NEW RESULTS ON RARE NEUTRAL KAON RARE DECAY FROM FERMILAB EXPERIMENT E-731[1]

Yau W. Wah

The Enrico Fermi Institute and the Department of Physics

University of Chicago

Chicago, Illinois 60637, USA

ABSTRACT

We reported preliminary results on i) the determination of the branching ratio of $K_L \rightarrow \pi^0\pi^\pm e\nu$; ii) a limit on the branching ratio of $\pi^0 \rightarrow e^+e^-$; and iii) an update of the decay $K_L \rightarrow \pi^0\gamma\gamma$.

Introduction

Fermilab Experiment E-731 was primarily designed to measure the parameter ϵ'/ϵ of the $K \rightarrow 2\pi$ decay. Because of the loose triggers and high kaon flux of the experiment, a number of rare neutral kaon decay modes could also be measured [2]. Fig.1 shows the apparatus schematics and its details were described elsewhere [3].

$K_L \rightarrow \pi^0 \pi^\pm e\nu$ (K^0_{e4})

The kinematics of K_{e4} is described by the following five variables [5] which are (i) $m_{\pi\pi}$, the effective mass of the dipion, (ii) $m_{e\nu}$, the effective mass of the dilepton, (iii) θ_π, the angle of the charged pion in the c.m system of the pions with respect to the direction of flight of the dipion in the K rest system, (iv) θ_l, the angle of the $e^\pm$ in the c.m. system of the

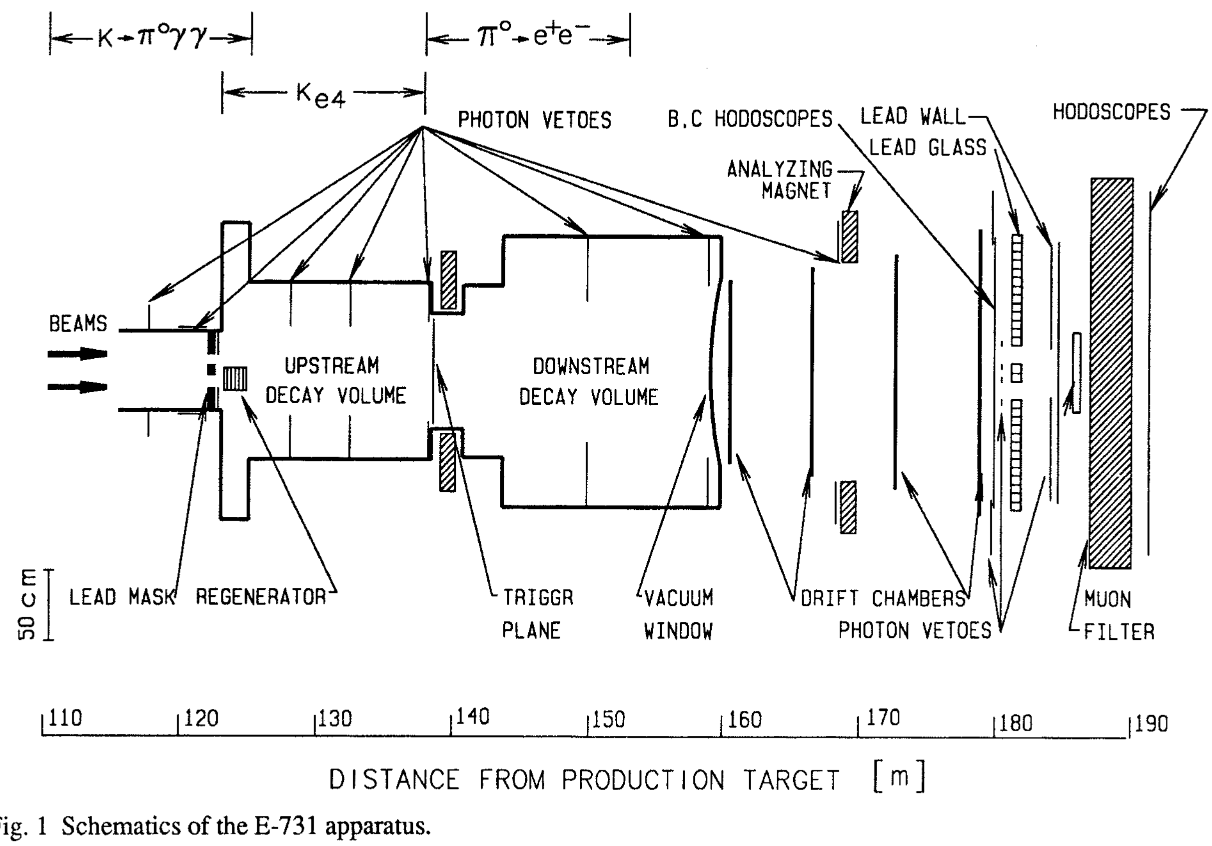

Fig. 1 Schematics of the E-731 apparatus.

leptons with respect to the direction of flight of the dilepton in the K rest system, and (v) ϕ, the angle between the plane formed by the pions in the K rest system and the corresponding plane formed by the leptons. The differential distribution of the above variables could be written as a linear sum of nine 'form factors'. It was pointed out [6] that one of them could be a test of non-standard model CP violation. Moreover, unlike the $K^+ \to \pi\pi e\nu$ which is s-wave dominated, K^0_{e4} is p-wave dominated and its branching ratio tests current algebra [4] in the limit of $m_\pi \to 0$.

In the E-731 data set, events with two opposite charged tracks (each with momenta >2.5 GeV) matched with calorimeter clusters and identified one as $\pi^\pm$ and the other as $e^\pm$ were selected. The π^0 was then recontructed from two extra calorimeter clusters each with > 2.1 GeV. To suppress backgrounds from $K_L \to \pi^+\pi^-\pi^0$ with misidentified pions and gamma accidentals overlapped with $K_{e3\gamma}$, candidate events were rejected if they were consistent as $K_L \to \pi^+\pi^-\pi^0$ from target or if $2p_e \cdot p_\gamma < 100$ MeV. A cut on the high solution of the calculated neutrino transverse momentum was also applied. Fig. 2 shows the $m_{\gamma\gamma}$ distribution after all cuts. There are 929 events within 10 MeV of m_π. The remaining background is estimated to be 149 events. With 1.38% acceptance (assume flat phase space) and 1.01×10^9 K_L exposure, the nett 780 events gives a B.R. of $(5.6 \pm 0.2) \times 10^{-5}$. The combined systematics from pion showers, uncertainty in accidental losses, background subtraction, flux, and acceptances is 7.6% which contributes ±0.4 error to the B.R.. Only 15 events were observed before [7]. Fig.3 shows the three angular distributions of the data compared to phase space Monte Carlo. It is quite possible to look for non-standard CP violation in this mode in a near future high flux kaon experiment [2].

$\pi^0 \to e^+e^-$

Basically this decay could be described by a fourth order electromagnetic diagram and its branching ratio is calculated to be $>4.8 \times 10^{-8}$, the so-called unitarity limit. The published measurements [8] were $(1.8 \pm 0.7) \times 10^{-7}$ and $< 1.2 \times 10^{-7}$. A new technique was used to search

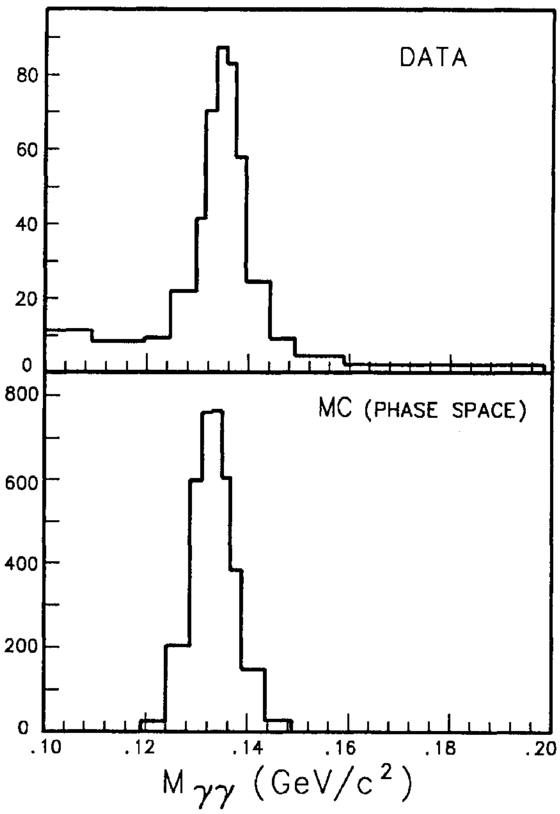

Fig. 2 Invariant $\gamma\gamma$ mass distribution of $K_L \to \pi^\pm\pi^0 e\nu$ after all cuts.

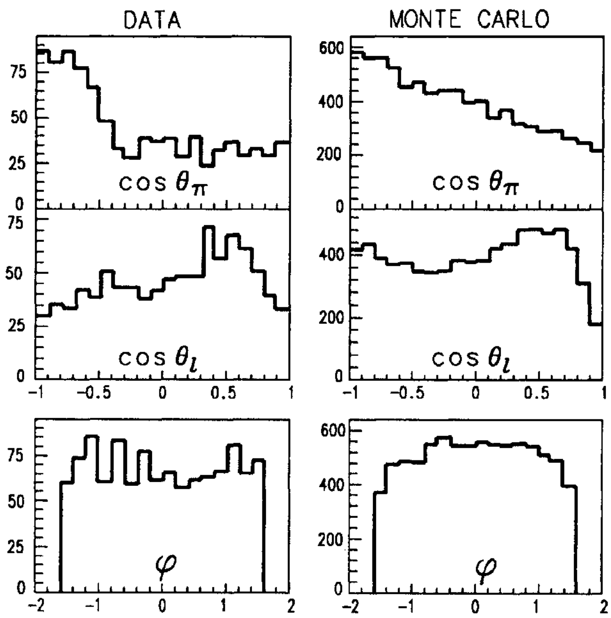

Fig. 3 Angular distributions of $K_L \to \pi^\pm\pi^0 e\nu$ after all cuts

for this decay in the $K \to 3\pi^0$ decay (the 'Neutral Sample'): any one of the three π^0 from the 15×10^6 reconstructed $K \to 3\pi^0$ decay could decay to a e^+e^- pair

Thus the signature of $\pi^0 \rightarrow e^+e^-$ is: two tracks (each with momentum >2.5 GeV) matched with the leadglass showers (>1 GeV) and reconstructed as a π^0; the remaining four extra clusters reconstructed as $2\pi^0$ (χ^2<8); and all six final state particles reconstructed as a kaon coming from the target (P_t^2 < 300 MeV^2).

We also search for the decay through our multi-track events (the 'Track Sample') via $K_L \rightarrow \pi^+\pi^-\pi^0$ if $\pi^0 \rightarrow e^+e^-$. The cuts are four tracks found with pions and a electron pair. The whole event has to be consistent with a kaon from target.

The background of the 'Neutral Sample' came from (i) single dalitz of π^0; (ii) double dalitz of π^0; and (iii) two π^0 single dalitz. In the 'Track Sample', only (i) and (ii) contributed. For (i) and (ii), $m_{ee} < m_\pi$ and for (iii), m_{ee} can be > m_π, but suppressed since $m_{\gamma\gamma}$ could have π^0 mass accidentally. With cuts that demand: no photon vetos activities; no extra track segment points to the reconstructed vertex; and no more than 120 MeV of energy in any of the leadglass blocks that do not belong to the electromagnetic showers, there is no background. Fig.4 and 5 shows the scatterplot of M_{ee} vs $M_{3\pi}$ of the 'Neutral Sample' and 'Track Sample' respectively after all cuts.

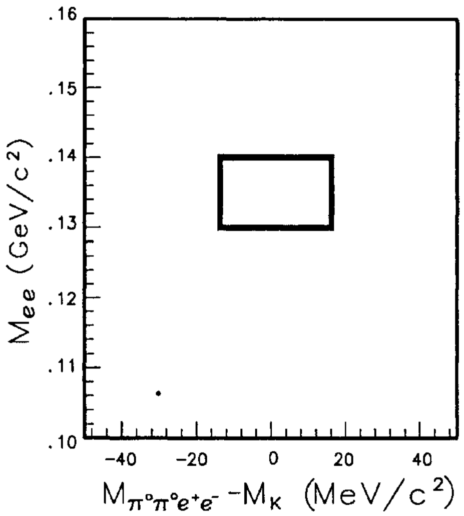

Fig.4 Reconstructed e^+e^- pair mass vs. kaon mass of the $\pi^0 \rightarrow e^+e^-$ search in $K_L \rightarrow 3\pi^0$ decay.

From the two samples, we found a combined limit of B.R. ($\pi^0 \rightarrow e^+e^-$) < 2.3 x 10^{-7} (90% c.l.). This 'tagged pion' technique has good acceptance (2.4%)

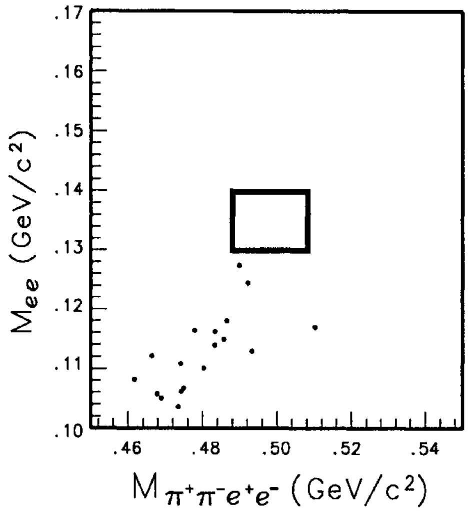

Fig.5 Reconstructed e^+e^- pair mass vs. kaon mass of the $\pi^0 \rightarrow e^+e^-$ search in $K_L \rightarrow \pi^+\pi^-\pi^0$ decay.

and is, unlike the previous experiments, background free. It will be exploited much further in another high kaon flux experiment [2].

$K_L \rightarrow \pi^0\gamma\gamma$

The decay $K_L \rightarrow \pi^0\gamma\gamma$ provides the handle on the size of the CP conserving amplitude of the most interesting decay $K_L \rightarrow \pi^0e^+e^-$, of which ε'/ε could be $O(1)$. Chiral perturbation theories[9] predict the decay with a branching ratio of ~ 6.8 x 10^{-7} with a characteristic $\gamma\gamma$ mass distribution peaked at ~ 325 MeV^2. If true, then the CP conserving amplitude of $K_L \rightarrow \pi^0e^+e^-$ could contribute as small as <10^{-13} to the branching ratio. This will ease the untangling of the CP violating amplitudes which may contribute as much as ~10^{-11}.

Details of the analysis can be found elsewhere [3,10]. Fig.6 shows the $\gamma\gamma$ invariant mass of the candidate events after all cuts. A characteristic feature in this distribution is a prominent 'double fusion' peak appearing at about $m_{\gamma\gamma}$ = 270 MeV which arises when two π^0's in $K_L \rightarrow 3\pi^0$ decays are superimposed: each photon from one π^0 overlaps with a photon from the other π^0 with an invariant mass threshold of twice the π^0 mass. Same figure also shows the Monte Carlo prediction of the backgrounds to $\pi^0\gamma\gamma$ decay coming from the $3\pi^0$ and $2\pi^0$ modes. The Monte Carlo distributions are normalized by means of a sample of fully reconstructed $K_L \to 2\pi^0$ decays observed simultaneously.

For $m_{\gamma\gamma} > 300$ MeV, 66 candidates events was found with a predicted background of 61 events from 2π and 3π decays. For $m_{\gamma\gamma} > 280$, 111 candidates events was found with 81 predicted background. The $\pi^0\gamma\gamma$ acceptance is 2.4% and the normalization is provided by the $K_L \to 2\pi^0$ decays observed. Assuming the $m_{\gamma\gamma}$ spectrum follows that predicted by chiral perturbation theory, then (i) B.R.($K_L \to \pi^0\gamma\gamma$, $m_{\gamma\gamma}$ >300 MeV) < 1.3×10^{-6} (90% c.l.) and; (ii) B.R.($K_L \to \pi^0\gamma\gamma$, $m_{\gamma\gamma} >$ 280MeV)=(1.7 ± 0.7) x 10^{-6} are obtained. Errors are statistical only. These results are preliminary and consistent with other recent measurements [10].

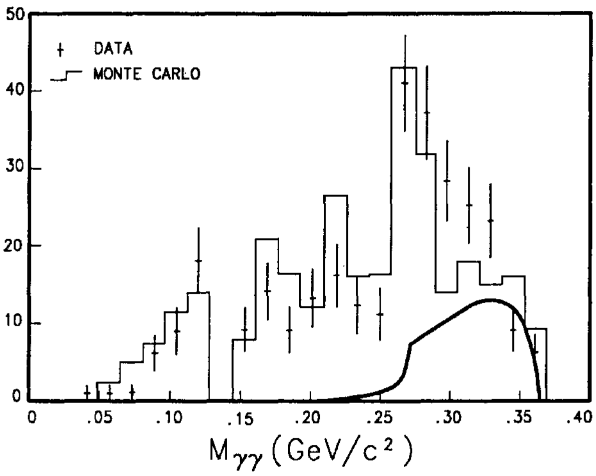

Fig. 6 Invariant $\gamma\gamma$ mass distribution of $K_L \to \pi^0\gamma\gamma$ distribution. The thick solid curve represents a chiral theory predicted[9] branching ratio of 2.7×10^{-6}.

Summary

Preliminary results are reported on:
i) B.R.($K_L \to \pi^0\pi^\pm e^- \nu$) = ($5.6 \pm 0.2 \pm 0.4$) x 10^{-5};
ii) a limit on B.R.($\pi^0 \to e^+e^-$)<2.3×10^{-7}; and
iii) B.R.($K_L \to \pi^0\gamma\gamma$, $m_{\gamma\gamma} >$ 280MeV)=(1.7 ± 0.7) x10^{-6}.

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DISCUSSION

Q. **K. Kleinknecht** *(Univ. Mainz)*: Concerning the branching ratio for $k_L \to \pi^0\gamma\gamma$ you reported: on how many observed events is it based?

A. **Y. W. Wah**: The reported branching ratio, $B(K_L \to \pi^0\gamma\gamma) = (1.7\pm0.7)\times10^{-6}\tau s$ is based on 111 events ($\mu_{ee} > 280$ MeV) with 81 events background, the error on background is 6.4 evts.

RECENT RESULTS ON CP-VIOLATION FROM FERMILAB EXPERIMENT E731

E731 Collaboration (*Chicago-Elmhurst-Fermilab-Princeton-Saclay*)
presented by

YEE B. HSIUNG
Fermi National Accelerator Laboratory, P.O.Box 500
Batavia, IL 60510, U.S.A.

ABSTRACT

We report the current status of the analysis for the CP-violation parameters, ε'/ε, from the entire data sample of Fermilab experiment E731. A new measurement on the CP-violation parameter $\eta_{+-\gamma}$ was extracted from the K_L-K_S interference of the $\pi^+\pi^-\gamma$ decay mode at downstream of a regenerator. The preliminary result, $|\eta_{+-\gamma}| = 0.0020 \pm 0.0002(\text{stat}) \pm 0.0003(\text{syst})$, is consistent with $|\eta_{+-}|$ from $\pi^+\pi^-$ decay mode. Some results on the measurements of the branching ratio for $K_L \to \pi^+\pi^-\gamma$ and $K_S \to \pi^+\pi^-\gamma$ are presented in here. For the very first time, the quadratic decay parameter of $K_L \to 3\pi^0$ has been measured from the $3\pi^0$ Dalitz plot. Our result, $b = (-0.6 \pm 1.4) \times 10^{-3}$, is consistent with zero, indicating a flat Dalitz distribution. It is inconsistent with the fitted result, $b = (-8.3 \pm 2.4) \times 10^{-3}$, from Devlin and Dickey's review.

STATUS OF ε'/ε

Fermilab experiment E731 was designed to measure the "direct" CP violation parameters ε'/ε, which can be expressed as a double ratio of four decay modes of the $K_{L,S} \to \pi^+\pi^-(\pi^0\pi^0)$, *i.e.*

$$R \equiv \frac{|\eta_{+-}|^2}{|\eta_{00}|^2} = \frac{\Gamma(K_L \to \pi^+\pi^-)/\Gamma(K_S \to \pi^+\pi^-)}{\Gamma(K_L \to \pi^0\pi^0)/\Gamma(K_S \to \pi^0\pi^0)}$$

$$\cong 1 + 6\,Re(\varepsilon'/\varepsilon),$$

where ε describes the K^0-$\bar{K}^0$ mixing and ε' describes the "direct" CP-violation from the decay amplitude. The E731 apparatus has been shown elsewhere [1]. Basically two nearly identical *side-by-side* K_L beams were brought into the decay region, where a regenerator was placed alternately in one of the K_L beams to produce "K_S". Therefore *simultaneous detection* of both K_L and K_S decays in the same decay region with the same detector became possible. The experiment was designed to minimize systematic uncertainty and in particular to be immune to changes in accelerator cycles and other parameters over which the experimenter has little control, such as *beam intensity and asymmetry, electronic drifts, phototube gain shifts as well as inefficiencies, etc.*

To get the double ratio R, *both* $2\pi^0$ and $\pi^+\pi^-$ decays need to be recorded for *both* K_L and K_S modes. The experiment had its major run in 1987/88. For the majority of the running, either both *charged* ($\pi^+\pi^-$) modes or both *neutral* ($2\pi^0$) modes were recorded simultaneously and this greatly reduced a variety of systematics. For the last 20% of the data, *all four modes* were taken at the same time, providing a best check on the systematics. It is on this data set that we have done the full analysis and published result [2] recently. We found $\varepsilon'/\varepsilon = -0.0004 \pm 0.0014$ (stat.) ± 0.0006 (syst.), which is consistent with zero and is two standard deviations below the earlier result from CERN NA31 [3]. From the same data, we also determined the CPT violating phase difference between η_{+-} and η_{00}: $\Delta\phi = -0.3° \pm 2.4°(\text{stat.}) \pm 1.2°(\text{syst.})$ [4]. This together with a result of similar precision from NA31 [5] in a dedicated experiment resolved a long-standing earlier discrepancy [6] on $\Delta\phi$ which was $12° \pm 6°$.

In the Standard Model, the value of ε'/ε is expected to be non-zero but its precise magnitude depends on, among other parameters, the value of the top quark mass: the higher the mass of the top, the smaller is ε'/ε [7]. Our result clearly favors higher values of m_t in ref. 7 as indicated by the top quark search ($m_t >$ 89 GeV) from the collider experiment [8].

We have made substantial progress in the analysis of the entire data. A result is expected soon in 1991. All of the remaining data has been reduced and the work at present is in the rather delicate calibration of the leadglass detector as well as in the understanding of the acceptance of the apparatus. In Fig. 1 we show the reconstructed $K_{L,S} \to \pi^0\pi^0$ invariant mass distribution for the entire data set, without the final calibration constants. The K_L signal has about 290k events in it (480 to 520 MeV/c^2) while the $3\pi^0$ background is less than 0.5% and is smooth and well understood. The K_S signal has about 970k events in it with a very small background. Since these events were taken simultaneously with the same leadglass detector, the line-shape of the mass peak is virtually identical between K_L and K_S. The incoherent K_S background in the K_L beam (the "crossover" K_S) was quite similar to the one in 20% of the data [2]. The exact amount of the background in the *neutral* mode will be determined once we have final leadglass calibration constants.

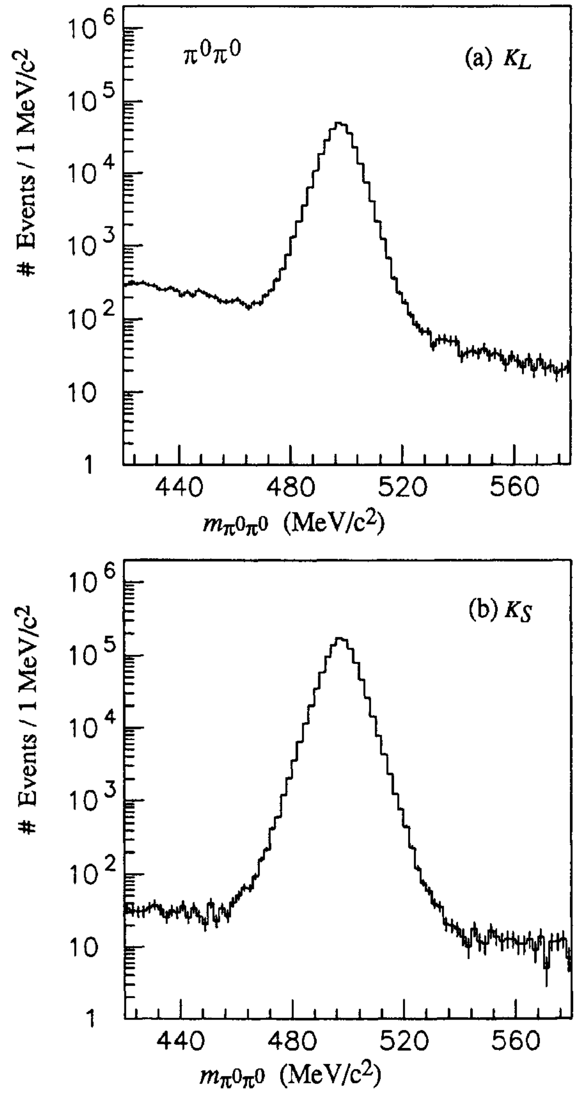

Fig. 1 $\pi^0\pi^0$ mass distribution (a) for K_L and (b) for K_S.

In Fig. 2 we show the reconstructed $K_{L,S} \rightarrow \pi^+\pi^-$ invariant mass distribution of all the data, in which there are 370k K_L and 1.2M K_S in between 484 to 512 MeV/c^2. Comparing this with the 20% of the data, the background level and the line-shape of coherent kaons in the *charged* mode are nearly identical to each other. The background in K_L, can be determined from p_t^2 distribution, is about 0.33% (mainly from the remaining K_{e3}) and in K_S is about 0.14% (from incoherently scattered K_S). With the above statistics, the final statistical error on ε'/ε will be about 0.0005. With better understanding of the detector acceptances, resolutions and energy scales from the current analysis by using the high statistics modes (80M K_{e3}, 10M $3\pi^0$ and 10M $\pi^+\pi^-\pi^0$), we would hope to achieve much smaller systematic uncertainty than that of the 20% data.

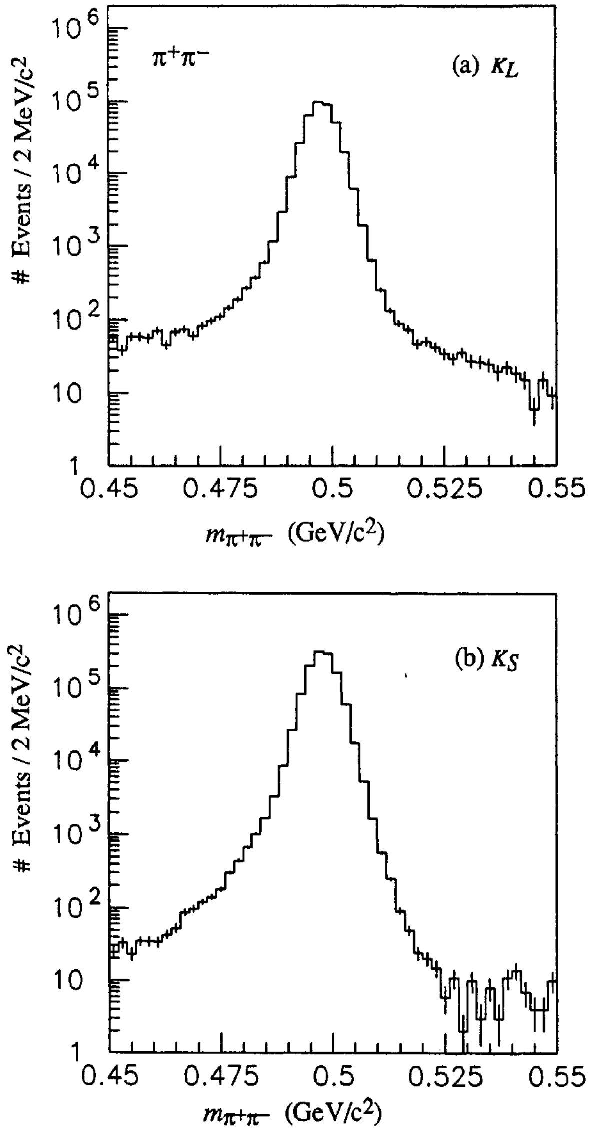

Fig. 2 $\pi^+\pi^-$ mass distribution (a) for K_L and (b) for K_S.

$K_{L,S} \rightarrow \pi^+\pi^-\gamma$

There are two processes by which the neutral kaons decay into $\pi^+\pi^-\gamma$: inner bremstrahlung (IB), where a decay occurs into two charged pions followed by the emission of a photon from one of the pion; and direct emission (DE) decay, where a photon emits from the $\pi\pi$ decay vertex (or one of the quark lines) [9]. Since it is essentially a two-pion decay, the IB process is CP-violating for K_L decays and CP-conserving for K_S, while in the DE process the opposite is true. In Fig. 2(b), the low side non-Gaussian tail of the $\pi^+\pi^-$

mass peak in K_S mainly comes from the radiative IB emission. Since the direct emission is CP-conserving in K_L decay. It is a possible background to ε'/ε. This lead us to measure it's branching ratio and studied it's contribution to ε'/ε [10].

Event reconstruction requires two opposite-sign charged tracks in the 4 sets of drift chambers and hodoscopes, as well as a separate photon cluster ($E_\gamma >$ 1.5 GeV) in the leadglass calorimeter. Pion selection was done by requiring that each track has $E/p < 0.80$ in the leadglass to reject event containing electrons. Each event was then required to pass fiducial and quality cuts as well as the following kinematic cuts:

(1) track momentum $p > 7$ GeV,
(2) kaon mass cut: $484 < m_{\pi\pi\gamma} < 512$ MeV/c^2,
(3) transverse momentum cut: $p_t^2 < 250$ (MeV/c)^2,
(4) kaon energy cut: $30 < E_K < 160$ GeV,
(5) decay z vertex within 27 m decay region,
(6) a kinematic variable cut: $P_0^2 < -0.025$ [11] to reject $\pi^+\pi^-\pi^0$ decays with one photon escaping the detection.

There were virtually no background remains after these selections.

Fig. 3 shows the center of mass photon energy ($E_\gamma{}^*$) spectrum overlayed for the K_L and K_S events, where the K_S spectrum has the inner bremstrahlung $1/k$ distribution and the K_L spectrum has the superposition of IB and DE ($\propto k^3$) distributions. The K_S spectrum has been properly normalized for the subtraction of IB in the K_L spectrum to extract the DE component. With $E_\gamma{}^* > 20$ MeV, we have 4620 K_S decays and 1552 IB and 2607 DE events in the K_L decays. The acceptance for K_S decay is 0.272 and for K_L decay

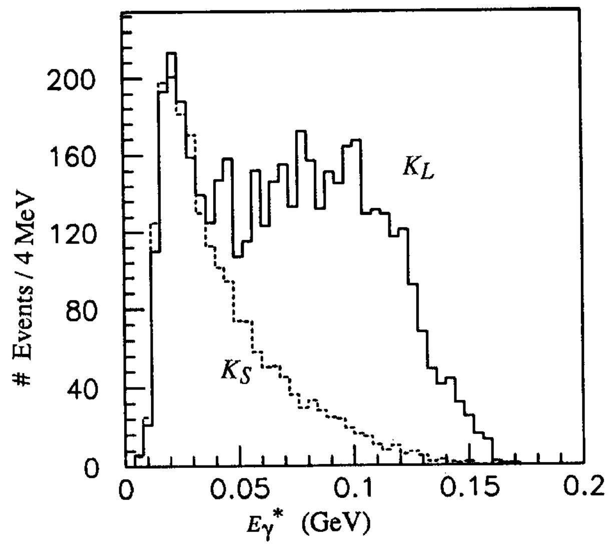

Fig. 3 $E_\gamma{}^*$ distribution of $\pi^+\pi^-\gamma$ for K_L decay (solid line) and K_S decay (dotted line).

is 0.130. The number of kaon decays into $\pi^+\pi^-$ was used to normalize the yield. The branching ratios of neutral kaon decay into $\pi^+\pi^-\gamma$ are $(4.35\pm.08) \times 10^{-5}$ for K_L and $(4.39\pm.07) \times 10^{-3}$ for K_S. The K_L decay has a branching ratio of $(2.98\pm.08) \times 10^{-5}$ into the DE mode and $(1.43\pm.08) \times 10^{-5}$ into the IB mode. This result is consistent with the previous measurement [11] and the theoretical predictions [9].

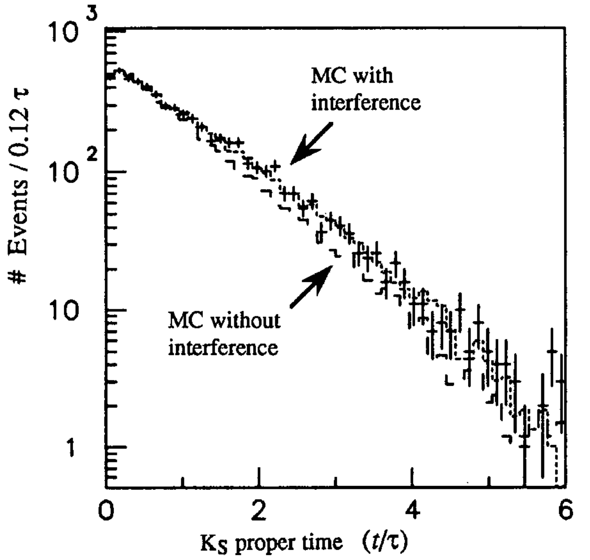

Fig. 4 Proper time distribution for $K_S \to \pi^+\pi^-\gamma$. Data with error bar were superimposed by a monte carlo with (dotted line) and without (broken line) interference term.

Since the coherent kaon beam behind the regenerator is the superposition of $K_L+\rho K_S$, one expects to see K_L and K_S interference in the proper time distribution, as $2|\rho||\eta_{+-\gamma}|\, e^{-t/2\tau} \cos(\Delta m t+\phi_\rho+\phi_\eta)$. The proper time distribution for the $K_S \to \pi^+\pi^-\gamma$ in the regenerated beam is shown in Fig.4. A monte carlo simulation with and without interference term (let $\eta_{+-\gamma} = \eta_{+-}$) are superimposed on the data. Clearly the data favors the monte carlo with interference. A preliminary fitting result has shown

$$|\eta_{+-\gamma}| = 0.0020 \pm 0.0002(\text{stat}) \pm 0.0003\ (\text{syst}),$$

which is consistent with the world average of the CP-violation parameter $|\eta_{+-}| = (2.268 \pm .023)\times 10^{-3}$ [12].

$K_L \to 3\pi^0$ QUADRATIC DECAY PARAMETER

The Dalitz plot parameters for the $K \to 3\pi$ decays can be parametrized by a series expansion as

$$|M|^2 \propto 1 + gY + \frac{g^2}{4}Y^2 + b(Y^2+\frac{X^2}{3}) + c(Y^2-\frac{X^2}{3}),$$

where $Y = (s_3-s_0)/m_\pi^2$ and $X = (s_2-s_1)/m_\pi^2$ are Dalitz variables [12]. In the $K_L \to 3\pi^0$ decay, the linear coefficient as well as the 1st and 3rd quadratic coefficient vanish, *i.e.* $g = 0$ and $c = 0$, because of the identical final state particles. The possible non-vanishing terms is then given by

$$|M_{000}|^2 \propto 1 + b(Y^2 + \frac{X^2}{3}).$$

Measurement on the $K_L \to 3\pi^0$ Dalitz plot has never been done before. Although the quadratic decay parameter b vanishes in the lowest order chiral perturbation theory, the combined fit for the $K^{\pm} \to 3\pi$ and $K_L \to \pi^+\pi^-\pi^0$ data by Devlin and Dickey [13] suggested an undesirable non-zero energy dependence $b = (-8.3 \pm 2.4)\times10^{-3}$, corresponding to a 2% drops from the center of Dalitz plot to the edge of the plot.

Based on a sample of 5M $K_L \to 3\pi^0$ decays in E731, we have performed an analysis on the $K_L \to 3\pi^0$ Dalitz distribution. Fig. 5 shows the data and a flat Dalitz monte carlo comparison of the event density distribution *vs* Dalitz angle Θ, where Θ is the angle between the event and Y-axis on the X-Y Dalitz plot. The good match between data and monte carlo gives us confidence on the proper reconstruction of the Dalitz variables in the $3\pi^0$ center of mass system. Fig. 6 shows the event density distribution *vs* R^2 for data and monte carlo, where $R^2 \equiv (Y^2 + X^2/3)$. The slope of the ratio of data *vs* monte carlo gives the direct measure of the quadratic decay parameter. The result is

$$b = (-0.6 \pm 1.4) \times 10^{-3}$$

with $\chi^2 = 36$ for 27 degrees of freedom, indicating a flat $3\pi^0$ Dalitz distribution.

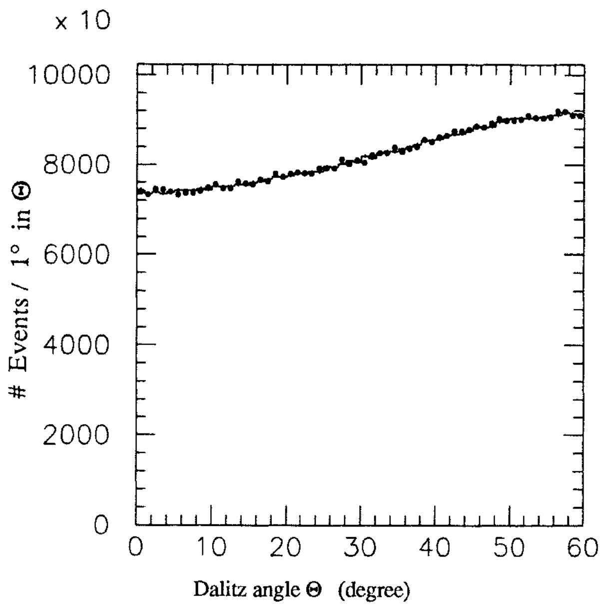

Fig. 5 Event density distribution in $K_L \to 3\pi^0$ Dalitz plot *vs* Dalitz angle Θ. Data with error bar were superimposed by a monte carlo with flat Dalitz distribution (dot).

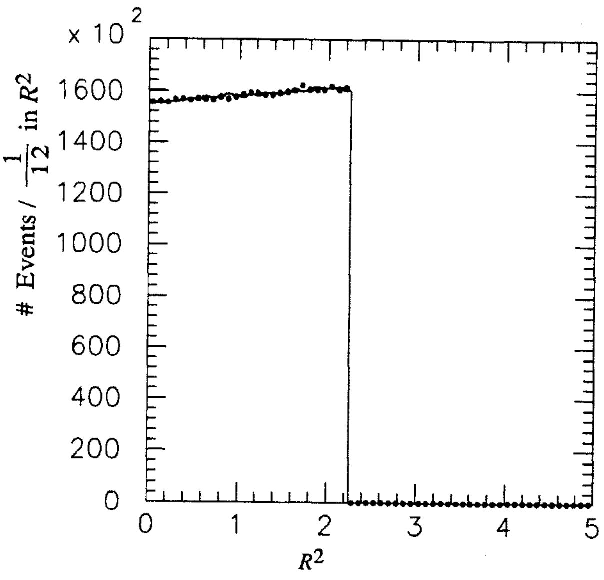

Fig. 6 Event density distribution for $K_L \to 3\pi^0$ *vs* R^2. Data were superimposed by a flat Dalitz monte carlo (dot).

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DISCUSSION

Q. **K. J. Peach** *(Univ. Edinburgh)*: In the comparison of data taken separately charged and neutral, how do you monitor the threshold and efficiency of the active regenerator?

A. **Y. B. Hsiung**: The efficiency of the active regenerator was monitored throughout the run (charged and neutral) with stray muon going through the regenerator veto-counters. The veto threshold was set to about half minimum ionized particle pulse height, so the inelastic interactions in the regenerator can be easily rejected. This threshold is far above the nuclear diffractive scattering and is only sensitive to the inelastic scattering which breaks up the nucleus apart into multi-particles. Therefore, it will not affect the remaining diffractive and inelastic background shape underneath the coherent peak. The background subtraction can be done separately in the charged mode and neutral mode, as long as the underline background shape remains the same between them, even through the background level may be different (say due to different veto inefficiency). Therefore, the background subtractions are quite insensitive to the variations of the threshold and inefficiency in the active regenerator throughout the run. This can be seen when comparing backgrounds between 20% and 100% data sample, they are virtually identical except the statistics.

Q. **K. Kleinknecht** *(Univ. Mainz)*: Concerning your table of systematic errors on ε'/ε for the old 20% data (where all 4 moves are taken simultaneously) and the ones expected for the 100% sample (taken in alternating runs): one would expect that at least one of the errors is much larger for the new 100% sample (e.g. the one from rate effects and the effect of the active regenerator). Why is that not taken into account in your table?

A. **Y. B. Hsiung**: Because of the use of double beam method, alternating regenerator between beams, and that we are always taking K_S and K_L decays ($\pi^+\pi^-$ or $\pi^o\pi^o$) at the same time, i.e. K_S and K_L decays are seen by the same detector simultaneously. Same accidentals would appear in both K_S and K_L decays, but different between charged and neutral mode, even though both charged and neutral modes are taken at the same time. Therefore, the systematic errors from the rate effect becomes a 2nd order effect. We can then evolute the accidental rate effect by overlaying the accidental events which were taken sparsely during the data run and were proportional to the proton intensity, to calculate the changes in yield ratio between K_S and K_L 2π decays. Hence, more statistics on accidental events in the entire data sample would help to reduce the uncertainty due to rate effect. Similar argument also applies to uncertainties in energy scale, resolution and acceptance, where other high statistic mode would also help to reduce the systematics. The effect of the background subtraction from the regenerators has been answered in question 1 (refer to Ken Peach's question), where more statistics would help to determine the background shape more accurately and hence reduce its uncertainty. Therefore in the experiment E731 (unlike NA31), we do not expect larger uncertainties for 100% data sample.

Q. **H. Nelson** *(CERN)*: The angular distribution of the γ of $\pi^+\pi^-\gamma$ in the $(\pi^+\pi^-)$ center of mass must be very different for direct emission and Inner Bremsstrahlung. Have you looked at this?

A. **Y. B. Hsiung**: Yes, we did look at the γ-angular distribution of $\pi^+\pi^-\gamma$ in the $\pi^+\pi^-$ center of mass for K_S (Inner Brem) and K_L (Inner Brem plus Direct Emission). We can see the differences between them, where there are more γ's emitted around 90% for $K_L \to \pi^+\pi^-\gamma$ than that for $K_S \to \pi^+\pi^-\gamma$. Our analysis on $\pi^+\pi^-\gamma$ is still in progress. I do not have plots for you to compare yet.

ε'/ε AND TOP QUARK

E.A. Paschos
Institut für Physik
Universität Dortmund
4600 Dortmund 50, Germany

and

Y.L. Wu
Institut für Physik
Johannes-Gutenberg-Universität
6500 Mainz, Germany

ABSTRACT

In this article we review recent theoretical developments on the CP-parameter ε'/ε in the standard model and its intimate connection with the mass of the top quark.

INTRODUCTION

During the last two years, there are two significant developments on studies of CP violation took place. Two experimental groups reported values /1/ for

$$\varepsilon'/\varepsilon = (3.3\pm1.1) \times 10^{-3} \quad \text{NA31 at CERN} \qquad (1)$$

and

$$\varepsilon'/\varepsilon = (-0.4\pm1.4(\text{stat})\pm0.6(\text{syst}) \times 10^{-3} \quad \text{E731 at FERMILAB .} \qquad (2)$$

On the theoretical side it becomes clear that the above ratio is closely correlated with the mass of the top quark. The ratio (ε'/ε) can be expressed as

$$\varepsilon'/\varepsilon = \frac{1}{\sqrt{2}} \frac{1}{|\varepsilon|} \left(\frac{\mathrm{Im}A_2}{|A_o|} - \frac{\mathrm{Im}A_o}{|A_o|}\right) . \qquad (3)$$

Even though the last term is suppressed by the factor $\omega = |A_2|/|A_o| = \frac{1}{22}$ it has been until recently the dominant term, with its contribution coming from the gluon-penguin diagram. However, when $m_t > m_w$, there are important terms from electroweak penguin diagrams where Z^o-bosons and photon are exchanged instead of gluons. Their importance is enhanced because

i) their contribution increases with m_t /2,3/,

ii) They contribute to $\mathrm{Im}A_2$ which is not suppressed in Eq.(3),

iii) the matrix elements of operators arising from electroweak penguin are not chirally suppressed.

The complete analysis of electroweak penguins was carried out by three groups /3,4,5/. The effects are calculated by deriving an effective Hamiltonian at the low energy scale p~μ

$$H_{eff}(\mu) = -\frac{G}{\sqrt{2}} \sum_{\substack{i=1\\ i\neq 4}}^{8} [\xi_c C_i^c(\mu) + \xi_t (C_i^t(\mu)] Q_i \quad (4)$$

with

C_i, C_j: Renormalized coefficient functions,

Q_i : Basis of operators, and

$\xi_q = V_{qd}V_{qs}$ with V_{ij} elements of the KM matrix.

The KM matrix elements contain phases which produce imaginary parts. The unitarity of the KM matrix implies

$$I_m\xi_c = -I_m\xi_t \approx -\beta\gamma\sin\delta'$$

and gives

$$I_m H_{eff}(\mu) = -(G/\sqrt{2}) I_m \xi_t \sum_i C_i(\mu) Q_i \qquad (5)$$

with $C_i(\mu) = C_i^t(\mu) - C_i^c(\mu)$. Substituting the effective Hamiltonian into the definition of ε', we obtain

$$\varepsilon'/\varepsilon = \frac{1}{\sqrt{2}} \frac{1}{|\varepsilon|} \frac{1}{|A_0|} \frac{G}{\sqrt{2}} \, \mathrm{Im}\xi_c \sum_{\substack{i=1 \\ i\neq 4}}^{8} (C_i(\mu) \times$$

$$[\langle Q_i\rangle_2 - \omega \langle Q_i\rangle_0]) \ . \qquad (6)$$

The problem for estimating the ratio is now divided into three parts:

1. determination of the coefficient functions $C_i(\mu)$ at a low energy scale μ,
2. estimates of the hadronic matrix elements $\langle\pi\pi|Q_i|K^0\rangle$, and
3. ranges for $I_m\xi_t$ to be determined from ε_K and the B^0-$\overline{B^0}$ mixing.

We shall discuss each of these problems in detail.

COEFFICIENT FUNCTIONS

The origin of the coefficient functions is two-fold: Some of them come from QCD and others from the electroweak theory. In addition they mix with each other when we integrate the renormalization group equations (RGE).

The standard method /6/ is to begin with initial conditions at $p\sim m_W$ and integrate the RGE from m_W to m_b then from $m_b\to m_c$... down to a scale μ which is relevant for the decay of K-mesons. Three groups reported values for the coefficients. Originally there were small differences of less than 10% which were mentioned in the articles. The differences have been understood in the meanwhile and I report on Table 1 the results of our calculation. The table is consistent with Flynn and Randall /7/ and Buchalla, Buras and Harlander /4/. A very small difference of a few percent, which comes from the two distinct methods of integration, is inessential for predictions of the ratio ε'/ε.

Table 1 Coefficients C_i at $\mu = 1\ GeV$ as functions of Λ and m_t.

Λ_4	m_t	C_1	C_2	C_3	C_5	C_6	C_7/α	C_8/α
0.1GeV	50	0.034	-0.035	-0.017	0.009	-0.051	-0.080	-0.005
	75	0.034	-0.036	-0.017	0.010	-0.054	-0.075	0.002
	100	0.032	-0.036	-0.017	0.010	-0.055	-0.058	0.016
	125	0.031	-0.036	-0.016	0.010	-0.056	-0.033	0.035
	150	0.029	-0.036	-0.016	0.010	-0.057	-0.001	0.058
	175	0.028	-0.036	-0.015	0.010	-0.057	0.036	0.085
	200	0.026	-0.036	-0.014	0.010	-0.058	0.079	0.116
	225	0.024	-0.035	-0.013	0.010	-0.058	0.126	0.151
	250	0.021	-0.035	-0.013	0.010	-0.059	0.178	0.188
0.2GeV	50	0.043	-0.044	-0.020	0.011	-0.070	-0.068	-0.003
	75	0.042	-0.044	-0.019	0.011	-0.074	-0.063	0.006
	100	0.041	-0.045	-0.019	0.011	-0.076	-0.046	0.026
	125	0.040	-0.045	-0.019	0.011	-0.078	-0.022	0.052
	150	0.038	-0.045	-0.018	0.011	-0.079	0.009	0.086
	175	0.036	-0.044	-0.017	0.012	-0.079	0.045	0.124
	200	0.035	-0.044	-0.017	0.012	-0.080	0.087	0.168
	225	0.032	-0.043	-0.016	0.012	-0.081	0.133	0.216
	250	0.030	-0.043	-0.015	0.012	-0.081	0.183	0.269
0.3GeV	50	0.051	-0.051	-0.021	0.012	-0.091	-0.057	0.000
	75	0.050	-0.052	-0.021	0.012	-0.096	-0.053	0.012
	100	0.049	-0.052	-0.020	0.012	-0.099	-0.037	0.037
	125	0.047	-0.052	-0.020	0.012	-0.101	-0.013	0.073
	150	0.046	-0.052	-0.019	0.012	-0.102	0.017	0.117
	175	0.044	-0.052	-0.019	0.012	-0.103	0.052	0.168
	200	0.042	-0.051	-0.018	0.012	-0.104	0.092	0.225
	225	0.040	-0.051	-0.017	0.013	-0.105	0.137	0.290
	250	0.038	-0.050	-0.017	0.013	-0.105	0.187	0.360

HADRONIC MATRIX ELEMENTS

At the low energy scale μ seven operators survive. Their matrix elements were computed in the chiral Lagrangian inspired by the 1/N expansion and factorizable amplitudes. The main interest here is on the elements $\langle Q_7\rangle_{0,2}$ and $\langle Q_8\rangle_{0,2}$. Keeping the O(1/N) terms we obtain

$$\begin{aligned}
\langle Q_6\rangle_0 &= -2Z(m_k^2-m_\pi^2)/\Lambda^2\chi \\
\langle Q_7\rangle_0 &= \frac{1}{\sqrt{2}} X + \frac{1}{N} Z \\
\langle Q_8\rangle_0 &= Z + \frac{1}{n}\frac{1}{\sqrt{2}} X \\
\langle Q_7\rangle_2 &= -X + \frac{1}{N}\frac{1}{\sqrt{2}} Z \\
\langle Q_8\rangle_2 &= \frac{1}{\sqrt{2}} Z - \frac{1}{N} X
\end{aligned} \qquad (7)$$

with

$$X = \sqrt{\frac{3}{2}}\, F_\pi (m_k^2 - m_\pi^2),$$

$$Z = 2\sqrt{\frac{3}{2}}\, (m_k^2/(m_s+m_d))^2\, F_\pi$$

and

$$\Lambda_\chi = 1.026 \text{ GeV} \ .$$

Our matrix elements $\langle Q_7\rangle_{0,2}$ and $\langle Q_8\rangle_{0,2}$ have more terms than those in ref./4/, because we consistently keep both terms of O(1) and O(1/N) in the expansion. Explicit substitution shows that for $\langle Q_8\rangle_{0,2}$ the O(1) terms are dominant and those proportional to 1/N give a subdominant contribution. The situation is reversed for the elements of

$\langle Q_7\rangle_{o,2}$ where the O(1) terms proportional to X are much smaller than the 1/N terms for N=3. Thus the expansion for this operator is questionable. Fortunately, we find its effect on $(\varepsilon'/\varepsilon)$ to be small for larger m_t, so that a large uncertainty can be tolerated. In our numerical calculation we introduce for Q_1 and Q_2 the next to leading corrections for the elements and their μ-dependence /8/ as given in Table 2.

Table 2

μ in GeV	$< Q_1 >_0$	$< Q_2 >_0$	$< Q_1 >_2$	$< Q_2 >_2$
0.6	-0.022	-0.045	0.013	0.013
0.7	-0.024	-0.050	0.011	0.011
0.8	-0.026	-0.053	0.010	0.010

The μ-dependence of other matrix elements is determined by the μ-dependence of the running strange quark mass /9/

$$m_s^2(\mu) = m_s^2(1\text{GeV}) \left(\frac{\alpha_s(\mu)}{\alpha_s(1\text{GeV})} \right)^{8/9} \qquad (8)$$

Combining the coefficients and matrix elements, we can study their relative contributions to the ratio $(\varepsilon'/\varepsilon)$. We introduce the following notation

$$h_{oct} = \sum_{i=1}^{2} C_i(\mu) \langle Q_i\rangle_0 ,$$
$$h_{27} = \frac{1}{\omega} \sum_{i=1}^{2} C_i(\mu) \langle Q_i\rangle_2 \qquad (9)$$
$$h_i = C_i(\mu) (\langle Q_i\rangle_0 - \frac{1}{\omega} \langle Q_i\rangle_2)$$
$$\text{for } i=3, \ldots, 8$$

and the Ω_i's are normalized to h_6 as follows

$$\Omega_{oct} = \frac{h_{oct}}{h_6} , \ldots , \Omega_p = \sum_{i=3}^{5} \frac{h_i}{h_6} , \qquad (10)$$
$$\Omega_{EWP} = \sum_{i=7}^{8} \frac{h_i}{h_6}$$

and

$$\varepsilon'/\varepsilon = (\varepsilon'/\varepsilon)_6 \, \Omega_{tot} \qquad (11)$$

With

$$(\varepsilon'/\varepsilon)_6 = \frac{1}{\sqrt{2}} \frac{1}{|\varepsilon|} \frac{1}{|A_0|} \text{Im}\xi_t h_6 > 1 \times 10^{-3}$$

as was established in earlier works /10/ Ω_{tot} is the sum of all Ω_i's including $\Omega_{\eta+\eta'}$. Our numerical values for the Ω_i's are plotted in Fig.1. The term $\Omega_{\eta+\eta'}$ is the long distance contribution from π-η-η' mixing /11/. The terms which show a strong dependence on m_t are Ω_{27} and Ω_{EWP}. The most important variation comes from Ω_{EWP} which includes effects of Q_7 and Q_8. We note that Ω_{EWP} starts positive, then at $m_t \sim O(m_W)$ is zero; it becomes negative for larger values of m_t and $m_t \sim 210$GeV produces a zero for Ω_{tot}, i.e. ε'/ε imitates the superweak theory /12/. These results are important to motivate us to study the variation of the results on the underlying parameters Λ_4, μ and m_s. The Λ_4 dependence of Ω_{tot} is very weak as shown in Fig.1. We find that Ω_{tot} has also a very weak dependence on μ and m_s. We also wish to point out that the above results are independent of the KM matrix elements which appear as a multiplicative factor to be discussed below. A further investigation for hadronic matrix elements is still necessary since there are non-factorizable terms and one should test the above expectations by lattice calculations.

K-M FACTOR $\text{Im}\,\xi_t$

A range for $\text{Im}\xi_t$

$$\text{Im}\xi_t \approx |V_{cb}| \, |V_{ub}| \sin\delta \qquad (12)$$

is obtained from fits of ε_k and B^0-$\overline{B^0}$. This analysis is well known and we do not include it here. It is presented in /13/. The results are consistent with /4/ and other careful analyses done by Kim, Rosner and Yua, /14/ and J. Maalampi and M. Roos /15/. The conclusions are that: by fitting ε_k, there are two solutions for δ one with δ <

$\pi/2$ and a second with $\delta > \pi/2$.

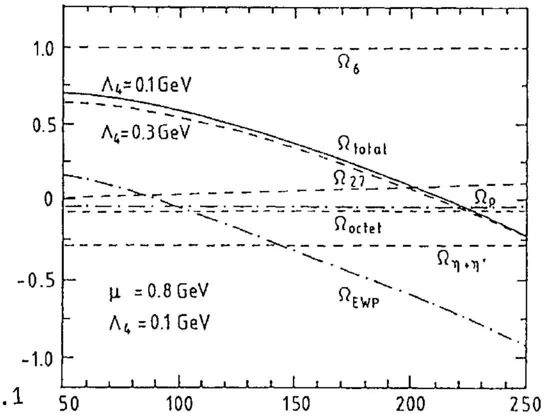

Fig.1

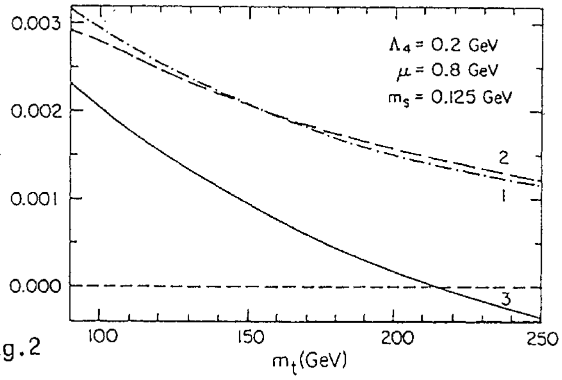

Fig.2

DISCUSSIONS AND CONCLUSIONS

Collecting the results from the above analysis, one can present results for ε'/ε. They are plotted in Fig.2 with central values from the parameters $B_k = 0.75$, $|V_{cb}| = 0.05$, $|V_{ub}| = 0.005$. The three curves correspond to

1. pure QCD case corresponding to $\alpha_{QED} = 0$, $\sim(\varepsilon'/\varepsilon)_6$,
2. the inclusion of $\Omega_{\eta+\eta'}$, and photon-penguin contributions,
3. Our full result after the Z^0-penguin and box diagrams are included, i.e. $(\varepsilon'/\varepsilon)_6 \Omega_{tot}$.

These three cases also show history of the development of $(\varepsilon'/\varepsilon)$ calculations during the past ten years. The cases 1) and 2) give very similar results which show that there is an approximate cancellation between $\Omega_{\eta+\eta'}$, and the photon-penguin contributions. If we use the tables and figures from Buchalla, Buras and Harlander /4/ we obtain identical curves.

Many features of ε'/ε are closely related with Ω_i and $I_m\xi_t$:

i) The superweak behaviour is determined by Ω_{tot} which has only a weak dependence on the underlying parameters μ, Λ_{QCD} and m_s. The dependence of Λ_{QCD} is shown in Fig.1.

ii) The variation of the ε'/ε on the underlying parameters μ, Λ_{QCD} and m_s can be described by $\Sigma = h_6\Omega_{tot}$. Studies of the μ-dependence noted that the variation is small /5/, especially for 170GeV < m_t < 220GeV. The dependence on Λ_{QCD} is somewhat bigger as shown in Fig.3. It varies by a factor 1.7. The m_s-dependence can be seen from Fig.2. It is very sensitive to m_s, which arises dominantly from the multiplicative factor $\langle Q_6\rangle_0$.

iii) Multiplying the curve 3 in Fig.2 by the allowed range of $\mathrm{Im}\xi_t$, we obtain the possible range of ε'/ε which is shown in Fig.4. We note that the curves show a substantial variation, but still predict a limited range for ε'/ε for each values of m_t. In paticular the ratio is measurable provided m_t < 180GeV. We notice, as expected, the zero value at $m_t \sim$ 210GeV and a band of values with an approximate 30% variation around the central values.

The three groups which analyzed CP phenomena for a heavy top agree that ε'/ε could become very small for larger $m_t > m_w$. The general conclusion is that it can imitate the superweak theory for the extreme case m_t > 200GeV.

From the various detailed analyses we can draw the following conclusions:

1. Significance of the electroweak terms is enhanced when the top quark is heavy /3,4,5/.
2. A complete renormalization analysis with the 1/N estimates for matrix elements gives the range $0.1 \times 10^{-3} < \varepsilon'/\varepsilon < 3 \times 10^{-3}$ for 200GeV > $m_t > m_w$.
3. There is still a large range predicted for ε'/ε, which indicates that it is impor-

tant to study and improve the theoretical estimates. We have tried to point out in this talk the underlying assumptions and show the variations they introduce.

4. It is crucial to improve the experiments and bring them in agreement with each other, because for the ranges allowed for the parameters.

i) The large value $\varepsilon'/\varepsilon \sim 1 \sim 2 \times 10^{-3}$ signifies an $m_t \sim 100$GeV.

ii) A smaller value $\varepsilon'/\varepsilon \sim 10^{-4}$ favours a heavier $m_t \sim 200$GeV or the superweak theory.

iii) Negative values occur for $m_t > 210$GeV.

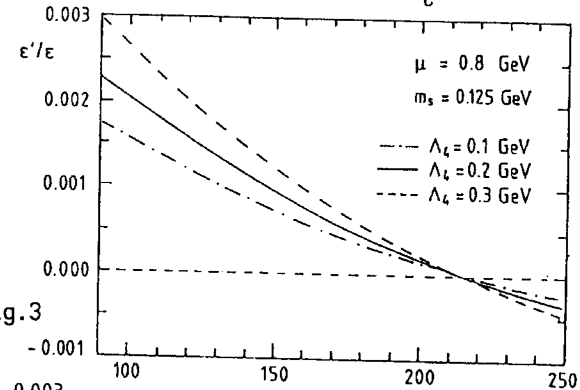

Fig.3

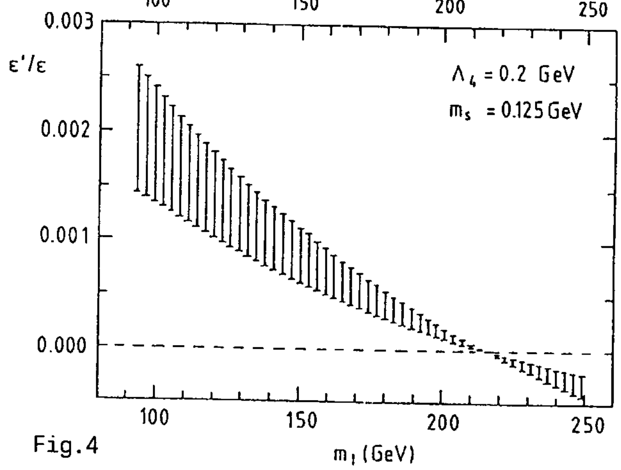

Fig.4

ACKNOWLEDGEMENT

We wish to thank T. Schneider and J.M. Schwartz for help with the calculations, and Dr. W. Bardeen for emphasizing the necessity for keeping lower order terms in the 1/N expansion (see Eqs.(7)).

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Q. **D. Chang** *(Northwestern Univ.)*: Is strong dependence on m_s a result of using the $\frac{1}{N}$ estimate of matrix elements?

A. **Y. L. Wu**: Yes.

Q. **E. Gabathuer**: If the top quark mass is determined from high energy colliders, can you use this to go back to determine more precisely the hadronic dynamics, e.g. matrix elements?

A. **Y. L. Wu**: Not yet. Since one also needs to know more precise KM matrix elements and ε'/ε. On the other hand, there are more than one matrix elements to be determined. One may use the data in $\Delta I = \frac{1}{2}$ rule to know information about the matrix elements.

EXPERIMENTAL STATUS AND PRELIMINARY RESULTS FROM THE CPLEAR COLLABORATION AT CERN

RAYMOND GAMET
Department of Physics, The University, P O Box 147,
Liverpool, L69 3BX, UK.

ABSTRACT

The experimental status and preliminary results from data taken in 1989 are presented from the CPLEAR collaboration at CERN. These show that it has been possible to identify a good sample of $K^o_s \to \pi^+\pi^-$ events, which yield a good lifetime measurement for the K^o_s.

INTRODUCTION

The experiment studies the reactions

$$\bar{p}p)_{rest} \to \begin{matrix} K^o K^- \pi^+ \\ \bar{K}^o K^+ \pi^- \end{matrix} \qquad (1)$$

The strangeness of the neutral kaon is tagged by observing the sign of the charged kaon.

The experiment will make precision tests of CP and CPT violation by studying the interference in the decays of K^o and $\bar{K}^o$ mesons to a given final state $|f\rangle$, by measuring the asymmetry

$$A_f(t) = \frac{R(\bar{K}^o \to f) - R(K^o \to f)}{R(\bar{K}^o \to f) + R(K^o \to f)} \qquad (2)$$

This will be done for as many final states as possible.

$$|f\rangle = |\pi\pi\rangle,\ |\pi\pi\pi\rangle,\ |\pi\ell\nu\rangle,\ |\gamma\gamma\rangle,\ \text{etc.}$$

For the 2π final state

$$\frac{\varepsilon'}{\varepsilon} = \frac{1}{3}\,(1 - |\eta_{00}|/|\eta_{+-}|) \qquad (3)$$

where

$$\eta_{2\pi} = R(K^o_L \to 2\pi)/R(K^o_s \to 2\pi) \qquad (4)$$

and is obtained from a fit to $A_{2\pi}(t)$.

To identify these channels it is clearly necessary to have good charged particle identification, accurate charged particle tracking and good photon energy and position resolution. Also, the desired "golden" channels (equation (1)) represent only ~ 4 x 10^{-3} of total $\bar{p}p$ annihilation rate so rapid online processing is necessary to check the quality of the data before recording them.

EXPERIMENTAL STATUS

A schematic diagram of the apparatus is shown in figure 1.

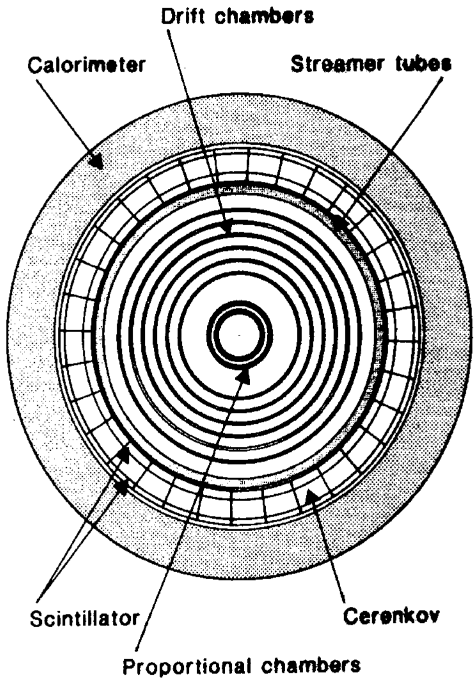

Fig 1

The detector is cylindrical with a radius of 1m and a length of 3m and is in a solenoidal magnetic field of 0.44T. The target is spherical of radius 7cm and contains hydrogen gas at 15 atmospheres pressure. All the apparatus is now working well, with the exception of the electromagnetic calorimeter. This is physically complete but has no readout electronics, so only charged track information is available.

The first three stages of online trigger processing for charged tracks are also working. The first of these is the "pretrigger". This demands that the number of charged tracks should be at least 2 and that at least one should be a kaon candidate, ie give no light in the Cerenkov but hits in both scintillators. This reduces the beam rate by about a factor of 5. The second stage is the "p_T cut". This uses the information from the drift chambers to place a rough cut on the minimum value of the momentum component in the transverse plane of any kaon candidate. This eliminates low momentum pions which have faked kaons in the pretrigger, and further reduces the event rate by a factor of about 6. The last stage which is presently working is the "track launcher". This associates wire chamber hits with each track and can be used to distinguish between "primary tracks", ie those with hits in the proportional chambers, and "secondary" tracks, ie those without such hits. Depending on the requirements placed on the numbers of primary and secondary tracks, this further reduces the event rate by factors of between 2 and 200.

PRELIMINARY DATA ANALYSIS

We shall here describe the analysis of about 5×10^7 events which were recorded during two running periods in November/December 1989. During the first stage of this data taking only the pretrigger was operational, and for the second stage the p_T cut was introduced. In the analysis of these data it was demanded that there should be exactly 4 charged tracks of total charge zero. It was required that just one of these should be a kaon, as identified by the Cerenkov and also by energy loss in the inner scintillator S1 and by "time-of-flight" to S1. It was then demanded that the kaon should form a good vertex with an opposite-sign track, whose time-of-flight should be consistent with a pion, and that the other two tracks should also form a good vertex. Cuts were then placed on the missing energy and momentum of the event to eliminate badly fitted tracks, misidentified tracks, and events with additional neutral particles. Some 4×10^4 events survived all these cuts, and in figure 2 is shown the invariant mass of the secondary pion pair.

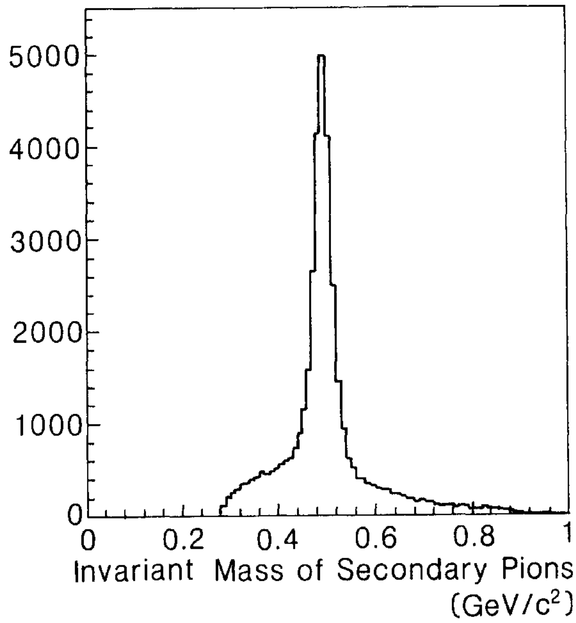

Fig 2

A clear peak is seen at the mass of the K^o on a small background. These events were then split into bins of K^o lifetime and, after estimated background subtraction, the number of $K^o \rightarrow \pi^+\pi^-$ events was obtained as a function of lifetime. This is shown in figure 3.

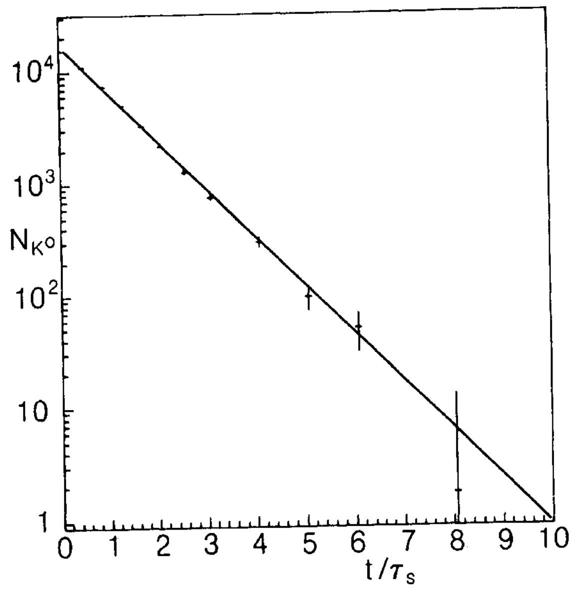

Fig 3

This shows a clear exponential behaviour and gives a measured lifetime of $(1.013 \pm 0.015)\tau_s$.

This event sample was then split into the two strangeness channels and the asymmetry $A_{+-}(t)$ (equation (2)) was calculated. This is shown in figure 4.

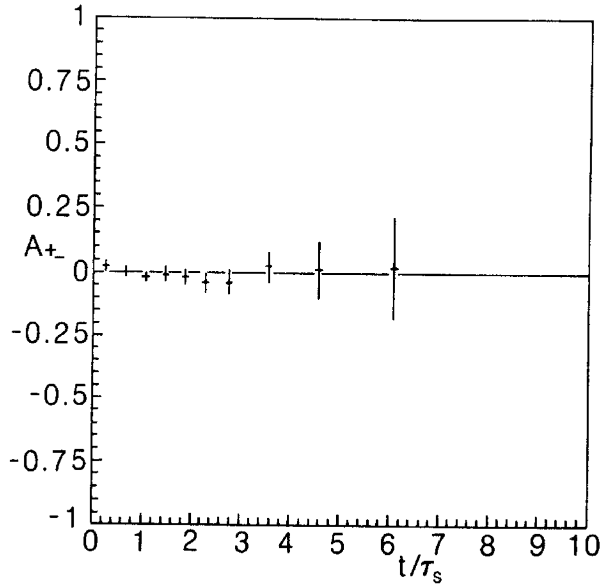

Fig 4

This is everywhere consistent with zero, but the asymmetry $A_{2\pi}(t)$ is not expected to become significantly different from zero until $t > 5\tau_s$, where the K^o_s and K^o_L interference becomes significant. This is just where the data run out of statistics. For the data taken in 1990 with the track launcher working, the trigger requirement has been made of 2 primary and 2 secondary tracks. This greatly enhances the proportion of decays beyond $t = 5\ \tau_s$ and should enable us to make a good measurement of $A_{+-}(t)$ from this year's data.

CONCLUSIONS

The apparatus, apart from the calorimeter, is working well, as are the first 3 stages of online trigger processing. We have shown we can isolate a good sample of $K^o_s \rightarrow \pi^+\pi^-$ events and measure a good lifetime. From the 1990 data we should be able to make a good measurement of $A_{+-}(t)$. In 1991 the calorimeter should be working and also further stages of trigger processing to analyse online the calorimeter data. We should then be able to start looking at other decay channels of the K^o and make a measurement of ε'/ε.

DISCUSSION

Q. **G. Gollin** *(Univ. Illinois)*: I realize the data are very preliminary. However, you seem to have a large background under your K° peak. Will this wash out the asymmetries you plan to measure? Will this cause you to confuse K° with $\bar{K}^\circ$?

A. **R. Gamet**: Almost all the background is in the bin 0–0.5 τs. There is almost no background beyond this. Hence this has no effect in the determination of the asymmetry.

Q. **K. Kleinknecht** *(Univ. Mainz)*: How many days of real data taking do these 15000 K° decays correspond to?

A. **R. Gamet**: Very few. We were also running with a preliminary trigger which meant that the beam rate had to be very low, being limited by the rate we could write events to tape.

Q. **I. Picek** *(Inst. Rudjer Bošković)*: What are the prospects to measure asymmetries for suppressed modes (e.g. $K^\circ(\bar{k}^\circ) \rightarrow 2\gamma$ or $K \rightarrow 3\pi$), because LEAR seems to be the best place to measure direct CP violation outside $K_L \rightarrow \pi\bar{\pi}$?

A. **R. Gamet**: We shall certainly look at CP violation in $K^\circ \rightarrow 3\pi$, and in as many other suppressed channels as possible, depending on the statistics we get in these channels.

RECENT DEVELOPMENTS IN NEUTRON ELECTRIC DIPOLE MOMENT AND RELATED CP VIOLATING QUANTITIES*

Darwin Chang
Department of Physics and Astronomy†
Northwestern University, Evanston, IL 60208
and
Fermi National Accelerator Laboratory,
P.O. Box 500, Batavia, IL 60510, USA

ABSTRACT

We summarize recent theoretical developments in CP violation related to the neutron electric dipole moment, chromo-electric dipole moments for quarks, chromo-electric dipole moment for gluon, and electric dipole moments for eletron and W boson.

Recent experimental improvements on measurements of the electric dipole moment of the neutron, D_N, and that of the electron, D_e, have inspired a lot of theoretical developments in CP violating models related to these quantities. Here I wish to summarize the developments since 1989. For reviews before 1989, see ref[1]. More recent reviews can be found in ref[2,3].

Experimentally, D_N was found to be $(-14\pm6)\times 10^{-26}$ e–cm by the Leningrad Group[4] and more recently $(-3\pm5)\times10^{-26}$ e–cm by the Grenoble Group[5]. Combining the two, we shall interprete this data as setting an upper bound on $|D_N| \leq 8\times10^{-26}$ e–cm. For D_e the most recently published data[6] gives $(-1.5\pm5.5\pm1.5)\times10^{-26}$. At this Conference, I was told that this limit has recently been improved by at least one order of magnitude. Therefore the upper bound on D_e is about 10^{-26}. As we shall see, due to the recent theoretical developments, these limits have imposed nontrivial constraints on various models of CP violation.

We shall start with a brief classification of various mechanisms of CP violation. In a gauge theory, the CP violation can be implanted typically into the following sources: (1)The left handed charged currents, like the standard Kobayashi-Maskawa (K-M) model[7]. The model predicts very small values for all the quantities that we are going to discuss in this talk. In a sense these quantities are good probes of the new physics beyond the standrad model. (2)The charged Higgs mixing, like the Weinberg-Branco model[8]. Since the CP violations in this type of model are typically suppressed by small Yukawa couplings, the charged Higgs does not have to be very heavy to suppress CP violation. (3)The neutral Higgs mixing between the scalar and pseudoscalar bosons, like the Lee model[9]. Typically, all models with extended Higgs sector contain a CP violating neutral Higgs subsector. (4)The right handed currents, like the left-right models[10]. Coexistence of left- and right-handed currents gives rise to a very interesting CP violating phase associated with the mixing between two currents. This CP violation does not need more than one generation of fermions. (5)Majorana mass, like the gluino mass or the neutrino masses. (6) Supersymmetric models[11] usually contain many sources of CP violation. Typically, it is a combination of colored, charged Higgs(the squarks) exchange and majorana masses of various neutralinos. (7)Strong CP θ parameter which we are not going to discuss here.

The recent attention on D_N was generated by Weinberg[12] who emphasized a new mechanism for D_N. He showed that there is an unique gauge invariant, P-odd, and T-odd operator of dimension 6, $\mathcal{O}_G$, involving solely the gluon field strength that can be written as

$$\mathcal{O}_G = -\frac{1}{3} f^{abc} g_{\alpha\beta} \tilde{G}^a_{\mu\nu} G^{b\mu\alpha} G^{c\nu\beta} \ . \qquad (1)$$

The operator can give potentially large contribution to D_N. Compared to another important operator, the color-electric dipole moment(CEDM) operator of quark q, $\mathcal{O}_q = \tilde{G}^a_{\mu\nu}\bar{q}\frac{1}{2}\sigma^{\mu\nu}T^a q$, the operator $\mathcal{O}_G$ should be identified as the color-electric dipole moment operator of the gluon[13,3].

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†Permanent address

Knowing the existence of $\mathcal{O}_G$, there are three issues to be investigated. The first one is to calculate its coefficient, $\mathcal{C}$, in a specific model of CP violation, and determine the scale at which the operator is induced. The second is to evolve, using QCD renormalization group (R.G.) equation, the operator to the low energy scale at which one can estimate its physical effect. The third is to evaluate its contribution to the D_N by calculating the corresponding hadronic matrix element. All three issues are coupled. Clearly, to include QCD corrections one needs to know the scale at which each operator is induced. The effect of the QCD R.G. correction is also very sensitive to the choice of the low energy hadronic scale. The proper hadronic scale to choose is presumably determined by the scale at which we think the hadronic matrix element can be evaluated with any confidence. The simplest way to estimate this is to use naive dimensional analysis[14]. It gives[12]

$$D_N \sim eM_X \zeta_{QCD}(\mu)(g_s(\mu)/4\pi)^{-3}\mathcal{C}(g_s(\mu)) \quad (2)$$

where M_X is the chiral symmetry breaking scale of $\sim 1.19GeV$, μ is a hadronic scale, and ζ_{QCD} is the QCD renormalization factor. However, the method was shown to be reliable[14] for the matrix elements which involve scales near the confinement scale ($\sim$ 250 MeV). Its reliability is not clear for scale near the neutron mass or M_X. In addition, near the confinement scale the QCD is known to be strong and the perturbative R.G. analysis is invalid there. Naive extrapolation of R.G. equation to lower energy gives a large and uncertain result. For this reason we shall take M_X as the low energy end of the R.G. evolution in order to make numerical comparison of different models. To get an idea of how uncertain the estimate of Eq(4) is, one can compare it with another recent estimate of this matrix element[15] which obtained a value 30 times smaller than Eq(4).

Weinberg [12] showed that $\mathcal{C}$ is induced, for models with CP violating mixing of the physical neutral Higgs bosons, through a two–loop diagram with top quark loop and neutral Higgs exchange. Dicus [16] showed that for models with CP violating charged Higgs mixing, similar two loop contributions can be obtained by replacing the neutral Higgs with the charged Higgs. In that case the fermion loop contains both top and bottom quarks. He noted that the contribution is not suppressed by the fact that b-quark is much lighter than the top quark. This is a little bit surprising because using the chirality argument one naively expects the calculation to be suppressed by two powers of lighter quark mass (m_b^2) with one from the Yukawa coupling and another one from the propagator. However this mass factor turns out to be cancelled by the fermion mass singularity in the loop integral. In ref.[17] we pointed out that Weinberg mechanism may also provide appreciable contribution to the D_N for models with CP violating left-right mixing. The diagrams can be obtained by replacing the charged Higgs in the previous model with the charged gauge bosons which contain the CP violating mixing. In this case, the m_b factor from the Yukawa coupling is replaced by the gauge coupling. As a result, left right models have the interesting feature that the lighter quark masses actually provide an enhancement factor due to the mass singularity in the loop integrals. In the supersymmetric model, the particles in the loops are the gluino, squarks, and quarks.[11]

The fermion mass singularity in the loop is of course a signal that the corresponding fermion loop should not be treated as a local operator above the the corresponding fermion mass scale. The proper way to treat this problem is to integrate out the t quark and the W boson first which induces a color electric dipole moment operator, as in eq(2), for the b quark. This operator is then QCD corrected using the R.G. technique down to the b-quark scale. The $\mathcal{O}_G$ operator is subsequently induced after the b quark is integrated out[19,18]. Therefore to calculate the QCD effect one needs the two by two anomalous dimension(a.d.) matrix of the operators $\mathcal{O}_G$ and $\mathcal{O}_q$.

A more useful basis of operators for the effective Hamiltonian are $\mathcal{O}_1(\mu) = g_s(\mu)^3\mathcal{O}_G(\mu)$ and $\mathcal{O}_2(\mu) = g_s(\mu)m_q(\mu)\mathcal{O}_q(\mu)$. The a.d. for $\mathcal{O}_1$ was calculate in ref[20]. Unfortunately, the sign of the answer was in error. Based on this wrong sign, Weinberg[12] concluded that any model with CP violation in the Higgs sector would require large finetuning to avoid the constraint from the D_N through this mechanism. The correct a.d.'s were later obtained[21] with $\gamma_{11} = \gamma_{GG} - 3\beta = -12C_A$ for $\mathcal{O}_1$ and $\gamma_{22} = \gamma_{qq} - \beta + \gamma_m = 4C_A - 16C_F$ for $\mathcal{O}_2$[22,21], where γ_{GG} is the a.d. for $\mathcal{O}_G$, γ_{qq} is the a.d. for $\mathcal{O}_q$ and $\gamma_m = -6C_F$ is the a.d. of the quark mass operator. For $SU(3)$, $C_A = 3$ and $C_F = \frac{4}{3}$. The operator $\mathcal{O}_G$ can induce the operator $\mathcal{O}_q$. This operator mixing is controlled by the a.d. $\gamma_{Gq} = 2C_A$. If one is only interested in $\mathcal{O}_G$ at low energy, this mixing is actually a higher order effect and can therefore be ignored[13].

It was later pointed out[15] that all the above a.d.'s have been calculated before by Morozov[23]. In fact he had calculated the a.d's of all the operators of dimension ≤ 8 except for the four fermion operators. Some of the other dim.6 or 8 operators can have effect as large as that of $\mathcal{O}_G$.[24]

The result $\gamma_{11} = -36$ indicates that the operator $\mathcal{O}_G$ is very suppressed by the QCD effect. This suppression is most severe in the neutral Higgs model of CP violation because $\mathcal{O}_G$ is induced at the highest scale. In spite of that, for resonable choice of CP violating parameter and Higgs masses, it can still result in D_N as large as the experimental limit.[2]

The operator $\mathcal{O}_q$ does not induce $\mathcal{O}_G$ through R.G. evolution. However, when one evolves below the threshold of the quark q, $\mathcal{O}_q$ induces an effective $\mathcal{O}_G$ through the matching condition at the threshold. In the charged Higgs models or the left right models, $\mathcal{O}_G$ is not induced until b quark is integrated out. Therefore, above the m_b scale, the QCD evolution affects only the operator $\mathcal{O}_q$ which has a much smaller suppression factor corresponding to $\gamma_{22} = \frac{-28}{3}$. In fact, for left right models, the m_q in $\mathcal{O}_2$ is m_t. Therefore setting $\gamma_m = 0$ we have $\gamma = \frac{-4}{3}$. Therefore, the QCD effect is least suppressive for the left right models.

If the CP-violation comes from the neutral Higgs boson mixing, the NEDM is estimated to be $D_N \sim 2.0 \times 10^{-21} \zeta^{NH}_{QCD}(\mu) h(m_t/M_H)(\text{Im } Z_2)$ e − cm. Here, Im Z_2 is the complex phase from the mixing between scalar and pseudoscalar Higgs bosons; the function $h(m_t/M_H)$ is defined in Refs.[12,16]. For $\mu \sim 1$ GeV, the QCD evolution factor, ζ^{NH}_{QCD}, is given by $\sim 3 \times 10^{-4}$, and D_N is about $6.0 \times 10^{-25} \text{Im} Z_2$ e − cm for $m_t \sim M_H$. For the charged Higgs boson case, $\zeta^{CH}_{QCD} \sim 10^{-3}$, D_N is about $3 \times 10^{-25} \text{Im } Z_2'$ $e - cm$ for $m_t \sim M_{H^+}$. Note that the case for susy models[11] is not very different from this case. In the left-right models, assuming the right-handed scale is around TeV, one has[17] $D_N \sim 1.59 \times 10^{-19} \zeta^{LR}_{QCD}(\mu) f(m_t/M_W) \sin\xi \sin\eta$ e − cm, where ξ and η are the left-right mixing angle and the CP-violation phase respectively. The function $f(x)$ is of order unity[18]. The QCD evolution factor is $\zeta^{LR}_{QCD}(1GeV) \sim 1.5 \times 10^{-3}$. We have $D_N \sim 2 \times 10^{-22} \sin\xi \sin\eta$ e − cm, assuming $m_t \sim M_W$. More numerical details can be found in ref.[3].

Another recent inspiring work is an observation by Barr and Zee[25]. In the neutral Higgs models of CP violation, the usual one loop mechanism for D_e is suppressed by three powers of electron mass and therefore negligible. Barr and Zee observed that this suppression factor is a one loop level accident which can be easily avoided by considering two loop contributions. They showed that the resulting D_e can be eight orders of magnitude larger than the traditional one loop mechanism. There are three follow-up calculations[26] of this mechanism with basically the same numerical conclusions.

Inspired by this observation, there are two recent works[27] which applied this mechanism to the CEDM of light quarks, D^c_q. It was found that this mechanism may be more important to D^c_q than to D_q because the diagram involves three powers of QCD coupling constant. Therefore it may also be the largest contribution to D_N in some neutral Higgs models of CP violation.

Finally we shall mention the recent developments about electric dipole moment of W, D_W. The most important physical consequence of D_W is again its contribution to D_N. There are two operators that can contribute to D_W. One of them is $SU(2)_L$ breaking and of dimension 4 and the other $SU(2)_L$ preserving and of dimension 8. Marciano and Queijeiro[28] tried to analyze this contribution model independently, they found that the upper bound for D_N of order 10^{-25} $e - cm$ could be used to place a very stringent upper bound on $D_W \leq 10^{-20}$ $e - cm$. However, they used only the dimension 4 term in their analysis for simplicity. In addition, in order to tame the divergence they have to postulate a form factor. It is interesting to ask whether such a form factor is realized in any model, or whether D_W can be larger than the above bound so that it may still be possible to measure directly in experiment.

Recent one loop analysis[29] of various models of CP violation showed that only models with right-handed current contribute and only the dimension 4 operator is induced in that case. However, it is clear that in models with two loop contributions[30], the dimension 6 operator will also be induced. It should be interesting to see how model independent are the form factors.

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DISCUSSION

Q. **X.-G. He** *(Univ. Melbourne)*: Comments: One should also take into account long range effects contribution to the neutron electric dipole moment because the Weinberg three gluon effect is not always dominant. (See our paper, He, McKellar & Pakvasa, Int. J. Mod. Phys. A, 1991.)

A. **D. Chang**: Yes, long range effects involving more than one quarks can be important. It is also more difficult to give precise predictions.

Q. **G. Gollin** *(Univ. Illinois)*: 1. The Barr-Zee work had $d_e = 4 \times 10^{-26}$ ecm. What modifications to the standard model are incorporated in their model?
2. It used to be the case that ε'/ε imposed strong constraints on extended Higgs models ($\sim$ 1983, as I recall). Is this no longer the case? Can you get large EDM's but small values of ε/ε now?

A. **D. Chang**: 1. Barr-Zee mechanism is particularly effective in the Higgs-exchange mechanism of CP violation in which the CP violation is due to charged or neutral Higgs exchange. More than one Higgs doublets are needed to implement this mechanism of CP violation.
2. The charged-Higgs-exchange model of CP violation is still close to being ruled out. It survives only because of the theoretical uncertainty involving hadronic matrix elements. Yes, it is possible to play with parameter to get larger EDM's but still small enough ε'/ε. For more constrained model like spontaneous Weinberg-Branco model, a new analysis is underway.

Q. **H. Nelson***(CERN)*: A group at LBL asked to report that the electron edm > $1.5.10^{-26}$ at 90% CL.
a) Is the upshot that experiments on edm's are now considered for more sensitive to non-standard models than before?
b) Is it possible to construct effectively superweak models that have large edms?

A. **D. Chang**: a) The experimental data on edm's of neutron and electron at the present level can already put nontrivial constraints on the non-standard models of CP violation.
b) It depends on what you mean by superweak model. If superweak model is defined, as usual, as a $\Delta s = 2$ interaction mediated by a heavy boson, then its contribution to the flavor neutral processes, like edm's of e or n, is usually very small. So far, no superweak model with large edm's exists.

PROBING CP VIOLATION WITH NEUTRAL B DECAYS

BORIS KAYSER
Division of Physics, National Science Foundation
Washington, DC 20550 USA

ABSTRACT

Neutral B decays can provide a decisive test of whether CP violation is due to the known weak interaction. If this interaction is indeed the origin of CP violation, then one expects large, cleanly-predicted CP-violating signals both in the decays to CP eigenstates, and in those to a variety of non-CP-eigenstates. Methods for extracting theoretically clean information from the non-CP-eigenstates are explained.

The experimental study of CP violation in the decays of neutral B mesons promises to be a clean test of whether CP violation is due to the known weak interaction. As is well-known, if this interaction is indeed the source of CP noninvariance, then this noninvariance arises from complex phase factors in the C(abibbo)-K(obayashi)-M(askawa) quark mixing matrix. When elements of this matrix occur at the vertices of some Feynman diagram, complex phase factors in these elements give the diagram as a whole an overall complex phase. *If* this diagram interferes with another one, this complex phase can make a difference and can lead to physical CP-violating effects. In short, if CP violation is due to the known weak interaction, then it always comes about through an interference between one amplitude and another.

In the decays of neutral B mesons, there is a particularly interesting interference which results from $B^o - \overline{B^o}$ mixing. Due to this mixing, a B born at time t = 0 as a pure $|B^o\rangle$ evolves in time t into a state $|B^o(t)\rangle$ which is a coherent superposition of $|B^o\rangle$ and $|\overline{B^o}\rangle$, and is given by

$$|B^o(t)\rangle = \exp[-i(M - i\tfrac{\Gamma}{2})t] \times \qquad (1)$$
$$\times \{\cos(\tfrac{\Delta m}{2}t)\,|B^o\rangle + i\eta_M \sin(\tfrac{\Delta m}{2}t)\,|\overline{B^o}\rangle\}.$$

Here, M, Γ, and Δm are, respectively, the average mass, common width, and mass difference of the neutral B mass eigenstates, and $\eta_M \equiv \exp(-2i\delta^M_{CKM})$ is a phase factor in which the $B^o - \overline{B^o}$ mixing phase, δ^M_{CKM}, depends only on CKM matrix elements. Now, if f is some final state into which both a pure $|B^o\rangle$ and a pure $|\overline{B^o}\rangle$ can decay, then as Eq.(1) makes clear, both the amplitude $w_f \equiv \langle f|T|B^o\rangle$ and the amplitude $\overline{w_f} \equiv \langle f|T|\overline{B^o}\rangle$ will contribute to the decay of the time-evolved particle $|\overset{(-)}{B^o}(t)\rangle$. The interference between these two contributing amplitudes can lead to CP violation.

If this interference is to be large, then w_f and $\overline{w_f}$ obviously must be of comparable size. Now, it was noted long ago [1] that if f is a CP eigenstate with CP parity η_f *and* $\langle f|T|B^o\rangle$ is dominated by a single weak decay mechanism with one CKM phase δ^f_{CKM}, then

$$\frac{\overline{w_f}}{w_f} = \eta_f \exp(-2i\delta^f_{CKM}). \qquad (2)$$

Thus, w_f and $\overline{w_f}$ are of *equal* size, and the interference is maximized.

From Eq.(1) and its analogue for the time-evolved state $|\overline{B^o}(t)\rangle$ which comes from an initially pure $|\overline{B^o}\rangle$, it follows that the rates for decay of $B^o(t)$ and $\overline{B^o}(t)$ into a CP eigenstate f are given (when one CKM phase dominates) by

$$\Gamma[\overset{(-)}{B^o}(t) \to f] \propto |w_f|^2 e^{-\Gamma t} \times$$
$$\times \{1 \overset{(-)}{+} \alpha_f \sin(\Delta m\, t)\}. \qquad (3)$$

Here, $\alpha_f \equiv \eta_f \sin[2(\delta^M_{CKM} + \delta^f_{CKM})]$. The sin $(\Delta m\ t)$ term arises from the interference between the w_f and $\overline{w_f}$ contributions to the decay. Since this term contributes oppositely to the $B^o(t)$ and $\overline{B^o}(t)$ decay rates, it clearly violates CP invariance. Note that since both δ^M_{CKM} and δ^f_{CKM} depend only on the CKM matrix, the coefficient of the CP-violating term, α_f, can be cleanly predicted from a knowledge of this matrix alone.

While the CP-violating effect in the decay rates (3) can be large, studying it experimentally will nevertheless require 10^8 to 10^9 B mesons, simply because the branching ratio for B decay to a typical CP eigenstate, like that for B decay to most states, is small. Thus, the challenge will be to accumulate the necessary statistics. In view of this fact, it is obviously of interest to ask in what ways one can relax the requirement that the final state be a pure CP eigenstate and still have a useful decay mode. In this way, we might find useful modes with bigger branching ratios. We might also find modes in which the CP-violating effects are predicted to be the same as in related CP-eigenstate decays. Combining the data on a non-CP-eigenstate decay with those on a CP-eigenstate decay which measures the same CP-violating parameter will permit one to determine this parameter with greater statistical accuracy from a given total number of B mesons.

Let us, then, consider classes of final states which are not CP eigenstates, but are close enough to being such states to contain large, cleanly-predicted CP-violating signals. The first class consists of two-body states such as $\psi\rho^o$ or $D^{*+} D^{*-}$, in which the particles are either their own antiparticles or else the antiparticles of each other, but in which both of them have spin. Here, the final state is a CP-self-conjugate collection of particles, but because of the spin it is not a CP eigenstate. To see this, consider the example where both of the final particles have spin S = 1. Since the B is spinless, its two-body decays always yield particles of equal helicity. Since CP reverses helicities, in our S = 1 example the CP-eigenstate final states are then $|0,0\rangle$, $|1,1\rangle + |-1,-1\rangle$, and $|1,1\rangle - |-1,-1\rangle$, where the numbers indicate the helicities. Clearly, the last of these three possible states has CP parity opposite to that of the other two. Thus, if the final state is a mixture of all three of these states, then it is not a CP eigenstate.

Nevertheless, a pure CP-eigenstate component of the final state can be isolated through angular distribution measurements [2]. To see how, consider the sequence of decays

$$\overset{(-)}{B^o}(t) \to \psi(\lambda) + K_+^*(\lambda) \quad \hookrightarrow K_S + \pi^o \quad , \qquad (4)$$

where λ denotes helicity. If we keep only those events in which the K* decays to $K_S\pi^o$, then, since the latter is a (CP-even) CP eigenstate and CP is conserved in the strong K* decay, the only component of the K* that can contribute is the CP-even component $K_+^* \equiv K^{*o} + \overline{K^{*o}}$. Thus, the particles produced in the primary decay of the sequence (4) do form a CP-self-conjugate set. However, the ψK_+^* state is not a CP eigenstate because of the ψ and K* spins.

Imagine, now, that one measures the distribution in the angle θ_K between the K_S momentum in the K_+^* rest frame, and the K_+^* momentum in the B rest frame. If the data are summed over the momenta and spins of the decay products of the ψ, the θ_K distribution will be given by

$$\frac{d\Gamma[\overset{(-)}{B^o}(t) \to \psi K_+^* \to \psi K_S\pi^o]}{d\cos\theta_K} = \qquad (5)$$

$$= \sum_\lambda \Gamma[\overset{(-)}{B^o}(t) \to \psi(\lambda)K_+^*(\lambda)] \, \frac{d\Gamma[K_+^*(\lambda)]}{d\cos\theta_K} \quad .$$

Since $d\Gamma[K_+^*(0)] \,/\, d\cos\theta_K = \cos^2\theta_K$, while $d\Gamma[K_+^*(\pm1)]/d\cos\theta_K = (\sin^2\theta_K)/2$, this distribution has the form $\overset{(-)}{a}(t)\cos^2\theta_K + \overset{(-)}{b}(t)\sin^2\theta_K$, where the coefficients a(t), b(t) correspond to a parent $B^o(t)$, and $\bar{a}(t)$, $\bar{b}(t)$ to a parent $\overline{B^o}(t)$. Note that in this distribution the $\cos^2\theta_K$ term comes entirely from the ψK_+^* state with $\lambda = 0$. Since this is a CP eigenstate, Eq.(3) applies, and we have

$$\overset{(-)}{a}(t) = \Gamma[\overset{(-)}{B^o}(t) \to \psi(0)K_+^*(0)] \propto$$

$$\propto e^{-\Gamma t}\{1 \overset{(-)}{+} \alpha_{\psi K^*}\sin(\Delta m\ t)\}, \qquad (6)$$

with $\alpha_{\psi K^*} \equiv \eta_{\psi(0)K_+^*(0)}\sin[2(\delta^M_{CKM} + \delta^{\psi K_+^*}_{CKM})]$ a cleanly-predicted CP-violating coefficient. Indeed, since ψK_+^* has the same quark content as ψK_S, $\delta^{\psi K_+^*}_{CKM} = \delta^{\psi K_S}_{CKM}$. Since $\eta_{\psi(0)K_+^*(0)} = -\eta_{\psi K_S}$, we then have $\alpha_{\psi K^*} = -\alpha_{\psi K_S}$. Thus, isolation of the $\cos^2\theta_K$ term in the θ_K distribution selects the decays into a CP eigenstate where the CP-violating coefficient is just the opposite of the one encountered in $B \to \psi K_S$. Hence, one may improve the statistical error on this coefficient by combining the ψK_+^* and ψK_S data.

The error on the coefficient $\alpha_{\psi K_S}$ as determined from the ψK_+^*data alone can be reduced by not using merely the distribution in θ_K but the joint distribution in all of the (three) angles which describe the decay sequence $B \to \psi + K_+^* \to (e^-e^+) + (K\pi^o)$. (Here we are assuming that the ψ will in practice be detected through its decay to a lepton pair.) It has been found [3] that by analyzing this joint distribution, one can determine $\alpha_{\psi K_S}$ from ψK_+^* data with an error just one-to-two times that associated with determining this same quantity from the simpler decay $B \to \psi K_S$.

The precise usefulness of the decay mode $B \to \psi K_+^*$ depends, of course, on the actual CP-eigenstate or helicity composition of the final state. If, for example, $B \to \psi K^*_+$ always yields the same one of the three possible final CP eigenstates, then this decay has all the advantages of any decay to a CP eigenstate. Let us, then, ask what helicity composition we expect. The quark-level process which dominates $B^o \to \psi K^*_+$ is the $\bar{b}$ decay $\bar{b} \to \bar{c} + c + \bar{s}$. The $\bar{c}$ and c then combine to make the ψ. To the extent that the $\bar{c}$ and c are relativistic, the $\bar{c}$ has positive helicity, and the c negative helicity, due to the left-handed character of the weak interaction. To make a ψ, the $\bar{c}$ and c must go off in the same direction, so the contributions of their helicities to the helicity of the ψ cancel. Since the ψ is a $\bar{c}c$ bound state with L = 0, there is no further contribution to its helicity. Thus, to the extent that this simplified picture is correct, the ψ produced in $B \to \psi K^*_+$ has helicity zero [2]. The final state is then $|\psi(0)K^*_+(0)\rangle$, which is a CP eigenstate, so that $B \to \psi K^*_+$ is a very useful decay mode.

To examine the ψ helicity structure experimentally, the ARGUS collaboration has studied the inclusive decays $B \to \psi + X$. Focussing attention on decays yielding a ψ of high momentum, in which the undetected system X will usually be a K* or K, the collaboration finds [4] that, to within errors, the ψ produced in $B \to \psi K^*$ does indeed have helicity zero 100% of the time! Thus, $B \to \psi K^*_+$ does yield a CP eigenstate. If this turns out to be an accurate statement (the errors are still significant), then the use of angular-distribution information will not be important in studies of $\overset{(-)}{B}{}^o(t) \to \psi K^*_+$. A measurement of the time-dependent rate will be enough [5]. Of course, in other decay modes where nature proves less generous, the analysis of angular distributions will still be a powerful tool.

A second class of final states which are not CP eigenstates but promise to be useful consists of three-body states such as $\psi K_S \pi^o$, where the $K_S\pi^o$ pair need not have come from a K*. Although ψ, K_S, and π^o are all their own antiparticles, the $\psi K_S \pi^o$ final state is not a CP eigenstate because of the various angular momentum possibilities. However, it can be shown [6] that when the projection of the ψ spin along the normal to the $B \to \psi K_S \pi^o$ decay plane has a definite value S_n, referred to as the transversity, the final state is a CP eigenstate with CP parity $-(-1)^{S_n}$. The B decays to the CP-even states with $S_n = \pm 1$ can be separated from those to the CP-odd one with $S_n = 0$ by using the angular distribution of the decay of the ψ [6,7], in much the same way as the B decays to the ψK^*_+ state with helicities zero can be isolated by using the angular distribution of the decay of the K* (cf. Eq. (5)). Once $\Gamma[\overset{(-)}{B}{}^o(t) \to \psi(S_n{=}\pm1)K_S\pi^o]$ has been separated from $\Gamma[\overset{(-)}{B}{}^o(t) \to \psi(S_n{=}0)K_S\pi^o]$, one may write either of these rates in the form of Eq. (3), with the appropriate η_f, and extract the CKM phase.

The transversity approach has now been applied to a variety of three-body final states [8].

The final class of promising final states I would like to discuss [9] consists of two-body states $a_1\overline{a_2}$ in which a_1 and a_2 are mesons which differ in J^{PC} or mass, but have the same quark content, $q_x\overline{q_y}$. Thus, the state $a_1\overline{a_2} = (q_x\overline{q_y})(\overline{q_x}q_y)$ contains a CP-self-conjugate collection of quarks. Examples of such a state are $\rho^+\pi^-$ and $D^{*+}D^-$. As we shall see, decays to such a state can involve a large, cleanly-predicted CP-violating signal, so that the state can be quite useful if it has a large branching ratio. Furthermore, the branching ratio for $\rho^+\pi^-$, for example, is indeed expected to be larger than that for $\pi^+\pi^-$. These two final states probe the same CP-violating CKM phase angle, but while $\pi^+\pi^-$ is a CP eigenstate and has been the most commonly cited probe of this angle, $BR(B \to \pi^+\pi^-)$ is predicted to be only $\sim 2\times10^{-5}$ [10].

As has been said, if the interference leading to CP violation in $B^o(t) \to f$ is to be large, then the amplitudes w_f and $\overline{w_f}$ for pure B^o and $\overline{B^o}$ decay into f must be of comparable size. Now, when f is a final state of the type $a_1\overline{a_2} = (q_x\overline{q_y})(\overline{q_x}q_y)$, so that it is a CP-self-conjugate collection of quarks, then each diagram for $B^o \to f$ has a simply-related analogue in $\overline{B^o} \to f$. This analogue is just the CP-conjugate of the $B^o \to f$ diagram, except for the way in which the outgoing quarks join together to make hadrons. For example, in the decay of B^o (a $\bar{b}d$ bound state) to $\rho^+\pi^-$, the dominant diagram is a quark decay diagram in which $\bar{b} \to \bar{u} + u + \bar{d}$. The u and $\bar{d}$ then bind to make the ρ^+, while the $\bar{u}$ joins the spectator d quark to make the π^-. The analogue of this diagram in $\overline{B^o} \to \rho^+\pi^-$ is one in which $b \to u + \bar{u} + d$. The $\bar{u}$ and d then bind to make the π^-, while the u joins the spectator $\bar{d}$ quark to make the ρ^+. Since any $B^o \to f$ diagram and its $\overline{B^o} \to f$ counterpart differ only in the last step where the outgoing quarks join up to make hadrons, we expect the two diagrams to be of comparable size. In particular, if some diagram dominates $B^o \to f$, then its $\overline{B^o} \to f$ analogue should be of comparable size and dominate $\overline{B^o} \to f$. Then w_f and $\overline{w}_f$ will be of comparable size, as desired for maximal CP-violating interference between them. Furthermore, since the weak interaction part of the $B^o \to f$ diagram and that of its $\overline{B^o} \to f$ counterpart are just CP conjugates of each other, the CKM phases of the two diagrams are just equal and opposite.

Let us suppose that one diagram, with magnitude A, CKM phase δ^f_{CKM}, and strong interaction phase α, does dominate $B^o \to f$. Then its analogue, with comparable magnitude $\bar{A}$, CKM phase $-\delta^f_{CKM}$, and some other strong interaction phase $\bar{\alpha}$, will dominate $\overline{B^o} \to f$. The amplitudes for the decays $\overset{(-)}{B^o} \to \bar{f}$, where $\bar{f}$ is the final state CP conjugate to f (e.g., if $f = \rho^+\pi^-$, $\bar{f} = \rho^-\pi^+$), then follow from those for $\overset{(-)}{B^o} \to f$ by applying CP. We have in all

$$
\begin{aligned}
<f|T|B^o> &= A \exp(i\,\delta^f_{CKM}) \exp(i\alpha) \\
<f|T|\overline{B^o}> &= \bar{A} \exp(-i\,\delta^f_{CKM}) \exp(i\bar{\alpha}) \\
<\bar{f}|T|\overline{B^o}> &= A \exp(-i\,\delta^f_{CKM}) \exp(i\alpha) \\
<\bar{f}|T|B^o> &= \bar{A} \exp(i\,\delta^f_{CKM}) \exp(i\bar{\alpha})\ .
\end{aligned} \quad (7)
$$

Inserting these expressions into Eq. (1) and its analogue for the time-evolved state $|\overline{B^o}(t)\rangle$, we find that the time-dependent decay rates $\overset{(-)}{\Gamma}_f(t) \equiv |\langle f|T|\overset{(-)}{B^o}(t)\rangle|^2$ and $\overset{(-)}{\Gamma}_{\bar{f}}(t) \equiv |\langle \bar{f}|T|\overset{(-)}{B^o}(t)\rangle|^2$ for decay into f and $\bar{f}$ are given by

$$
\begin{aligned}
\Gamma_f &= e^{(-\Gamma t)}\{c^2A^2 + s^2\bar{A}^2 + 2csA\bar{A}\sin(\delta_W + \alpha_S)\} \\
\bar{\Gamma}_f &= e^{(-\Gamma t)}\{c^2\bar{A}^2 + s^2A^2 - 2csA\bar{A}\sin(\delta_W + \alpha_S)\} \\
\Gamma_{\bar{f}} &= e^{(-\Gamma t)}\{c^2\bar{A}^2 + s^2A^2 + 2csA\bar{A}\sin(\delta_W - \alpha_S)\} \\
\bar{\Gamma}_{\bar{f}} &= e^{(-\Gamma t)}\{c^2A^2 + s^2\bar{A}^2 - 2csA\bar{A}\sin(\delta_W - \alpha_S)\}
\end{aligned} \quad (8)
$$

Here, $c \equiv \cos(\frac{\Delta m}{2}t)$, $s \equiv \sin(\frac{\Delta m}{2}t)$, the weak CKM phase δ_W is $2(\delta^M_{CKM} + \delta^f_{CKM})$, and the strong phase α_S is $\alpha - \bar{\alpha}$. Note that, apart from their overall normalizations, the four decay rates in Eqs. (8) depend on three unknowns: $\bar{A}/A$, $s_+ \equiv \sin(\delta_W + \alpha_S)$, and $s_- \equiv \sin(\delta_W - \alpha_S)$. By measuring the four decay rates, one can determine these three unknowns. From s_+ and s_-, one can then determine $\sin\delta_W$, the cleanly-predicted CKM quantity of interest, by using the relation

$$
\sin^2\delta_W = \frac{1}{2}[1 + s_+s_- \pm \sqrt{(1-s_+^2)(1-s_-^2)}]\ . \quad (9)
$$

To be sure, this relation only determines $\sin\delta_W$ up to a four-fold ambiguity, but one can still see whether any of the four candidate values agrees with the theoretical prediction. Alternatively, one can resolve the ambiguity by studying a related decay to a CP eigenstate, which may yield a value for $\sin\delta_W$ which is less precise, due to limited statistics, but is unambiguous. Note that while the four decay rates of Eqs. (8), unlike the rates for decay into a CP eigenstate, do depend on an unknown strong phase α_S, the combined use of all four of these rates enables one to test the theoretical prediction for $\sin\delta_W$ without knowing α_S.

In summary, CP violation in at least some neutral B decays is theoretically very clean in the sense that it depends only on CKM matrix elements. To study this clean CP violation will require a lot of B mesons, but it is now clear that cleanly-predicted CP-violating signals can be extracted not only from neutral B decays to CP eigenstates, but from their decays to non-CP-eigenstates as well. By using both the CP-eigenstate and non-CP-eigenstate decays, one will be able to get more accurate results from a given number of B mesons.

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DISCUSSION

Q. **Kam-Biu Luk** *(UC, Berkeley/LBL)*: What is the size of the CP violating effect in those $B^\circ \rightarrow$ CP-odd final state decays?

A. **B. Kayser**: In the standard model, the CP violating effect is at least 8%.

CP-VIOLATING EFFECTS IN THE LEPTON SECTOR

Werner Bernreuther

Institut für Theoretische Physik
Universität Heidelberg
D-6900 Heidelberg, FRG

ABSTRACT

Theoretical predictions for the electric dipole moment d_e of the electron are briefly reviewed in view of a new improved limit on d_e. Moreover, the CP-odd weak dipole form factor of the τ lepton and observables which are sensitive to it are discussed.

INTRODUCTION

If the Kobayashi-Maskawa (KM) mechanism [1] is the only source of CP nonconservation in Nature then CP-violating and/or T-violating effects among leptons are bound to be unobservably small. However, sizeable CP-violating interactions among leptons are conceivable since they occur quite naturally in extensions of the standard model (SM) (cf., e.g., [2] for a review). This, of course, is the motivation for experimental searches: CP-odd or T-odd leptonic observables, respectively, are "sensors" for new CP-violating interactions. By far the best observable for this purpose is the electric dipole moment (EDM) of the electron, because it can be determined very precisely from experiments with heavy atoms. T-odd quantities involving the muon can be measured much less precisely (for a compilation, see [3]). As to the τ lepton: So far almost nothing is known about T-odd quantities involving the tau. Since one gets plenty of τ's at the Z peak at LEP one can do, however, some CP tests and measure CP-odd correlations in $Z \to \tau^+\tau^-$ in the near future.

ELECTRON EDM

The electron EDM d_e is defined as the electromagnetic form factor of the electron which is P- and T-odd in the static limit. A nonzero d_e induces a CP-violating effective EDM interaction in an electromagnetic field

$$H_{eff} = \frac{i}{2} d_e \bar{e}(x)\sigma_{\mu\nu}\gamma_5 e(x)F^{\mu\nu}(x) \qquad (1)$$

which reduces in the non-relativistic limit to $H_I = -d_e\vec{\sigma}\cdot\vec{E}$. This interaction, which induces a linear Stark effect, is the basis of experimental searches for d_e. The best experimental limits on d_e were obtained from the search for permanent atomic EDMs, in particular of cesium and of thallium. The cesium [4] and thallium [5] experiments make use of the following results from atomic theory: If the electron has a nonzero d_e, it contributes to the EDM of Cs and of $T\ell$ in an enhanced way. Schematically one has

$$d_A = R_A d_e \quad + \quad \text{nuclear contributions.} \qquad (2)$$

Atomic physics calculations, using various approaches yield $R_{Cs} \approx 120$ [6-9] and $R_{T\ell} \approx -600$ [7,10-12] with an estimated error on these enhancement factors of order 10 %.

Obviously, if a nonzero d_A for some atom should be found, elaborate theoretical input would be necessary to pin down its origin. So far only d_A's consistent with zero have been measured. It is customary to deduce from these measurements upper bounds on the electron EDM and on the nuclear contributions barring accidental cancellations between the different contributions in (2). In this way the Amherst group [4] obtained $d_e = (-1.5 \pm 5.5 \pm 1.5) \times 10^{-26} ecm$. The Berkeley group using $T\ell$ recently reported [5]

$$d_e = (-2.7 \pm 8.3) \times 10^{-27} ecm \tag{3}$$

which yields a limit on $|d_e|$ which is almost an order of magnitude smaller than the previous limit. It is the most precise laboratory determination of a particle EDM so far. In order to appreciate this number one should note the following. Eq. (4) implies that $|d_e|$, if nonzero at all, must be smaller than about 10^{-10} times the length scale set by the inverse Z boson mass. $M_Z^{-1} \simeq 2 \times 10^{-16} cm$ is roughly the scale up to which the electron is known to be pointlike. Below that scale the electron might have an asymmetric charge contribution, i.e., a nonzero d_e. At this level of sensitivity d_e provides — as the experimental limits on the neutron EDM — an important constraint on non-standard models of CP violation.

Finally, an experiment using $T\ell F$ [13] obtained $d_e = (-1.4 \pm 2.4) \times 10^{-25} ecm$. This group also published in [13] new limits on P- and T-odd (semi-)hadronic interactions and on d_{proton}.

MODELS OF LEPTONIC CP VIOLATION

Which CP-violating forces can generate an observable $d_e \neq 0$? First two rather general statements. i) In renormalizable theories of CP violation d_e must be induced by loop diagrams as the interaction (2) is nonrenormalizable. ii) Since the interaction (2) flips the electron chirality, $d_e \neq 0$ requires chirality-flipping terms (notably fermion mass terms) in the Lagrangian of the theory.

As to models: Generation mixing among leptons has not been observed yet. Therefore CP violation à la KM in the lepton sector of the 3 generation SM is not expected to yield observable effects. In fact this mechanism — if it exists at all — leads only to unobservably small leptonic EDMs. In the SM with mass-degenerate neutrinos there are no elementary CP-violating couplings. Nevertheless CP-violation may be transferred from the quark sector to the lepton sector via W boson exchange. The effects are expected to be tiny as they necessarily involve higher loop orders. It has been argued [14] that this leads to a nonzero d_e at 3-loop order and a calculation yields [14]

$$|d_e|_{SM} \lesssim 1.4 \times 10^{-37} ecm \tag{4}$$

using that the mass of the t quark $m_t < 200\ GeV$.

In many extensions of the SM CP-violating forces among leptons occur quite naturally (in particular, there can be CP-violating interactions which have nothing to do with generation mixing) and generate a nonzero d_e at 1-loop order through the amplitudes depicted in Fig. 1. The boson B must couple both to e_L and e_R with complex couplings g_L and g_R, respectively, such that $Im(g_L g_R^*) \neq 0$. Moreover, F must provide the chirality flip. One may distinguish between two types of models:

Flavor-changing models

In these models d_e is generated through the amplitudes of Fig. 1 where F is a fermion from a higher generation. Examples include leptoquark models, where B is a leptoquark boson and F is a quark

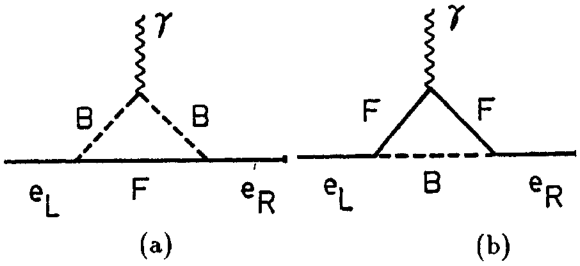

Fig. 1: Generic 1-loop diagrams which can generate a nonzero d_e. $F(B)$ denotes a fermion and a boson of spin zero or one, respectively.

and flavor changing Higgs model where B is a neutral Higgs and F is the τ lepton. These models can generate $|d_e| \gtrsim 10^{-27} ecm$; therefore, (3) provides a non-trivial constraint. It is amusing to note that in leptoquark models of CP-violation it is likely that $|d_e| \gtrsim |d_{neutron}|$. (Cf., e.g., [15,2].)

Flavor-conserving models

This denotes models where the most significant 1-loop contribution to d_e comes from Fig. 1 with F being not necessarily from a higher generation. Popular examples are i) Supersymmetric models: Here F is e.g. a neutralino (i.e., photino $\tilde{\gamma}$, zino,...) and B a selectron $\tilde{e}$ (or F=chargino, B= sneutrino). CP violation specific to SUSY can arise as follows. Neutralinos, being Majorana particles can have complex mass terms, e.g. $m_{\tilde{\gamma}} = |m_{\tilde{\gamma}}| \exp(i\phi_\gamma)$. In the basis where the $\tilde{\gamma}$ mass is real, this phase is shifted into the $\tilde{e}e\tilde{\gamma}$ interaction. Moreover the $\tilde{e}_L \leftrightarrow \tilde{e}_R$ selectron mixing parameter A will be in general complex: $A = |A| \exp(i\phi_A)$. There is not general reason why these phases should be small. These and other "SUSY phases" generate fermion EDMs [16] at 1 loop irrespective of generation mixing. Assuming that $|m_{\tilde{\gamma}}| \approx |A| \approx m_{\tilde{e}} \approx 100 GeV$ (suggested by "naturalness") one gets for the photino contribution to d_e:

$$d_e \simeq 10^{-25} \left(\frac{100 GeV}{|m_{\tilde{\gamma}}|}\right)^3 \sin(\phi_A - \phi_\gamma) ecm. \quad (5)$$

This and other contributions to d_e, and moreover the more stringent corresponding analysis for the neutron EDM, where gluino contributions dominate [16-18] indicates a problem for SUSY models: SUSY phases must be small for phenomenological reasons but what is the theoretical reason for it?

ii) Left-right-symmetric models. Here F is the electron-neutrino (more precisely: the light ν_{eL} slightly mixed with a heavy N_{eR}) and B is the lightest mass-eigenstate W boson W_1 or W_2 which are mixtures of W_L and W_R [19,20]. A nonzero $d_e$at 1 loop requires in addition to $W_L - W_R$ mixing also a complex Dirac neutrino mass μ_D but no generation mixing. One gets

$$d_e = 2 \times 10^{-24} F \left(\frac{m_{N_e}}{m_W}\right) \sin(2\xi) \left(\frac{Im \mu_D}{1 MeV}\right) ecm \quad (6)$$

where $F = O(1)$ and for the $W_L - W_R$ mixing parameter ζ one has the phenomenological constraint $|\sin 2\zeta| < 0.008$ (cf., e.g., [2]). That is, d_e could be as large as a few $times 10^{-27} ecm$ in these models, but then right-handed charged weak currents should show up somewhere else, too. For Higgs contributions to d_e in these models see [19,21].

iii) Higgs models of CP violation. The simplest of these models which is phenomenologically viable is the 2 Higgs doublet extension of the SM with natural flavor conservation at the tree level. The spectrum of physical Higgs particles consists of a charged Higgs $H^\pm$ and 3 neutrals: h, H with CP parity +1 and A with CP parity -1. In general the scalar self interaction will contain terms which induce CP-violating mixing [23] of h, H with A (unless forbidden "by hand" by imposing a discrete symmetry).

No a priori reason is known why this mixing should be suppressed. It leads to 3 neutral mass eigenstates $\phi_i (i = 1,2,3)$ with undefined CP parity; i.e. they couple both to scalar and pseudoscalar quark and lepton densities with coupling proportional to the fermion mass. That is, besides the KM mechanism, this model has also CP violation which is not connected to generation mixing and which leads to larger effects for heavy flavors. At 1-loop d_e is generated by Fig. 1b, where $B = \phi_i$ and $F = e$. It is very small, $d_e \sim G_F m_e^3/16\pi^2 m_\phi^2 \approx 10^{-35} ecm$ (where ϕ is the lightest Higgs); 2 powers of m_e resulting from the ϕee couplings and one power from the chirality flip. However, it has been recently observed [24] that this suppression is overcome for a class of amplitudes at 2 loops: At 1-loop $\gamma\gamma\phi$ vertices are induced by virtual t and W particles which in turn generate a 2-loop EDM $d_e \sim m_e$! This effect was calculated in [24] for the model where one doublet couples to up-type quarks, the other one to down-type quarks and leptons. In subsequent publications the W-loop contribution was treated more carefully [25] and other relevant contributions were calculated [25-27]. Definite numerical predictions are hampered by unknown parameters. Even if one assumes that the lightest ϕ dominates the effect d_e depends on 4 unknowns (asuming $m_t = O(120\ GeV)$) : m_ϕ, the ratio of vacuum expectation values $|v_2/v_1|$ and the dimensionless parameters $Im\ Z_0$ and $Im\ \tilde{Z}_0$ measuring CP violation in Higgs mixing [28]. Taking $|v_2| >> |v_1|$ and $|ImZ_0|$ near its upper bound of about $|v_2/2v_1|$ [28,26] one gets $ImZ_0 \simeq -Im\tilde{Z}_0$[28].Choosing e.g. $|v_2/v_1| = 10$ and $m_\phi \gtrsim 3m_W$ results in $|d_e| \gtrsim 2\times10^{-26} ecm$ [25,26] which is in conflict with the experimental upper limit resulting from eq. (3). Ref. [26] points out that for other parameter values $|d_e|$ is likely to be somewhat smaller even if CP violation is maximal: namely, if $2m_w \lesssim m_{\phi i} \lesssim 1TeV$ (i=1,2,3), $|d_e|_{max}$ is of order $a\ few \times 10^{-27}\ ecm$ at $|v_2/v_1| = 10$. This is because $d_e = 0$ if the $\phi_i (i = 1,2,3)$ are mass-degenerate, due to sum rules for the CP-violating parameters [28]. Nevertheless, improvement of the limit (3) by about a factor of 10 would place a stringent constraint in particular on this kind of CP violation.

DIPOLE MOMENTS OF μ AND τ

The EDMs of the μ and the τ could be substantially larger than d_e depending on whether $d_\ell \sim m_\ell^p$ with $p \geq 1$ [2,22]. In the Higgs model discussed above $(d_\ell)_{1-loop} \sim m_\ell^3$, whereas $(d_\ell)_{2-loop} \sim m_\ell$. Even if one assumes that $d_e = 10^{-26} ecm$, this model then does not predict d_μ and d_τ large enough for being measurable in the foreseeable future. A value somewhat larger than for d_τ results for the CP-violating weak dipole form factor (WDM) $\tilde{d}_\tau = \tilde{d}_\tau(m_z^2)$ which can be present in the $Z\bar{\tau}\tau$ vertex and can be measured at LEP with some precision [29]. One obtains in the 2-doublet model that $\tilde{d}_\tau \simeq -2.6 \cdot 10^{-22} \sin^2\beta ImZ_0\ ecm$ for $m_\phi = 100\ GeV$ [30]. Values $|\tilde{d}_\tau| > 10^{-20}\ ecm$ can be obtained in specific leptoquark-type models of CP violation. So far one knows from comparing the contribution $\Delta\Gamma(Z \to \tau^+\tau^-) \simeq 2(m_z \tilde{d}_\tau s_\theta c_\theta/e)^2/3$ (in GeV) [31] with (Γ_{exp} – Γ_{SM}) one obtains $|\tilde{d}_\tau| < 6.5\times10^{-17}\ ecm$ (1s.d.)[32]. A much more sensitive determination of $\tilde{d}_\tau$ can be made by measuring CP-odd tensor correlations [29] in $Z \to \tau^+\tau^-, \tau \to \pi\nu_\tau, \tau \to \ell\bar{\nu}_\ell\nu_\tau$, etc., at the Z peak at LEP; e.g. the CP-odd tensor

$$T_{ij} = (\hat{q}_+ - \hat{q}_-)_i \hat{n}_j + (i \leftrightarrow j)$$

where $\hat{q}_{\pm i}$ are the Cartesian components of the momentum directions of $\pi^\pm$ or $\ell^\pm$ and $\hat{\mathbf{n}} \sim (\hat{\mathbf{q}}_+ \times \hat{\mathbf{q}}_-)$. The expectation values $< T_{ij} > \sim \tilde{d}_\tau$ have been calculated for $\tau \to \pi, \ell, \rho$ [33] and one expects a sensitivity $\delta(\tilde{d}_\tau) \sim (2-3) \times 10^{-18}$ ecm if $10^7 Z$'s are produced.

SUMMARY

Experimental data on the electron EDM have become — like those for the neutron EDM — an important touchstone for non-standard models of CP violation; notably for supersymmetric and Higgs models. Further improvement of the experimental sensitivity would be highly desirable, as it will help to answer the question whether CP-violating forces other than those generated by the KM mechanism play any role in Nature or not. The EDMs d_μ, d_τ and $WDM\tilde{d}_\tau$ might be considerably larger than d_e. However, the precision with which these moments can be measured at present is 7-8 orders of magnitude smaller than for d_e. Nevertheless, it seems worthwhile to (continue to) measure these moments, too, and try to find means which enhance the sensitivity to them.

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DISCUSSION

Q. **B. McKellar** *(Univ. Melbourne)*: Comment: CP violation of nucleon is also important. One needs to study these effects in order to get correct bound on the electron electric dipole moment.

A. **W. Berneuther**: Yes. If there are enhancements on the CP violating nucleon effects. There are difficulties to extract electron electric dipole moment.

TRANSVERSE POLARIZATION AS A PROBE FOR CP-VIOLATION FROM NEW PHYSICS*

C.P. BURGESS and J.A. ROBINSON
Physics Department, McGill University, 3600 University Street
Montréal, Québec, H3A 2T8, Canada.

ABSTRACT

We use the effective-lagrangian formalism to examine the effects of CP-violation from outside the standard model at LEP and HERA. We then define a simple observable that may be defined when the initial electron beam is transversely polarized, and systematically determine which effective interactions can contribute to it. We find that for a wide class of theories, the observable only gets contributions from a single CP-violating electron-gauge boson coupling within the effective theory. We conclude that detectable CP-violation from this source may be observable at LEP or HERA.

I. INTRODUCTION

The spectacular success of the standard model suggests that the energy scale of the next level of new physics is likely to be high compared to those currently experimentally accessible. This invites the application of effective-lagrangian[1] techniques to exploit this large ratio of scales. Since the standard model is itself the most general renormalizable theory consistent with its particle content and gauge symmetries, its success can be understood as simply being the consequence of the scale of this new physics being so large.

The biggest uncertainty faced when writing down a general effective lagrangian is the question of whether to include the Higgs particle. The least model-dependent attitude would be to not introduce an elementary Higgs particle, but to merely require that the effective theory include the Goldstone bosons associated with the spontaneous breaking of $SU_L(2) \times U_Y(1)$. In this case only the unbroken $SU_c(3) \times U_{em}(1)$ subgroup need be linearly realized. The alternative approach is to linearly realize the full standard-model gauge group, $SU_c(3) \times SU_L(2) \times U_Y(1)$, and include symmetry-breaking effects through the explicit dependence of the effective lagrangian on the Higgs field(s). This method has the advantage of incorporating the additional information that the order parameter of the electroweak symmetry breaking is one (or several) $SU_L(2)$-doublet(s).

The apparent reliance on the existence of a physical Higgs boson in the second of these techniques is illusory to the extent that propagating Higgs degrees of freedom are not important for the processes under consideration. In what follows the Higgs field only contributes through its vacuum-expectation-value (*v.e.v.*), and so just plays the role of order parameter for $SU_L(2) \times U_Y(1)$ breaking.

We ask here whether upcoming experiments at HERA[2] or at LEP/SLC[3] are likely to detect any of the CP-violating terms that can appear as dimension-five and -six effective interactions. The goal is to determine whether an visible signal (in as simple an observable as possible) can be consistent with other bounds, in particular those due to limits on electric dipole moments for elementary fermions. More details of the arguments and calculations described here may be found in Refs. [2] (HERA) and [3] (LEP/SLC). A summary of these results with a review of the effective-lagrangian formalism is given in Ref. [4]. Further applications to $W^{\pm}$ and Z^0 moments appear in Refs. [5].

In searching for signatures of CP violation, we take advantage of the fact that transversely polarized electron beams are expected to be available at both ep and e^+e^- machines in the near future.[6] Consider, then, a two-body process such as pair production, $e^+e^- \to f\bar{f}$, at SLC or LEP, or deep-inelastic scattering, $e^{\pm}p \to \ell f$, at HERA. Given initially transversely polarized electrons or positrons a convenient observable consists of the asymmetry defined as the difference: $dA(\mathbf{p}_i, \mathbf{s}_i) =$

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$d\sigma(\mathbf{p}_i, \mathbf{s}_i) - d\sigma(-\mathbf{p}_i, -\mathbf{s}_i)$, in the differential cross-section before and after reversing all momenta and spins. An integrated version of this observable is given in terms of the integrated luminosity, L, by: $\mathcal{A} = LA$. It may be defined operationally as the difference between the number of final-state fermions, f, detected on either side of the plane defined by the initial transverse spin direction and the colliding-beam axis—or, equivalently the difference between the number of such fermions appearing out of one side of this plane as the initial electron polarization is reversed.

This observable is sensitive to time-reversal-odd interactions because time-reversal reverses all three-momenta, $\mathbf{p}_i$, and spins, $\mathbf{s}_i$. Strictly speaking, however, the observation of a nonzero dA need not signal the presence of T-violation because the action of time reversal would run the reaction backwards as well as flipping all momenta and spins. As a result the potential effects of any T-preserving interactions must be carefully considered when considering such triple products. This type of background turns out to be negligibly small within the standard model.

A great simplification arises once we focus only on those terms in the effective lagrangian that can contribute to dA. In this case all but two of the many effective operators are suppressed in their tree-level contribution to $dA/d\theta d\phi$ by powers of light quark or lepton masses. This claim is justified in more detail in refs. [2] and [3]. The two effective interactions that can appear may be written (up to integrations by parts and field redefinitions) as a linear combination of: $\mathcal{L}_W = \lambda_W \left[\left(\overline{L} P_R D^\mu E\right) D_\mu \phi\right] +$ c.c. and $\mathcal{L}_\gamma = \frac{i}{2}\lambda_\gamma \left[g_1 B_{\mu\nu} \left(\overline{L}\gamma^{\mu\nu} P_R E\right)\phi\right] +$ c.c.. In our conventions, the CP-odd part of these interactions is proportional to the imaginary part of the coefficients, λ_W and λ_γ.

Im λ_γ is strongly bounded by the electron e.d.m.(see below). The contribution of the remaining operator[3] to dA (in the CM frame) are easily computed, and are summarized for LEP in the following table:

Fermion:	$\tau^+\tau^-$	$c\bar{c}$	$b\bar{b}$
SM Events, ($N \times 10^{-8}$):	6.5	24	31
Asymmetry ($\mathcal{A} \times 10^{-5}$):	−6.1	−97	−170
$\mathcal{A}/N \times 10^3$:	0.94	4.0	5.5

Table (1): Integrated Asymmetries at LEP

This table assumes $\lambda_W = (400\ GeV)^{-2}$ and an integrated luminosity of [7] $L = 4.8 \cdot 10^5\ pb^{-1}$ at LEP. The smaller luminosity at SLC makes it is impossible to distinguish a signal from standard-model statistical fluctuations.

In order to quantitatively estimate how big λ_W must be to be detectable at LEP let us assume that the systematic error in the measurement of the asymmetry, A, in the lepton- and heavy-quark-production cross-section can be reduced to 0.1% of the total cross-section. Systematic errors of roughly this size are expected[8] for LEP measurements of the forward-backward asymmetry, A_{FB}. In this case the two-jet asymmetry would be just visible at the 2-σ level above background provided that: $\lambda_W > (570\ GeV)^{-2}$, for up-type quarks; and $\lambda_W > (660\ GeV)^{-2}$, for down-type quarks.

The tree-level expression for the asymmetry in e^-p collisions may be similarly found[2]. We have numerically integrated these expressions using $\lambda_W = (400\ GeV)^{-2}$, and a HERA luminosity of 200 pb^{-1}. The charged-current signal is greater than twice the standard-model statistical background if: Im $\lambda_W > \left(\frac{1}{600\ GeV}\right)^2$ in e^-p collisions and Im $\lambda_W > \left(\frac{1}{330\ GeV}\right)^2$; for e^+p.

Should a nonzero asymmetry be observed there are several ways to determine whether the effective interaction, $\mathcal{L}_W$, is responsible: (a) The angular (θ- and ϕ-) dependence of the asymmetry at LEP, or the x- and Q^2-dependence at HERA, must be as predicted since the only unknown parameter, λ_W, enter into these expressions through an overall multiplicative factor. (b) If both the electron and positron beams at LEP are transversely polarized then the asymmetry must be proportional to the *difference* between electron and positron polarization, $\mathbf{s}_{e^-} - \mathbf{s}_{e^+}$. At both HERA and LEP A must change sign with the initial polarization. (c) $SU_L(2) \times U_Y(1)$-invariance implies that the effective interaction, $\mathcal{L}_W$, also predicts a CP-violating electron-photon-Z vertex which might be detectable through asymmetries in the process $e^+e^- \to f\bar{f}\gamma$ or $e^\pm p \to \ell f\gamma$.

The main question is whether Im λ_W of this size is consistent with other experimental bounds. By far the most restrictive bound is the present limit on the intrinsic electric dipole moment of the electron: $d_e < 1.5 \cdot 10^{-26}$ (Berkeley[9]), $d_e < 1.2 \cdot 10^{-25}$ (Amherst[10]). Other observables, such as limits on triple-product correlations in μ-decays, or limits on deviations from standard-model predictions of Z^0-partial widths, $\Gamma(Z \to f\bar{f})$, are not sufficiently sensitive to significantly constrain λ_W.

The e.d.m.-bound completely eliminates the coupling Im λ_γ. This is because although many effective operators can contribute to the electron e.d.m., they cannot cancel the effect due to Im λ_γ without undermining perturbation theory in our small expansion parameters—the small gauge and Yukawa couplings. Since Im λ_γ is the coefficient of the sole operator that contributes to d_e at tree level and to leading order in m_e/v it must be bounded to be smaller than: Im $\lambda_\gamma <$ $(10^3\, TeV)^{-2}$.

The remaining coupling, Im λ_W, is only one of several interactions that contribute to d_e at one-loop and at leading order in m_e/v and so its contribution to the electron e.d.m. could potentially cancel with others in at this order. If we assume that each of these operators is separately bounded by its individual contribution to the electron e.d.m. we derive a bound of: Im $\lambda_W <$ $(10\, TeV)^{-2}$. This would be sufficiently strong to rule out any observable asymmetry at all at either LEP or HERA.

It should be borne in mind, however, that this bound would not apply if the contribution of Im λ_W to d_e should cancel with that of another operator of the effective lagrangian. Indeed, the occurence of such cancellations tell us a great deal about the symmetries and selection rules of the underlying theory. In the present instance this would require a cancellation of the contribution to within a little better than the percent level. Although a cancellation to this accuracy appears unnatural in the effective lagrangian approach, it need not be within the context of a specific model where the effective couplings are all determined in terms of a small number of underlying parameters. The most famous example in which this occurs is the standard model itself where the cancellation reflects the GIM mechanism of the underlying theory. The observation of the asymmetry described here would point to similar cancellations at work at the scale of new physics.

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Q. **Kam-Biu Luk** *(UC, Berkeley, LBL)*: If I understand you correctly, you are using triple-product to look for CP violation. If so, how will you compare your approach with semi-leptonic decays, for instance, beta decay of kaon?

A. **C. P. Burgess**: Yes, it is a triple-product test for CP violation. Semi-leptonic decay of kaon involves light masses that do not give good limit on the parameters in the lagrangian. Other decays like Z^0 decay give comparable constraints on the parameters as from the electric dipole moment of the electron.

DIRECT CP VIOLATION IN $K_{L,S} \to 2\gamma$ TRANSITIONS FOR AN ARBITRARY TOP MASS *

J. O. EEG

Institute of Physics, University of Oslo, N-0316 Oslo 3, Norway

B.NIŽIĆ and I. PICEK

Theoretical Physics Department, Rugjer Bošković Institute, Zagreb, Croatia, Yugoslavia

We study short-distance contributions to $K^0(\overline{K^0}) \to \gamma\gamma$ within the standard model for large m_t, including QCD corrections. Owing to a relaxed GIM suppression, CP violation in the decay amplitude (direct CP violation) turns out to be more important than previously thought and might be of order 20% of that coming from the mass matrix (indirect CP violation). This leads to rare $K_{L,S} \to \gamma\gamma$ decays as potentially interesting candidates for testing direct CP violation at LEAR and Φ factories. Still, it does not seem to be achievable by the luminosities foreseeable for forthcoming Φ factories.

1. INTRODUCTION

The results of two existing experiments from CERN and FNAL [1] on direct CP violation in $K_L \to \pi\pi$ decays disagree. Also, theoretically, the SM prediction for $\frac{\epsilon'}{\epsilon}$ is rather sensitive to the yet unknown value of the top mass [2]. Therefore, a study of other rare decay-modes might decide both on the existence of direct CP violation and possibly give an insight into "new physics". Some time ago, two of us [3] discussed $K_{L,S} \to 2\gamma$ as a possible candidate in this respect. More recently [4], this process has been revived in the light of the current proposals for Φ factories and a significant increase in the lower bound on the top mass.

The possibilities of observing CP violation in $K_{L,S} \to 2\gamma$ were discussed before [5] and after the advent of gauge theories [6,7]. Quite surprisingly, both the rates and the CP-violating amplitudes were attributed to long-distance (LD) pole contributions. Since these suffer from a substantial uncertainty (owing to $\eta - \eta'$ mixing [8]), we continue the previously initiated [3] study of short-distance (SD) CP-violating amplitudes.

The hope that one could tell direct CP violation from signals in $K_{L,S} \to 2\gamma$ arose from the measurement of the partial $K_S \to 2\gamma$ lifetime which does not overwhelm $K_L \to 2\gamma$. Measuring the time-dependent asymmetry of K^0 and $\overline{K^0}$ intensities at LEAR and Φ factories could then determine the direct CP-violating parameter $\epsilon'_{\gamma\gamma(-)}$, an analogue of the parameter $\epsilon'_{\pi\pi} \equiv \epsilon'$ of the 2π decay mode:

$$\eta_- \equiv \eta_{\gamma\gamma(-)} = \frac{A(K_S \to \gamma\gamma(-))}{A(K_L \to \gamma\gamma(-))} = \epsilon + \epsilon'_{\gamma\gamma(-)}. \quad (1.1)$$

A similar parameter for $CP = +1$ photons is esentially determined by indirect CP violation [9].

2. 1PI AMPLITUDE FOR AN ARBITRARY TOP MASS

Here we consider the CP-violating part of the electroweak $s\bar{d} \to \gamma\gamma(-)$ amplitude for an arbitrary heavy top quark. For this purpose, we take under scrutiny the full set of electroweak diagrams shown in Fig.1.

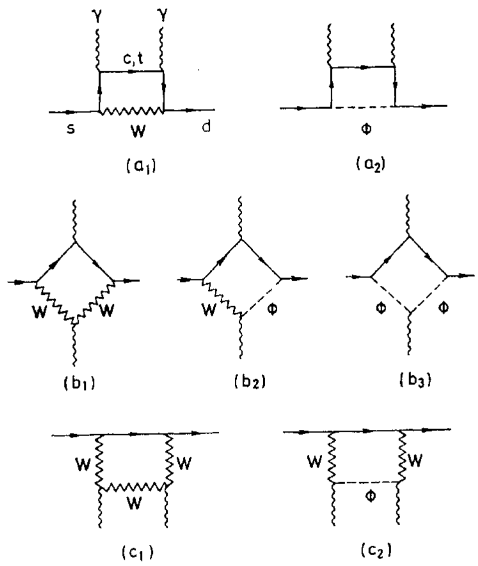

Fig.1 Electroweak Feynman diagrams contributing to $s\bar{d} \to \gamma\gamma(-)$ in the R_ξ gauge .

Presented by I. Picek

These diagrams were considered by Gaillard and Lee [6] for the CP-conserving amplitude. In the light-quark approximation $m_q^2 \ll m_W^2$, they found that diagram (a1) of Fig.1 was most important. In ref. [7] the result of ref. [6] was implemented for the CP-violating case and a negligible SD CP-violating amplitude obtained. By calculating diagram Fig.1(a1) for arbitrary quark mass, we find that this result of ref. [7] is completely changed. We then calculate all the other diagrams in Fig.1 and, finally, consider the role of QCD corrections.

The amplitude to lowest order in photon momenta, with $CP = -1$, has the form

$$M(s\bar{d} \to \gamma\gamma(-)) = \frac{G_F}{\sqrt{2}} \frac{\alpha_{em}}{\pi} \left(\sum_{q=u,c,t} \lambda_q A_q \right) \epsilon^{\mu\nu\alpha\beta} \epsilon^*_{1,\mu} \epsilon^*_{2,\nu} (k_1 - k_2)_\alpha (\bar{d}_L \gamma_\beta s_L) , \quad (2.1)$$

where $\lambda_q = V_{qs} V^*_{qd}$ $(q = u, c, t)$ is the appropriate KM factor and G_F is the Fermi coupling constant. The polarization vectors and the momenta of the photons are denoted by ϵ_j and k_j $(j = 1, 2)$, respectively. Using the relation $\lambda_u + \lambda_c + \lambda_t = 0$, we decompose the contributions into CP-conserving and CP-violating parts as follows:

$$\sum_{q=u,c,t} \lambda_q A_q = \lambda_u (A_u - A_c) - \lambda_t (A_c - A_t) . \quad (2.2)$$

The CP-violating quantity $\epsilon'_{\gamma\gamma(-)}$ defined in eq.(1.1) requires calculation of the one-particle irreducible (1PI) contributions to A_c and A_t via the electroweak Feynman diagrams represented in Fig.1.

The result for the CP-violating amplitude $A_{EW} = A_c - A_t$ in eq.(2.2) is depicted in Fig.2 for an arbitrary top mass.

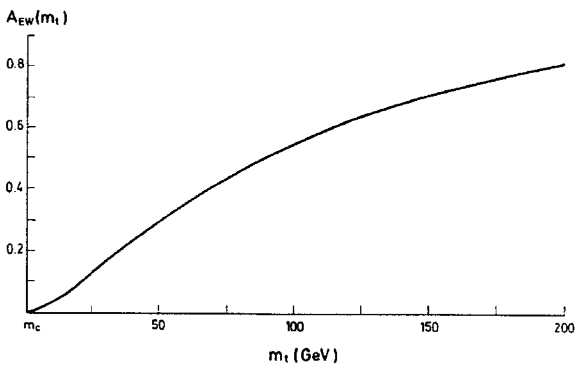

Fig.2 Behaviour of the 1PI amplitude $A_{EW}(m_t)$.

For m_t in the range 70 GeV $< m_t <$ 200 GeV, we find

$$A_{EW}(m_t) \simeq +0.4 \ \text{to} \ +0.8 . \quad (2.3)$$

Using the 't Hooft-Feynman gauge for the W-boson propagator, we obtain for the first diagram (a1) in Fig.1 for $m_t^2 = M_W^2 \gg m_c^2 \gg m_K^2$:

$$A_{EW}^{(a1)} = [(-\frac{4}{9}) - (-\frac{4}{27})] . \quad (2.4)$$

The two terms in the parentheses correspond to the c- and t-contributions respectively. This result differs from the result $\sim 10^{-1} m_K^2/m_q^2$ from ref.[6].

In the neutral kaon decay at hand, the QED gauge invariance for the physical, on-shell process at the hadronic level , can be achieved by employing the matrix element of the quark current in eq. (2.1). For the underlying quark process under consideration, within the 't Hooft-Feynman gauge, diagram (c1) is most important and gives a positive contribution, such that the total contribution is positive, as shown in Fig.2.

It has been explained in a series of papers [10] that when the expression for a diagram is written as a loop integral over virtual (squared) loop momentum p^2, QCD corrections may be obtained in a rather intuitive way. Using these methods gives us that the dominant QCD correction is accounted for by effectively making the following replacement of the charm contribution $(-4/9)$ in eq.(2.4):

$$-\frac{4}{9} = +\frac{4}{9}(1-2) \to \frac{4}{9}(1-2\eta(m_c^2)) = -\frac{4}{9} + \frac{8}{9}(1-\eta(m_c^2)) \quad (2.5)$$

where

$$\eta(\mu^2) = 2C_+(\mu^2) - C_-(\mu^2) , \quad (2.6)$$

and $C_\pm(\mu^2)$ are coefficients of the well-known effective non-leptonic hamiltonian [11], giving $\eta(m_c^2) \simeq -0.1$. Thus A_c^{a1} changes sign when QCD corrections are taken into account.

We find that the important contributions from other diagrams are all dominated at momenta $p^2 \sim M_W^2 \sim m_t^2$. Thus, for all contributions except the one for (a1), QCD corrections should be rather small for $m_t \simeq M_W$. For $m_t = M_W$ we obtain

$$A_{EW+QCD}(m_t = M_W) = A_{EW}(m_t = M_W) + \frac{8}{9}(1 - \eta(m_c^2)) \simeq 1.4, \quad (2.7)$$

what represents an increase by a factor of $\simeq 3$ over the pure electroweak amplitude, $A_{EW}(m_t = M_W) = \frac{25}{54} \simeq 0.46$.

Replacing $(A_u - A_c)$ in (2.1)-(2.2) by the empirical value $A_{emp} \simeq 3.66$, needed to reproduce the total rate for $K_L \to \gamma\gamma$, we obtain

$$(\epsilon'_{\gamma\gamma(-)})_{SD} = \frac{\text{Im}(\lambda_t)}{\text{Re}(\lambda_u)} \frac{A_{EW+QCD}}{A_{emp}} \simeq 0.39 \frac{\text{Im}(\lambda_t)}{\text{Re}(\lambda_u)} , \tag{2.8}$$

where the numerical value corresponds to $m_t \simeq M_W$. This implies that even for the modest value $m_t = M_W$ the direct CP-violating amplitude might be 20% to 30% of that coming from the mass matrix.

To conclude, we have presented a calculation of SD contributions to the $s\bar{d} \to \gamma\gamma(-)$ transition for a heavy top. The diagrams in Fig.1 are shown to be much more important than previously thought. The example of the rather on-shell $b\bar{s} \to \gamma\gamma(-)$ quark process in the B meson shows the cancellation of the 1PI and 1PR contributions[12]. However, the off-shellness of light quarks in the kaon leads to survival of the unsuppressed $K_{L,S} \to \gamma\gamma(-)$ amplitude presented here.

In our opinion, therefore, the processes $K_{L,S} \to 2\gamma(-)$ considered in this paper are potentially good candidates to be tested at LEAR and forthcoming Φ factories. However, there are essential differences between these two facilities. Firstly, there is an enormous difference in tagging efficiency. Secondly, two-photon decays of kaons are suppressed (almost rare) modes, which further compensates the original nonsuppression by the $\Delta I = 1/2$ rule. In order to achieve the measurement accuracy of 1/10 in $\epsilon'_{\gamma\gamma(-)}$, one would need $10^{14}\Phi$ particles per year (in comparison with 10^6 events needed at LEAR). In contrast, the optimal integrated luminosity of 10^{40} over 100 cm can give only $10^{10}\Phi$ particles per year. Thus, it is a pity that the advent of the LHC facility at CERN seems to close the search of $K_{L,S} \to \gamma\gamma$ at LEAR, the best place to measure direct CP violation in $K_{L,S} \to \gamma\gamma$.

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