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INTRODUCTION

The intense and clean beams of antiprotons which will become a reality at LEAR (the Low-Energy Antiproton Ring) open up the possibility of resolving many of the problems in existing low-energy antiproton physics, which are largely brought about by a lack of flux. It is interesting to ask the question, "What else can be done at LEAR apart from a study of antiproton interactions?"

Together with Veneziano we have already pointed out¹ that the intense flux of pseudoscalar mesons at LEAR allows an experimental study of the strong CP problem through precise tests of the chiral predictions of the standard model by the mixing of pseudoscalars with ghosts and axions. A review of the strong CP problem and related physics has already been presented at this conference by Rossi² and will therefore not be considered further.

However, we would like to discuss the possibility of studying the weak CP violation through the neutral-kaon system at LEAR, because the annihilation of antiprotons at rest produces large fluxes of K^0 and $\bar{K}^0$ mesons of equal magnitude. The neutral-kaon system has been studied in great detail over the last twenty years, therefore it is important to ask what new physics can be learned from a more precise study of this system.

SOURCE OF CP VIOLATION

The parameters of CP violation can be found in many standard textbooks³⁻⁵ or reviews⁶⁻⁸, and only a brief outline will be presented

for clarity and completeness. In the neutral-kaon system, we can distinguish four states: K^0 and $\bar{K}^0$ are eigenstates of the strong interaction, while K_1^0 and K_2^0 are the $CP = \pm 1$ eigenstates. Allowing for CP non-conservation the *observed* decay eigenstates K_L and K_S are

$$|K_S\rangle = \{p|K_0\rangle - q|\bar{K}_0\rangle\} \quad (1)$$

$$|K_L\rangle = \{p|K_0\rangle + q|\bar{K}_0\rangle\} , \quad (2)$$

where, assuming TCP invariance, we have defined the strong eigenstates

$$|K_0\rangle = \frac{1}{2p} \left(|K_L\rangle + |K_S\rangle \right) \quad (3)$$

$$|\bar{K}_0\rangle = \frac{1}{2q} \left(|K_L\rangle - |K_S\rangle \right) , \quad (4)$$

with $p, q = (1 \pm \varepsilon)/\sqrt{2(1 + |\varepsilon|^2)}$. Here ε is the small complex amplitude specifying the CP impurity of the observed eigenstates (e.g. for $\varepsilon = 0$, $|K_1\rangle \equiv |K_S\rangle$).

Before the development of the electroweak theories there were two main approaches to the violation of CP invariance. The so-called "superweak" predictions postulated a new form of interaction with coupling strength $\sim 10^{-9} G_F$, where all the CP violation effects are entirely due to the impurity in the decay eigenstates K_S, K_L ($\langle K_S | K_L \rangle \neq 0$). In this case there is no direct CP violation in the decay of the CP eigenstates, i.e. $K_2(CP = -1)$ does not decay directly into two pions and the violation occurs only in the mass matrix due to the mixing of $|K_S\rangle$ and $|K_L\rangle$ ($\varepsilon \neq 0$). In terms of actual measurable amplitudes, it can be seen from Eqs. (1) and (2) that

$$\eta_{\pi\pi} = \frac{A(K_L \rightarrow 2\pi)}{A(K_S \rightarrow 2\pi)} = \varepsilon$$

giving

$$\eta_{+-} = \eta_{00} = \varepsilon ,$$

where η_{+-} and η_{00} are the charged and neutral 2π decay ratios.

In the second approach, the "milliweak" prediction is that direct CP violation is allowed, and therefore the source of CP violation occurs in both the mass matrix and directly in the decay matrix, i.e. $K_2^0 \rightarrow 2\pi$. In order to consider this case in terms of measurable quantities, it is convenient to decompose the two-pion final states into $I = 0$ and $I = 2$ components as given in Ref. 4, thus introducing ε' which denotes the ratio of the imaginary part

of the $I = 2$ to the real part of the $I = 0$ amplitude. We can then define the experimental decay rates in terms of

$$\begin{aligned}\eta_{+-} &= \epsilon + \epsilon' \\ \eta_{00} &= \epsilon - 2\epsilon'\end{aligned}\tag{5}$$

where the real part of the $I = 2$ $\pi\pi$ amplitude in the K_1 -decay ($\sim 5\%$ of the $I = 0$ amplitude) has been neglected.

Experimentally the present limits on $|\epsilon'/\epsilon|$ are approximately 2 per cent^{9,10} and therefore it is necessary to ask what we expect for this ratio since the introduction of the electroweak interaction and the standard model, which of course have no direct relation to CP violation in weak decays.

In present-day theories, CP violation can be introduced in several ways; in particular, (a) by enlarging the Higgs sector, (b) by including right-handed currents (RHC), and (c) by adding at least a new family of fermions. The first possibility includes two distinct classes of models. In one, flavour-changing couplings are allowed, requiring two Higgs doublets^{8,11,12}. However the additional neutral Higgs particles have to be heavy^{13,14} to prevent flavour-changing neutral currents. In the other, a minimum of three Higgs doublets are required, since no flavour changing couplings are allowed¹⁵⁻¹⁸. The second possibility introduces CP violation by the presence of RHC¹⁹⁻²¹ where the violation is brought about by the phase clash between the parity-conserving and parity-violating parts of the effective weak Hamiltonian. The third possibility is the six-quark Kobayashi-Maskawa^{22,23} model, where the CP violation occurs naturally as a consequence of weak mixing between the quarks analogous to the extension of the Cabibbo angle to the Glashow-Iliopoulos-Maiani (GIM) model. The $2/3$ quarks and $1/3$ quarks are then related via a matrix containing three real angles, and one complex phase $e^{i\theta}$ which accounts for CP non-conservation. It is possible to determine in which quadrant δ lies from a measurement of the sign of $\text{Re } \epsilon'$.

The theories and their predictions are summarized in Table 1, and it is obvious that a precise determination of the CP-violation parameter $|\epsilon'/\epsilon|$ to a level of ~ 0.002 will be required in order to quantify the source of CP violation in our present understanding of weak decays.

From our present knowledge of the production and decay of c and b quark states in known and planned accelerators and storage rings, it will be extremely difficult to determine the source of CP violation in heavy quark decays, since the asymmetry is predicted to be $\sim 10^{-4}$ for leptonic decays. The GIM mechanism keeps the mixing small in these heavy quark systems. In addition, the knowledge of

Table 1. Predictions of CP-violating parameters

Model	$ \epsilon'/\epsilon $
Superweak	0
Kobayashi-Maskawa	0.002 to 0.02
Higgs A	~ 0
Higgs B	0.02 to 0.05
RHC	0

the origin of CP non-conservation in Nature will increase the understanding of physics beyond the standard model, since all CP models have an impact on problems related to grand unification. Following these considerations, schematically presented in Fig. 1, it is extremely interesting to note that low-energy machines like LEAR can play an important role in our understanding of grand unification.

Fig. 1. Unification in particle physics.

CP VIOLATION AT LEAR

It is our contention that by using the high-intensity source of LEAR, it is possible to produce a well-defined source of K^0 and $\bar{K}^0$ mesons through the reaction at rest

$$\bar{p} + p \text{ (annihilation at rest)} \rightarrow \left. \begin{aligned} &K^+ \pi^- \bar{K}^0 \\ &\rightarrow K^- \pi^+ K^0 \end{aligned} \right\} \quad (6)$$

where the total reaction rate in reaction (6) occurs at 4×10^{-3} of stopped antiprotons. Assuming a production rate of $10^6 \bar{p}/s$ ($\sim 10^{11} \bar{p}/\text{day}$) it is possible to produce $\sim 2 \times 10^8 K^0$ and $\bar{K}^0$ per day. The efficiency of measurement of the K^0 and $\bar{K}^0$ decays can be treated identically and the signature of the K^0 and $\bar{K}^0$ can be subsequently tagged by a determination of the sign of the charged kaon. In this way it is possible to minimize many of the systematic errors by measuring differences and sums relating to the K^0 and $\bar{K}^0$ initial states for the neutral and charged two-pion decays of the K_L^0 and K_S^0 systems.

The decay rate of a K^0 into two pions is given as a function of time by

$$R_{\pi\pi}^0(t) = \frac{1}{4|p|^2} \left\{ R_{K_S} e^{-\gamma_S t} + R_{K_L} e^{-\gamma_L t} + 2|\eta_{\pi\pi}| R_{K_S} e^{-[(\gamma_S + \gamma_L)/2]t} \cos(\Delta m t - \theta_{\pi\pi}) \right\}, \quad (7)$$

where the rate is defined for either $\pi^+\pi^-$ or $\pi^0\pi^0$. Δm is the $K_L - K_S$ mass difference and $\theta_{\pi\pi}$ is the phase of the complex parameter $\eta_{\pi\pi}$. The corresponding expression for $\bar{K}^0$ is obtained by reversing the sign of the last term and by replacing $|p|^2$ by $|q|^2$ in Eq. (7). We then obtain for the sum and difference of the K^0 and $\bar{K}^0$ rates the following expressions

$$\begin{aligned} S_{\pi\pi}(t) &= R_{\pi\pi}^0(t) + \bar{R}_{\pi\pi}^0(t) = R_{K_S} e^{-\gamma_S t} + R_{K_L} e^{-\gamma_L t} - \\ &\quad - 4 \operatorname{Re} \epsilon |\eta_{\pi\pi}| R_{K_S} e^{-[(\gamma_S + \gamma_L)/2]t} \cos(\Delta m t - \theta_{\pi\pi}) \quad (8) \\ &\approx R_{K_S} e^{-\gamma_S t} + R_{K_L} e^{-\gamma_L t} \end{aligned}$$

$$\begin{aligned}
D_{\pi\pi}(t) &= R_{\pi\pi}^0(t) - \overline{R_{\pi\pi}^0}(t) = 2|\eta_{\pi\pi}| R_{K_S} e^{-[(\gamma_S + \gamma_L)/2]t} \cos(\Delta mt - \theta_{\pi\pi}) - \\
&- 2 \operatorname{Re} \epsilon \left[R_{K_S} e^{-\gamma_S t} + R_{K_L} e^{-\gamma_L t} \right] \\
&= 2 R_{K_S} |\eta_{\pi\pi}| e^{-[(\gamma_S + \gamma_L)/2]t} \cos(\Delta mt - \theta_{\pi\pi}) - 2 \operatorname{Re} \epsilon S_{\pi\pi}(t)
\end{aligned} \tag{9}$$

In Eq. (8) the last term is of the order of a few 10^{-6} and thus neglected. Starting from K^0 and $\overline{K}^0$ states it is therefore possible to define a time-dependent asymmetry factor between K^0 and $\overline{K}^0$ rates given by

$$\begin{aligned}
A_{\pi\pi}(t) &= \frac{R_{\pi\pi}^0(t) - \overline{R_{\pi\pi}^0}(t)}{R_{\pi\pi}^0(t) + \overline{R_{\pi\pi}^0}(t)} = \frac{D_{\pi\pi}(t)}{S_{\pi\pi}(t)} \\
&= 2 \left[\frac{|\eta_{\pi\pi}| e^{(\gamma_S/2)t} \cos(\Delta mt - \theta_{\pi\pi})}{1 + |\eta_{\pi\pi}|^2 e^{\gamma_S t}} - \operatorname{Re} \epsilon \right]
\end{aligned} \tag{10}$$

where we have neglected γ_L compared to γ_S . Finally, we can define a time-dependent factor $U(t)$ giving the deviation from unity of the K^0 and $\overline{K}^0$ rates by

$$\begin{aligned}
U_{\pi\pi}(t) &= 1 - \frac{R_{\pi\pi}^0(t)}{\overline{R_{\pi\pi}^0}(t)} \\
&= 4 \left[\operatorname{Re} \epsilon - \frac{|\eta_{\pi\pi}| e^{(\gamma_S/2)t} \cos(\Delta mt - \theta_{\pi\pi}) [1 - 4 \operatorname{Re} \epsilon]}{1 + |\eta_{\pi\pi}|^2 e^{\gamma_S t} - 2|\eta_{\pi\pi}| e^{(\gamma_S/2)t} \cos(\Delta mt - \theta_{\pi\pi})} \right]
\end{aligned} \tag{11}$$

where we have neglected γ_L compared to γ_S and higher orders of $\operatorname{Re} \epsilon$ compared to unity.

The time-dependent rates $R_{\pi\pi}^0(t)$, $\overline{R_{\pi\pi}^0}(t)$, $A_{\pi\pi}(t)$, and $U_{\pi\pi}(t)$ are illustrated in Figs. 2 to 5 from which it is possible to extract $\eta_{\pi\pi}$, $\theta_{\pi\pi}$, and ϵ , since $\eta_{\pi\pi}$ and $\theta_{\pi\pi}$ determine the amplitude and position of the interference effect, respectively, and ϵ determines the time-independent deviation from zero. It can also be seen that any measurement should cover the time interval out to $\sim 20 K_S^0$ lifetimes (τ_S) in order to cover the region of maximum interference at $\sim 12 \tau_S$ and later times where the K_L decays dominate. We now con-

Fig. 2 Rates of the decay $K^0 \rightarrow \pi^+ \pi^-$ as a function of time for a 10-day run at LEAR ($\sim 2 \times 10^9 K^0$'s).

Fig. 3 Rates of the decay $\bar{K}^0 \rightarrow \pi^+ \pi^-$ as a function of time for a 10-day run at LEAR ($\sim 2 \times 10^9 \bar{K}^0$'s).

Fig. 4 The asymmetry factor $A_{\pi^+ \pi^-}(t)$ [Eq. 10]. The statistical errors correspond to a 10-day run at LEAR ($\sim 2 \times 10^9 K^0$'s and $\bar{K}^0$'s).

Fig. 5 The factor $U_{\pi^+ \pi^-}(t)$ [Eq. 11]. The statistical errors correspond to a 10-day run at LEAR ($\sim 2 \times 10^9 K^0$'s and $\bar{K}^0$'s).

sider a detector consisting of a magnetic spectrometer with a momentum resolution of a few per cent, which will enable us to define the momentum and direction of the K^0 and detect the decays over a distance of ~ 60 cm path length. This will require a photon detector of good energy and spatial resolution which can detect photons down to low energies, e.g. scintillation glass or BGO. It is important to define at the vertex of the stopped antiprotons the charged kaon in the trigger and this can be done over the momentum range required ($\lesssim 750$ MeV/c) using dE/dx chambers close to the small target and conventional Čerenkov detectors, e.g. lucite, H_2O , at the outer radius of the detector. The magnetic analysis of the charged decay pions is not necessary for the decay-vertex determination or for identification of the $\pi^+\pi^-$ decay of the K^0 , $\bar{K}^0$, but it does provide a constrained fit which will help in the understanding of background and systematics. A further experimental constraint is imposed by the number of events which can be handled in the analysis, and therefore three different methods are presented here, which indicate the feasibility of obtaining the necessary precision in $|\epsilon'/\epsilon|$ without any consideration of systematic uncertainties. Two of these measure the time dependence and the third measures the total integrated rates.

Method I: The important point is that the sum $S_{\pi\pi}(t)$ of K^0 and $\bar{K}^0$ rates [Eq. (8)] represents the pure exponential decays of K_S and K_L . The neutral-kaon rate, without distinguishing K^0 and $\bar{K}^0$, allows a precise determination of $|\eta_{\pi\pi}|^2 = R(K_L \rightarrow \pi\pi)/R(K_S \rightarrow \pi\pi)$.

A resolution of a few millimetres in defining the decay vertex for the charged decays results in a definition of $|\eta_{+-}|^2$ with an accuracy of 5×10^{-3} for a 10-day run at LEAR (2×10^9 K^0 's).

The decay-vertex reconstruction for the neutral π^0 decays is different. Here the positions where the gammas enter the calorimeter, as well as their energies, must be measured with good resolution. The knowledge of the neutral kaon momentum vector allows then the determination of the vertex by a 5C fit. The resolution in the vertex reconstruction obtained by Monte Carlo simulation is shown in Fig. 6 for a scintillation-glass and BGO calorimeter. Fitting techniques developed for time spectra of nuclear excited states provide a precise determination of the two exponential functions [Eq. (8)], even if the spectra are smeared. Thus the precision in defining $|\eta_{00}|^2$ is close to that for $|\eta_{+-}|^2$. The desired accuracy for $|\eta_{\pi\pi}|^2$ can be achieved by using only the time region $\geq 6 \tau_S$ and thus overcoming any problem arising from the number of events to be handled in the analysis.

The definition of ϵ' in Eq. (5) gives

$$6 \left| \frac{\epsilon'}{\epsilon} \right| = \left(1 - \frac{|\eta_{00}|^2}{|\eta_{+-}|^2} \right) \quad (12)$$

Fig. 6. Resolution, obtained by Monte Carlo simulation, in the vertex reconstruction of the decay $K^0 \rightarrow \pi^0 \pi^0$ as a function of the distance from the target and for different detectors.

and thus a precision in $|\epsilon'/\epsilon|$ of $\sim 2 \times 10^{-3}$ is expected for a 10-day run at LEAR.

Method 2: A determination of both the interference amplitude $|\eta_{\pi\pi}|$ and the phase $\theta_{\pi\pi}$ could be obtained by fitting the asymmetry factor $A_{\pi\pi}(t)$ [Eq. (10)] or the factor $U_{\pi\pi}(t)$ [Eq. (11)]. The determination of $|\eta_{\pi\pi}|$ and $\theta_{\pi\pi}$ is sensitive (Figs. 4 and 5) at later decay times ($\sim 14\tau_S$), where a smaller fraction of the events contributes and the determination of the vertex in the neutral decays is most precise (Fig. 6). Moreover, the width of the interference pattern is much bigger than the resolution in the neutral-vertex reconstruction.

This method is more attractive since we can define the magnitude and the phase of each CP-violation parameter. A precision of 2×10^{-3} in $|\epsilon'/\epsilon|$ could be expected for a running time of 10 days at LEAR.

Method 3: In this approach we abandon the time-dependent measurement of the neutral decays and we use only the integral rates in a time

interval, i.e. from $t = 0$ to $t \approx 20 \tau_S$. We can define an integral asymmetry $I^{00}(t_0 = 20 \tau_S)$ given by:

$$I^{00}(t_0 = 20 \tau_S) = \frac{\int_0^{t_0} R_{\pi^0 \pi^0}^0(t) dt - \int_0^{t_0} \overline{R}_{\pi^0 \pi^0}^0(t) dt}{\int_0^{t_0} R_{\pi^0 \pi^0}^0(t) dt + \int_0^{t_0} \overline{R}_{\pi^0 \pi^0}^0(t) dt} \approx 4 \operatorname{Re} \eta_{00} - 2 \operatorname{Re} \epsilon. \quad (13)$$

Then the ratio $|\epsilon'/\epsilon|$ as a function of $|\eta_{+-}|$ and I^{00} is given by the expression

$$10 \left| \frac{\epsilon'}{\epsilon} \right| = \left[2 - \frac{I^{00}}{\operatorname{Re} |\eta_{+-}|} \right]. \quad (14)$$

In a run of one day at LEAR the neutral asymmetry $I^{00}(t_0 \approx 20 \tau_S)$ can be defined to an accuracy of $\sim 5\%$, while the accuracy in $\operatorname{Re} \eta_{+-}$ is much smaller and thus does not contribute to the precision of $|\epsilon'/\epsilon|$. Consequently after 1 day of running at LEAR we will be able to measure $|\epsilon'/\epsilon|$ to an accuracy of 5×10^{-3} , whereas an increase of the running time to 10 days would allow an accuracy of 2×10^{-3} in $|\epsilon'/\epsilon|$.

Finally we summarize some advantages of this approach with initially pure K^0 and $\overline{K}^0$ states as opposed to the standard experimental configuration with K_S and K_L . The initial $\bar{p}p$ state of a precisely defined energy enables us to define the neutral kaon direction and momentum from the $K^\pm \pi^\mp$ kinematics, while the signature of the K^0 and $\overline{K}^0$ can be tagged by the sign of the charged kaon thus giving us an absolute determination of the flux. The experimental results are not strongly dependent on the efficiency for charged kaon identification, since all observables are normalized independently for K^0 and $\overline{K}^0$. In first order, systematic errors such as detector efficiencies, solid angles, small contributions from other neutral kaon decays cancel out [Eqs. (10), (11), (13)] and thus the factors $A^{\pi\pi}(t)$, $U^{\pi\pi}(t)$ as well as the integral asymmetry $I^{00}(t_0 \approx 20 \tau_S)$ are free from systematic errors.

T VIOLATION AT LEAR

According to the hypothesis of TCP invariance, the observed violation of CP invariance must be associated with a corresponding departure from T invariance^{24,25}. In spite of many searches there

is as yet no experimental evidence in support of this theoretical conclusion. At the same time, it has been shown that the observed parameters of K^0 decay cannot be easily reconciled with the hypothesis of time-reversal symmetry. Since weak interactions do not conserve strangeness, a K^0 meson can, and does occasionally, transform itself into a $\bar{K}^0$ meson in the course of time. Similarly, a state produced initially as $\bar{K}^0$ may be found later to have changed into K^0 . Time-reversal invariance, or microscopic reversibility, would require all details of the second process to be deducible from the first. From Eqs. (1) to (4) we can calculate the probability $P_{\bar{K}K}(t)$ of finding a state prepared initially as K^0 to be in a $\bar{K}^0$ state at time t , and similarly the probability $P_{K\bar{K}}(t)$ for the inverse transformation. The time-dependence of the two transition rates is found to be the same, so that their ratio is independent of time, and we can define a time-independent time-asymmetry factor

$$a_T = \frac{P_{K\bar{K}}(t) - P_{\bar{K}K}(t)}{P_{K\bar{K}}(t) + P_{\bar{K}K}(t)} \approx 2 \operatorname{Re} \langle K_L | K_S \rangle \quad (15)$$

A convenient way to make the measurement of the time asymmetry a_T is to take advantage of the now well verified $\Delta S = \Delta Q$ rule. Then $\pi^- \ell^+ \nu$ and $\pi^+ \ell^- \bar{\nu}$ decays of a neutral kaon provide indicators of the K^0 and $\bar{K}^0$ content, respectively. Consequently, our experimental philosophy of starting with pure K^0 and $\bar{K}^0$ states is ideally suited for defining the parameter a_T^{26} and for providing that CP non-invariance is accompanied by T non-invariance.

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