which is inside the error of a recent determination in the same spirit which includes radiative corrections,[9] and which is the only result compatible with the phase-shift analysis, summarized in Ref. 9. We do not use the Wu-Yang[10] gauge, but rather let[11,8]

$$\begin{aligned}&\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K^0_1\rangle=a_0e^{i\delta_0},\\&\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K^0_2\rangle=ia_0\gamma e^{i\delta_0}\end{aligned} \qquad (2.12)$$

instead of the more standard former choice $\gamma=0$. The parameters a_0 and γ are real numbers because of CPT. K^0_1 (K^0_2) is the even (odd) CP eigenstate. Note also that the CP impurity parameter in the Wu-Yang gauge, ϵ', can be expressed as

$$\epsilon'=\epsilon_0, \qquad (2.13)$$

whereas in our phase convention the corresponding impurity parameter, denoted by ϵ [see Eq. (2.6)], is given by

$$\epsilon=\epsilon_0-i\gamma, \qquad (2.14)$$

where ϵ_0 is the ratio

$$\epsilon_0=\frac{\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K_L\rangle}{\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K_S\rangle}. \qquad (2.15)$$

The relation between the impurities in both gauges is

$$\epsilon'=\epsilon+i\gamma. \qquad (2.16)$$

Note also that transforming the K^0 and $\overline{K}^0$ states as

$$\begin{aligned}&K^0\to(K^0)'=e^{i\xi}K^0\simeq(1+i\xi)K^0,\\&\overline{K}^0\to(\overline{K}^0)'=e^{-i\xi}\overline{K}^0\simeq(1-i\xi)\overline{K}^0\end{aligned} \qquad (2.17)$$

(in going from one phase convention to another), one obtains for the infinitesimal ξ the value

$$\xi=\gamma=\frac{\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K^0_2\rangle}{\langle2\pi,T=0|\mathcal{H}_{\mathrm{NL}}|K^0_1\rangle}, \qquad (2.18)$$

keeping the same value for ϵ_0 in both gauges. The expressions above then give the connection between the two gauges.

III. COMPUTATION OF THE AMPLITUDES

For the dynamical computation of the relevant meson decays we apply the work of Böhm, Joos, and Krammer, extended to include symmetry breaking in a simple way. The mesons are considered as quark-antiquark bound states, represented by Bethe-Salpeter (BS) amplitudes $\chi(r,P)$, where r is the relative momentum and P the bound-state momentum, which satisfy the (BS) equation

$$S^{-1}(P_1)\chi(r,P)\bar{S}^{-1}(P_2)=i\int d^4r'K(r,r',P)\chi(r',P). \tag{3.1}$$

The basic assumptions for the former BS equation are large quark masses ($m>1$ GeV), unrenormalized propagators $S^{-1}=\gamma\cdot P-m$, and a convolution-type, energy-independent kernel. From the requirement that the spectrum derived from Eq. (3.1) should reproduce the nonrelativistic one, the spin structure turns out to be pseudoscalar + vector − scalar.[12] The equation is solved in the Wick-rotated form, for a smooth kernel approximated by an oscillator form, as an expansion in $(M/m)^2$ of the bound-state zero-mass solution. The (BS) amplitudes for the pseudoscalar mesons are

$$\chi_{0^-}(P,r)=\frac{4}{\sqrt{3\beta}}\left(1+\frac{\not P}{M}\right)\gamma_5\exp\left(-\frac{r^2}{2\sqrt{\beta}}\right)|q\bar{q}\rangle. \tag{3.2}$$

$2\sqrt{\beta}=1$ GeV^2 is obtained from the meson level spacing and $|q\bar{q}\rangle$ is the SU(3) quark content of the mesons. For the SU(3) breaking in weak meson decays, it has been shown[5] recently that it can be explained by an additional term

$$\delta u(K)(K-K')^\mu s(K'). \tag{3.3}$$

The coefficient δ of the SU(3) symmetry-breaking term is proportional to the mass difference between the strange and nonstrange quarks. δ vanishes in the limit of exact SU(3) symmetry, which implies that (3.3) is a current of the first kind. The numerical value of δ which modifies the bare Cabibbo current in hadronic decays can be fixed from the experimental[13] decay rate $\Gamma(K^+\to\pi^+\pi^0)=1.12\times10^{-17}$ GeV, which yields

$$\delta=-0.144\ \text{GeV}^{-1}. \tag{3.4}$$

We would like to remark that the subsequent analysis is rather weakly dependent on the numerical value of δ since only ratios of amplitudes enter.

The amplitudes for $K_S\to\pi^+\pi^-$ and $K_S\to\pi^0\pi^0$ are then given by the traces of the diagrams shown in Figs. 1(a) and 1(c) and Figs. 1(b) and 1(c), respectively, which are

$$g_a=-e^{i\phi}\left[f_K(m_K^2-m_\pi^2)+\frac{16\delta\sqrt{\beta}}{3\sqrt{3}\,\pi}m_\pi^2\right]\exp\left(\frac{m_\pi^2}{16\sqrt{\beta}}\right)\frac{G}{\sqrt{2}}\sin\theta_C\cos\theta_C\,, \tag{3.5}$$

$$g_b=\left[\frac{e^{i\phi}+e^{-i\omega}}{2}f_K(m_K^2-m_\pi^2)-\frac{16\sqrt{\beta}M}{3\sqrt{3}\,\pi}(e^{i\phi}-e^{-i\omega})+\frac{4\sqrt{\beta}\,\delta}{3\sqrt{3}\,\pi}\left(2\sqrt{\beta}-\frac{5m_K^2}{8}+\frac{m_\pi^2}{16}\right)e^{i\phi}\right]\exp\left(\frac{m_\pi^2}{16\sqrt{\beta}}\right)\frac{G}{\sqrt{2}}\sin\theta_C\cos\theta_C\,, \tag{3.6}$$

$$g_c=\left\{\frac{16M\sqrt{\beta}}{3\sqrt{3}\,\pi}(e^{i\phi}-e^{-i\omega})+\frac{\delta\sqrt{\beta}}{3\sqrt{3}\,\pi}\left[\tfrac{3}{4}(m_K^2-4m_\pi^2)+8\sqrt{\beta}\right]e^{i\phi}\right\}\exp\left(\frac{m_K^2}{16\sqrt{\beta}}\right)\frac{G}{\sqrt{2}}\sin\theta_C\cos\theta_C\,. \tag{3.7}$$

The quark model predicts the existence of a $T=0$ $\pi\pi$ resonance (ϵ meson), with (BS) amplitude

$$\chi_{0^+}(P,r)=\frac{4\pi\sqrt{2}}{3\beta^{3/4}}\left[\not r-\frac{\not P(P\cdot r)}{m^2}+\frac{\not P\cdot\not r}{M}-\frac{P\cdot r}{M}+\frac{r^2}{M}-\frac{(P\cdot q)^2}{Mm^2}\right]\exp\left(-\frac{r^2}{2\sqrt{\beta}}\right)|q\bar{q}\rangle. \tag{3.8}$$

Experimentally there is in fact indication of the existence of such a (broad) resonance in the 600–700 MeV mass region, but with a huge uncertainty. We need such a resonance to account for the total $K_S\to2\pi$ width and the values which best reproduce this width and the K_S-K_L mass difference turns out to be $m_\epsilon\sim\Gamma_\epsilon=600$ MeV. The contribution from the diagram shown in Fig. 1(d) is given by

$$g_d=\left[\frac{16\beta^{5/4}}{3\sqrt{3}^{\,2}}\left(1-\frac{m_K^2}{4m_\epsilon^2}\right)e^{i\phi}\delta-\frac{4}{M}(e^{i\phi}-e^{-i\omega})\right]\frac{g_{\epsilon\pi\pi}}{m_K^2-m_\epsilon^2}\frac{G}{\sqrt{2}}\sin\theta_C\cos\theta_C\,, \tag{3.9}$$

where a real propagator $1/(m_\epsilon^2-m_K^2)$ has been used to satisfy *CTP* and unitarity.

The physical amplitudes for the decays $K_S\to\pi^+\pi^-$ and $K_S\to\pi^0\pi^0$ are

$$A_{+-}=\left(\frac{2}{3}a_{+-}-\frac{\sqrt{2}}{3}a_{00}\right)e^{i\delta_0}+\tfrac{1}{3}(a_{+-}+\sqrt{2}\,a_{00})e^{i\delta_2}, \tag{3.10}$$

$$A_{00}=\left(\frac{\sqrt{2}}{3}a_{+-}+\frac{a_{00}}{3}\right)e^{i\delta_0}+\frac{1}{\sqrt{3}}(a_{+-}-\sqrt{2}\,a_{00})e^{i\delta_2},$$

where a_{+-} and a_{00} are the "bare" amplitudes determined from Eqs. (3.5), (3.6), (3.7), and (3.9):

$$a_{+-}=(g_a+g_b+g_d),$$

$$a_{00}=\frac{1}{\sqrt{2}}(g_b-g_c-g_d). \tag{3.11}$$

Their numerical expressions in terms of the phase angles ϕ and ω are (in keV)

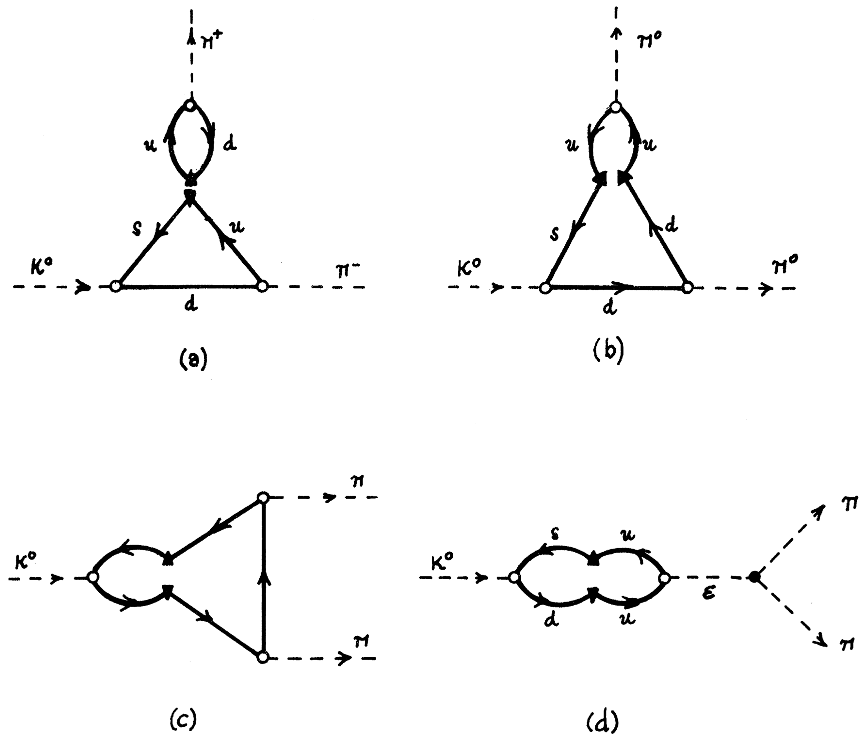

FIG. 1. Diagrams for the decay $K \to \pi\pi$. Solid lines represent the quarks u, d, and s and dashed lines the mesons. Circles stand for BS amplitudes and triangles for the currents.

$$A_{+-} = [(-1.63 - i1.13)e^{i\phi} + (1.35 + i0.77)e^{-i\omega}], \tag{3.12}$$

$$A_{00} = [(0.50 - i0.80)e^{i\phi} + (-0.68 + i0.54)e^{-i\omega}].$$

IV. APPLICATION TO THE PHASE-ANGLES *CP*-VIOLATION MODELS

The amplitudes for the $K \to 2\pi$ decays of the preceding section have been cast in the form

$$A_{+-} = A[(m+in)e^{i\phi} - (c+id)e^{-i\omega}], \tag{4.1a}$$

$$\overline{A}_{+-} = A[(m+in)e^{-i\phi} - (c+id)e^{i\omega}], \tag{4.1b}$$

$$A_{00} = A[(m_0+in_0)e^{i\phi} - (c_0+id_0)e^{-i\omega}], \tag{4.1c}$$

$$\overline{A}_{00} = A[(m_0+in_0)e^{-i\phi} - (c_0+id_0)e^{i\omega}], \tag{4.1d}$$

where, as we said before, ϕ and ω are the *CP*-violating phases of the Glashow-type models and $m, \ldots, d_0$ are given by (3.12).

In order to fix the *CP*-violating theoretical parameters ϕ and ω we shall express the amplitudes $A(K_S \to \pi^i\pi^j)$ and $A(K_L \to \pi^i\pi^j)$ and the related experimental parameters η_{ij} as a handy function of ϕ, ω, and ϵ. Then we use the best measured quantities $|\eta_{+-}|$, Φ_{+-}, and $\mathrm{Re}\,\epsilon$ to get a system of equations which fixes the phases ϕ and ω and $\mathrm{Im}\,\epsilon$, and with these values the remaining experimental parameter η_{00} is predicted.

The experimental parameter η_{ij} is given, according to (2.10), by

$$\eta_{+-} = \frac{(\epsilon\cos\phi + i\sin\phi) - (\alpha+i\beta)(\epsilon\cos\omega - i\sin\omega)}{(\cos\phi + i\epsilon\sin\phi) - (\alpha+i\beta)(\cos\omega - i\epsilon\sin\omega)}, \tag{4.2}$$

where

$$\alpha + i\beta = \frac{c+id}{m+in} = 0.78 + i0.07. \tag{4.3}$$

For η_{00} one gets the same expression with α_0 and β_0 instead of α and β, given by

$$\alpha_0 + i\beta_0 = \frac{c_0+id_0}{m_0+in_0} = 0.87 + i0.38. \tag{4.4}$$

Neglecting second-order terms in the *CP*-violating phases, we have from (4.2)

$$\epsilon = \eta_{+-} - i\frac{\phi + Z\omega}{1-Z}, \tag{4.5}$$

where $Z = \alpha + i\beta$.

The real and imaginary parts of Eq. (4.5) lead to

$$(\phi+\omega)=\frac{(1-\alpha)^2+\beta^2}{\beta}(\mathrm{Re}\,\epsilon-\mathrm{Re}\,\eta_{+-})$$

and (4.6)

$$\mathrm{Im}\,\epsilon-\omega=\mathrm{Im}\,\eta_{+-}-\frac{1-\alpha}{\beta}(\mathrm{Re}\,\epsilon-\mathrm{Re}\,\eta_{+-}).$$

Any conclusion based on values of $\mathrm{Re}\,\epsilon$ and $\mathrm{Re}\,\eta_{+-}$ is of course independent of the phase convention, whereas the value of $\mathrm{Im}\,\epsilon$ is gauge-dependent. Since we are considering *CP* violation in the weak Hamiltonian, the most natural convention is not the one of Wu and Yang,[10] where the amplitude $\langle 2\pi, T=0|\mathcal{H}_{\rm NL}|K_2^0\rangle$ is set equal to zero, reducing in fact all *CP* violation to the mass matrix. We want instead in our analysis a nonzero value for the former *CP*-violating amplitude.

One can see that $\phi+\omega$ is proportional to the difference $(\mathrm{Re}\,\epsilon-\mathrm{Re}\,\eta_{+-})$, where the proportionality factor, given by the values of α and β quoted in (4.3), is 0.77. The experimental value[8] of

$$(\mathrm{Re}\,\epsilon-\mathrm{Re}\,\eta_{+-})=(1.7\pm0.8)\times10^{-4}$$

has a 50% uncertainty, so that we have at least the same indeterminacy for the sum of the phases, which would be

$$\phi+\omega=(-1.3\pm0.6)\times10^{-4}. \tag{4.7}$$

The phases are then nearly opposite, but as long as $\mathrm{Re}\,\epsilon\neq\mathrm{Re}\,\eta_{+-}$, the exact relation $\phi=-\omega$ is impossible. In fact, the conclusion that $\phi+\omega=0$ is excluded if $\mathrm{Re}\,\epsilon\neq\mathrm{Re}\,\eta_{+-}$ holds for any current-current model. Should experiments confirm that $\mathrm{Re}\,\eta_{+-}=\mathrm{Re}\,\epsilon$, then $\phi+\omega=0$ necessarily, and then our Eq. (4.5) would become simply

$$\eta_{+-}=\eta_{00}=\epsilon-i\omega. \tag{4.8}$$

In this case one can illustrate again the difference between our phase convention and the Wu-Yang gauge. In our case, using Eq. (2.14) and $\epsilon_0\simeq\frac{1}{3}(2\eta_{+-}+\eta_{00})$, one gets

$$\eta_{+-}=\eta_{00}=\epsilon_0-i(\gamma+\omega)=\epsilon_0, \tag{4.9}$$

i.e.,

$$\gamma=-\omega, \tag{4.10}$$

whereas in the Wu-Yang case

$$\eta_{+-}=\eta_{00}=\epsilon_0=\epsilon', \tag{4.11}$$

where ϵ_0 and γ were defined in Eqs. (2.15) and (2.12), respectively, and where ϵ' is the *CP* impurity parameter in the Wu-Yang gauge. So the real parts of η_{+-} are given by the same expression in both gauges,

$$\mathrm{Re}\,\eta_{+-}=\mathrm{Re}\,\epsilon=\mathrm{Re}\,\epsilon_0 \quad \text{(in our gauge)},$$

$$\mathrm{Re}\,\eta_{+-}=\mathrm{Re}\,\epsilon'=\mathrm{Re}\,\epsilon_0 \quad \text{(in the Wu-Yang gauge)}$$

whereas the imaginary part of η_{+-} is expressed by different expression, but which give of course the same value

$$\mathrm{Im}\,\eta_{+-}=\mathrm{Im}\,\epsilon-\omega=\mathrm{Im}\,\epsilon_0 \quad \text{(in our gauge)},$$

$$\mathrm{Im}\,\eta_{+-}=\mathrm{Im}\,\epsilon'=\mathrm{Im}\,\epsilon_0 \quad \text{(in the Wu-Yang gauge)}.$$

On the other hand, the situations $\phi=\omega$, $\phi=0$ and $\omega\neq0$, and $\omega=0$ and $\phi\neq0$ are allowed with $\phi=\omega=-6.6\times10^{-5}$, $\omega=-1.3\times10^{-4}$, and $\phi=-1.3\times10^{-4}$, respectively. In these cases the equations can be solved exactly, and the results of the model are, for all three cases,

$$\mathrm{Im}\,\epsilon=2.23\times10^{-3},\quad \mathrm{Re}\,\eta_{00}=2.04\times10^{-3},\quad \phi_{00}=46.4^\circ, \tag{4.12}$$

yielding for $|\eta_{00}/\eta_{+-}|$ the value 1.28, which does not rule out the model in view of the large errors involved, as shown in (2.11).

Finally, we remark that a recent theoretical analysis[14] of the K° meson self-energy matrix gives a value for $\phi+\omega$ one order of magnitude smaller than in Eq. (4.7).

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