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WEAK INTERACTION EFFECTS IN e^+e^- ANNIHILATION WITH POLARISED BEAMS

Richard Simard
(Ph. D. thesis)

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WEAK INTERACTION EFFECTS IN e^+e^- ANNIHILATION WITH POLARISED BEAMS*

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ABSTRACT

Although the standard gauge model of weak and electromagnetic interactions based on the work of Salam and Weinberg has met with great success, there are experimental facts that will require its extension or its modification to a new gauge model; we refer here to the discovery of a heavy lepton at SLAC and to the absence of parity violation in atoms that is expected from the neutral weak current coupling to electrons. We propose 3 tests that bear on these questions.

First, we consider heavy lepton production in e^+e^- annihilation when one of the incident beams is longitudinally polarised and we examine the purely leptonic decay of this heavy lepton. An asymmetry in the inclusive angular distribution of one charged lepton (electron or muon) will be important in determining the structure of weak interactions of the heavy lepton. In fact, this angular asymmetry will easily distinguish between the cases $V - A$ and $V + A$ for the heavy lepton current.

Then, we consider the decay channel $L \to \nu_L +$ one hadron (L = heavy lepton) under the same experimental set-up, and examine the inclusive one-hadron angular distribution. Parity nonconservation in the decay of the heavy lepton will cause a conspicuous forward-backward asymmetry in the $\cos\theta$ distribution of the inclusive hadron spectrum near the high-energy end that can be distinguished easily from other sources of asymmetry. It will be easy then to discover the chirality ($V - A$ or $V + A$) of the heavy lepton current.

Finally we propose a test which will provide unambiguous and clear evidence for parity violation in e^+e^- annihilation. It consists in measuring a possible left-right asymmetry of inclusive hadron production with highly transversely polarised e^+e^- incident beams. If observed, this asymmetry will provide evidence of a parity violating neutral current coupling to electrons.

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INTRODUCTION

Probably the most significant development in theoretical particle physics in the last few years has been the gauge theories. These theories have solved many longstanding difficulties in the weak interactions. Not the least of their attractiveness is that they allow a unification of weak and electromagnetic interactions.

Already, as far back as 1957, Schwinger[1] had proposed a unification of weak and electromagnetic interactions. Later Glashow[2], Salam and Ward[3] and others attempted the same. It is true that there are many properties common to weak and electromagnetic interactions. Both obey a universality principle; all weak processes seem to occur with the same strength described by the universal Fermi coupling constant G_F and all electromagnetic processes involve the universal charge of the electron e. Both weak and electromagnetic currents are Lorentz vectors. A very important step towards unification was taken by Feynmann and Gellmann[4] with their "isotriplet vector current" hypothesis, namely that the isovector part of the hadronic electromagnetic current is the 3rd component of an isovector whose 1st and 2nd components are the strangeness-conserving polar vector hadronic weak current and its charge conjugate.

On the other hand, there are important differences too. Weak processes violate parity and charge conjugation maximally while electromagnetism conserves these quantum numbers. Furthermore, one interaction (electromagnetism) has an infinite range because of the massless character of the photon while the other interaction (weak) seems to have zero range experimentally. Up to now, the point-like current-current interaction successfully describes all observed weak processes and no propagator effect of a weak intermediate boson [the mass of such a boson being essentially the inverse of the range of the interaction] has been observed. The lower bound for the weak boson mass has been set at 20 GeV in the neutrino-inclusive reactions.

If one wants to replace the current-current weak interaction by one involving an intermediate boson W^{μ}, then one must have approximately the relation

$$G_F \approx \frac{g^2}{M_W^2}$$

between the dimensionless coupling g and the Fermi coupling G_F with M_W being the W mass. In a unified theory of weak and electromagnetic interactions, one must have $g \approx e$ so that

$$G_F \approx \frac{e^2}{M_W^2} \quad .$$

One obtains for the mass M_W

$$M_W \approx 50 \text{ GeV}$$

which is well above the experimental lower bound. What kind of theory then would have massless particles and particles with 50 times the mass of the proton in the same family? And the answer came from spontaneous symmetry breaking.

It was discovered by Goldstone[5] and elaborated by others[6] that the noninvariance of the vacuum under a continuous symmetry of the Lagrangian [this is called a spontaneously broken symmetry] implies the existence of massless scalar bosons, the so-called Goldstone bosons. But Higgs[7,8] and others showed later that when the broken symmetry is a gauge invariance, something totally different happens: the would-be massless Goldstone bosons combine with the would-be massless gauge bosons with two transverse polarisations to produce a set of massive vector bosons with three polarisations.

All the necessary ingredients being then available, Weinberg[9] and Salam[10] were able to build the first successful unified model of weak and electromagnetic interactions for the leptons. The hadrons were incorporated in 1970 by introducing a charmed quark.[11] t-Hooft's proof of the renormalisability of these nonabelian gauge theories with spontaneous broken symmetry[12] put nonabelian gauge theories at the forefront of physics, where they belong.

There exists now a multitude of gauge models for the weak and electromagnetic interactions where leptons, quarks, intermediate vector bosons proliferate freely. There exists also a nonabelian gauge theory of strong interactions which many believe to be correct. The motivation for such a theory comes from the observation of Bjorken scaling in deep inelastic lepton-hadron processes.[13] It was recognised that only an asymptotically free-field theory could explain such a behavior. Gross and Wilezek[14] and also Politzer[15] proved that non-belian gauge theories are asymptotically free. Then it was shown[16] that no quantum field theory which does not include nonabelian gauge fields can be asymptotically free.

There exist even unified gauge theories of strong, electro-magnetic and weak interactions; but these, being much more ambitious, are on less secure grounds. And, of course, it has been known for a long time that Einstein's theory of gravitation is a gauge theory. Thus all interactions seem to require the concept of gauge symmetry which has acquired a central place in theoretical physics today.

In this work, we will be interested only in weak and electro-magnetic interactions. We will now review the success of the standard model and point out the new experimental discoveries that seem to re-quire its extension or modifications. The standard model of Salam[11] and Weinberg[10] is based on the gauge group $SU(2) \otimes U(1)$. It con-tains four gauge bosons: the massless photon interacting with the

electromagnetic current, and 3 massive weak bosons, one neutral (Z^0) and two charged (W^+ and W^-) each coupling to its appropriate current. The theoretical lower bound on the mass of the charged weak bosons is 37.3 GeV, well outside the reach of present-day accelerators.

The fundamental fermions representation is as follows:

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L \quad \begin{pmatrix} u \\ d_c \end{pmatrix}_L \quad \text{and} \quad \begin{pmatrix} c \\ s_c \end{pmatrix}_L$$

$$e_R^- \qquad \mu_R^- \qquad u_R \qquad d_R \qquad s_R \qquad c_R.$$

All left-handed components (L) are isodoublets of the weak SU(2) and all right-handed components (R) are isosinglets of SU(2) except for the neutrinos ν_e and ν_μ which are messless and purely left-handed. d_c, s_c are the Cabbibo rotated quarks $d_c = d \cos \theta_c + s \sin \theta_c$ and $s_c = -d \sin \theta_c + s \cos \theta_c$. The first 3 doublets describes very well all known charged currents up to 1975. The motivation for the last doublet involving the charmed quark c was introduced first on aesthetic considerations (symmetry in lepton-quark).[17] It attained its full significance in 1970 when Glashow, Iliopoulos and Maiana [18] showed its usefulness in alleviating the problems of strangeness-changing neutral weak currents and higher orders of weak interactions; that is, it is the simplest and most elegant extension of the usual quark model (with u, d, s quarks only)

that will not violate the extremely small experimental bound on the strangeness-changing neutral weak process $K_L \to \mu^+\mu^-$, while at the same time being in agreement with the very small $K_L - K_S$ mass difference. The recent discovery of a new family of hadrons (ψ) with the subsequent discovery of what appears to be charmed mesons 2 years ago has made charm part of the standard folklore in weak interactions, although it is not as fully examined as the old quarks and leptons. According to the latest experimental results in e^+e^- annihilation,[19] the charm changing current is consistent with the above form

$$\bar{c}\,(-d \sin\theta_c + s \cos\theta_c) = \bar{c}\, s_c \ .$$

The conclusion then is that all charged weak currents are consistent with the standard model. Let us now consider the neutral weak currents first discovered at CERN in 1973.

(1) First there seems to be no off-diagonal flavor-changing neutral weak current as required in the standard model:

$\bar{s}d$ or $\bar{d}s$ currents are suppressed ;

this was the reason for introducing charm in the first place.

$\bar{c}u$ or $\bar{u}c$ currents are suppressed; absence of $D^0 - \bar{D}^0$ mixing[19], where D^0, $\bar{D}^0$ are charmed pseudoscalars.

(2) Leptonic neutral currents have been observed

$\bar{\nu}_\mu e^- \rightarrow \bar{\nu}_\mu e^-$ [20,21],

$\nu_\mu e^- \rightarrow \nu_\mu e^-$,

$\bar{\nu}_e e^- \rightarrow \bar{\nu}_e e^-$ [22].

The existence of these leptonic processes appears to be established and to have a reasonable magnitude. It is hard to draw definite quantitative conclusions from the meager data, but the results appear to be consistent with $\sin^2 \theta_W < 0.5$ where θ_W is the Weinberg angle, the only free parameter of the standard model.

(3) Deep inelastic neutral currents, inclusive reactions

$\nu + \text{Nucleons} \rightarrow \nu + \text{hadrons}$,

$\nu + \text{Nucleons} \rightarrow \nu + \text{hadrons}$.

There exist 3 measurements for these reactions[20, 23, 24]. Assuming the quark-parton model, these measurements yield an estimate of $\sin^2 \theta_W$ between 0.2 and 0.4.

(4) Elastic neutrino-proton scattering $\nu_\mu p \rightarrow \nu_\mu p$,

$\nu_\mu p \rightarrow \nu_\mu p$. [25, 26]

Although the experimental errors are large and many assumptions on form factors etc. go into the theoretical calculations, the measurements are not inconsistent with the standard model,

However, the limits on the Weinberg angle are not as restricting as in the inclusive reactions results.

In conclusion, it can be said that all neutral-currents experiments involving incident neutrinos are consistent with the standard model and point to a value of $\sin^2 \theta_W$ of 0.2 - 0.4.

However, there exist now several reasons why an extension or a modification of the standard model seems necessary. They are:

(1) CP violation: this extremely weak effect has been known for more than 10 years. Some gauge models attempting to explain it have been constructed and we will not consider it further.

(2) There are some indications for a new quark flavor besides charm.[27]

(3) The absence of parity violation in atoms:[28] the tests carried out for atoms of Bismuth were sensitive to the vector part of the hadronic weak neutral current and the axial-vector part of the neutral current coupling to electrons. The present experimental evidence is consistent with no parity violations in atoms; the reported errors allow an effect of the order of 10% of that predicted by the standard model. However, the interpretation of the atomic physics experiments could be wrong since one relies on very detailed calculations

of the atomic wave functions[29,30] but some experts consider these calculations to be reliable at least to a factor of 2.

It is thus very important to try to measure these electronic neutral weak currents in an independent reaction and we propose a test to measure parity violation in deep inelastic electron-positron annihilation in Chapter III of this work.

(4) There exists now very good evidence for the existence of a heavy lepton[31] of mass 1.8 - 2.0 GeV. The properties of this particles [lepton number, weak decay characteristics...] are very important as they will influence crucially how the standard model has to be extended. It will provide us with new insights on the ultimately correct gauge theory of weak and electromagnetic interactions. We address ourselves to the properties of this heavy lepton in Chapters I and II.

CHAPTER I: μe EVENTS IN ELECTRON-POSITRON ANNIHILATION WITH LONGITUDINALLY POLARIZED BEAMS

The μe events discovered at SPEAR[31] indicate the existence of heavy leptons. Details of their decay properties are of special interest from the viewpoint of weak-interaction theory. At the present moment little is known about their identity beyond the fact that they decay into a muon or an electron through three-body decay. The distributions of the collinearity angle and the momenta of μe are insensitive to whether the weak current of the heavy leptons is V - A or V + A.[32-35] It is rather difficult to distinguish between V - A and V + A through the collinearity angle and the momentum spectra.

Meanwhile, large transverse polarizations have been obtained at SPEAR in the vicinity of $\sqrt{s} = 7.4$ GeV. It is expected that large polarizations will be obtained more easily at PEP and PETRA.[36] Since transverse polarizations are not very useful for the purpose of exploring weak interactions or parity violation,[37] it is desired to rotate one or both of the beam polarizations into the longitudinal direction.[38] The parallel longitudinal polarizations would be most appropriate for our purpose, but it is almost equally useful if one of the beams is polarized longitudinally and the other beam remains unpolarized or transversely polarized. We examine in this chapter the angular distributions of muons and electrons in the

μe events when one of the initial beams is polarized longitudinally. With the initial beam polarized longitudinally, properties associated with parity will manifest themselves clearly, for instance, in the asymmetry of the angular distribution. This asymmetry in the emission angles of muons and electrons is large enough to distinguish easily between $V - A$ and $V + A$ for the heavy-lepton urrents.

I.1 The New Lepton Number

We must first of all discuss the lepton number assignment. The new lepton L can be either of the muon-type, or of the electron-type, or else it can be a sequential heavy lepton with its own lepton number and its own massive or massless neutrino ν_L.

(a) $L^+ \equiv M^+$ has the same lepton number as μ^-;

in this case, the heavy lepton could be produced in deep inelastic inclusive neutrino reactions and would subsequently decay rapidly into a pure leptonic final state, amongst others.

$$\nu_\mu + \text{Nucleons} \rightarrow M^+ + \text{anything}, \quad M^+ \rightarrow \mu^+ \bar{\nu}_\mu \nu_\mu \qquad \text{(I1.1)}$$

The net effect would be a "wrong-sign" muon in the final state, compared to the normal reaction

$$\nu_\mu + \text{Nucleons} \rightarrow \mu^- + \text{anything} \qquad \text{(I1.2)}$$

Barish et al.[41] have set an experimental lower bound on the mass of such a lepton. Assuming a reasonable branching ratio[32] of 30% for the decay $M^+ \to \mu^+ \nu_\mu \nu_\mu$, and a weak coupling constant G_M of M^+ equal to the usual Fermi coupling of $\mu (G_M = G_F)$, they obtain

$$\text{mass}_{M^+} > 8.4 \text{ GeV} \quad . \tag{I.1.3}$$

Conversely, assuming mass $(M^+) = 2$ GeV, which is the experimental upper bound measured at SLAC, one obtains

$$|G_M|^2 < .03 |G_F|^2 \quad . \tag{I.1.4}$$

Thus, if one assumes that there is no more than one new charged lepton of the muon type, keeping in mind the universality of the weak interactions, then the neutrino inclusive reactions (1.1) would already imply that the SLAC lepton cannot be of this type. But if one is willing to introduce more than one such lepton, then one can bypass the mass lower bound (I.1.3) set by the neutrino reactions. For instance, in a unified gauge theory of weak and electromagnetic interactions which has $SU(2)_{\text{Left-handed}}$ as a subgroup, one would have $SU(2)$ weak isospin multiplets like

$$\begin{pmatrix} M^+ \cos\theta + {}^*M^+ \sin\theta \\ \text{neutrals} \\ \mu^- \end{pmatrix}_{\text{Left-handed}}$$

where $^*M^+$ is another heavy lepton and the small mixing angle θ is restricted by (1.4) to be

$$\cos^2 \theta < .03 .$$

However it is nice to know that the SLAC results themselves exclude all such leptons independently of the neutrino reactions, as we will see in a moment.

(b) $L^+ \equiv E^+$ has the same lepton number as e^-.

Since beams of ν_e neutrinos are very hard to obtain, the inclusive neutrino reactions have nothing to say about the existence of this kind of lepton.

The most important fact about these types of leptons is, assuming the new neutral lepton ν_L emitted in the decay of L to be also massless, that the 2 neutrinos in the final state are identical for some decay channels. For instance in case (a) we have the possible final state

$$L^+ \rightarrow \mu^+ \nu_\mu \nu_L \qquad \text{(I.1.5a)}$$

where $\nu_L = \nu_\mu$

and in case (b)

$$L^+ \rightarrow e^+ \nu_e \nu_L \qquad (I.1.5b)$$

where $\nu_L = \nu_e$.

The Pauli exclusion principle forbids 2 fermions to exist in the same state which is insured by requiring the S-matrix amplitude to be antisymmetric under permutation of the 2 identical neutrinos. Consequently, not every 4 fermion coupling will be allowed in the interaction Lagrangian and in fact most coupling constants in Eq. I2.1 below will vanish identically. Some configurations will be thus forbidden and the charged lepton spectra will have some characteristic zeros. As an example, the decay amplitude of L must vanish at the high-energy end of the charged lepton energy spectrum as we can see from Fig. I. Since the 2 neutrinos, which are assumed to exist in a unique helicity state, would be in an identical momentum state.

Although the SLAC experimental data cannot discriminate yet between identical and nonidentical neutrinos on the basis of differential spectra, the measurements on branching ratios can eliminate cases (a) and (b).

For example, let's assume that L^+ has the same lepton number as e^-; then the total decay rates satisfy: [we are assuming that ν_e exists only as a left-handed ν]

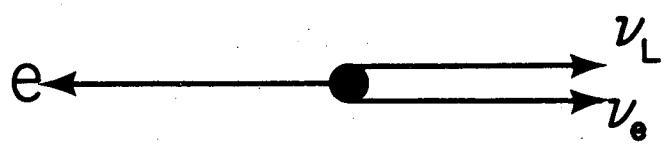

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Fig. 1

$$\Gamma(L^+ \to e^+\nu_e\nu_e) = 2\Gamma(L^+ \to \mu^+\nu_\mu\nu_e) \qquad (I.1.6)$$

Equation (I1.6) can be understood as follows: the electron amplitude is twice the muon amplitude because of the extra contribution coming from the permuted identical neutrinos, and the permuted amplitude contributes as much as the direct amplitude. Squaring the amplitude would give a factor 4 in the rates except that we must divide by 2 to avoid counting each neutrinos twice. So then, in the SLAC experiments

$$e^+e^- \longrightarrow L^+ \quad L^- \longrightarrow \ell^-\nu_\ell\nu_L \qquad (I.1.7)$$

$$L^+ \hookrightarrow \ell^+\nu_\ell\nu_L$$

(where ℓ is the lepton e or μ), theory predicts if $\nu_L = \nu_e$ from Eq. (I1.6)

$$\frac{\sigma(ee)}{\sigma(e\mu)} = 1$$

and (I.1.8)

$$\frac{\sigma(\mu\mu)}{\sigma(e\mu)} = \frac{1}{4}$$

If instead, $\nu_L = \nu_\mu$ the ratios (I1.8) would be interchanged. The experimental results[42] are:

$$\text{for } 4.8 < E_{CM} < 6.8 \text{ GeV} \qquad \frac{\sigma(ee)}{\sigma(e\mu)} = .54 \pm .3$$

$$\frac{\sigma(\mu\mu)}{\sigma(e\mu)} = .50 \pm .2$$

$$\text{for } 6.8 < E_{CM} < 7.8 \text{ GeV} \qquad \frac{\sigma(ee)}{\sigma(e\mu)} = .59 \pm .3 \tag{I.1.9}$$

$$\frac{\sigma(\mu\mu)}{\sigma(e\mu)} = .47 \pm .2$$

It is true that there are large experimental errors, but even then the cases (a) and (b) of identical neutrinos $\{\nu_L = \nu_\mu$ or $\nu_L = \nu_e$ respectively$\}$ can be ruled out with some confidence.

In all other cases (c,d,e below), the final neutrinos in the decay of L are nonidentical; consequently there are no restrictions on the 4-fermion coupling due to the Pauli exclusion principle. If there are no neutral currents contributions, then the theory predicts for these 3 last cases:

$$\Gamma(L \to e\nu_L\nu_e) = \Gamma(L \to \mu\nu_L\nu_\mu) \tag{I.1.10}$$

so we get

$$\frac{\sigma(ee)}{\sigma(e\mu)} = \frac{\sigma(\mu\mu)}{\sigma(e\mu)} = \frac{1}{2} \qquad \text{(I.1.11)}$$

This is perfectly consistent with the experimental results (I.1.9) Furthermore, when the experimental errors can be made sufficiently small, one will be able to conclude from (I.1.9)that the neutral currents contributions, if any, are small, because of the near equality of the electron and muon rations.

Assuming the conserved vector hypothesis (CVC), asymptotic chiral SU(3), the partially conserved axial current hypothesis (PCAC) and other standard lore in weak interactions, Tsai[32] has estimated the branching ratio

$$\frac{\Gamma(L \to \nu_L e \nu_e)}{\Gamma(L \to \text{all})} = 33\% \qquad \text{for identical neutrinos } \nu_L = \nu_e ,$$

$$= 20\% \qquad \text{for nonidentical neutrinos and no neutral currents contributions to this decay.}$$

The SLAC experimental results, obtained by assuming the equality of the L decay rates to electrons and muons, V - A coupling, and mass $(\nu_L) = 0$, are[43]

$$\frac{\Gamma(L^- \to e^-\nu_L\nu_e)}{\Gamma(L^- \to \text{all})} = \frac{\Gamma(L^- \to \mu^-\nu_L\nu_\mu)}{\Gamma(L^- \to \text{all})} = 0.186 \pm .030.$$

Again, this is consistent with the theoretical prediction for nonidentical neutrinos, and also seems to indicate very small neutral currents contributions.

In the rest of this work, we will assume that the neutrinos in the decay $L \to \ell\nu_\ell\nu_L$ are nonidentical, and so as we said, there are 3 possible lepton-number assignments, although the calculations in this work are independent of which case is realised in the actual world.

(c) $L^- \equiv M^-$ has the same lepton number as μ^-.

Then the simplest weak isospin multiplet in a unified gauge theory with SU(2) group is

$$\begin{pmatrix} \nu_\mu \\ \mu^- \cos\theta + M^- \sin\theta \end{pmatrix} \quad \text{Left-handed} \quad . \tag{I.1.12}$$

We would have to introduce an equal mixing angle and a new e^--type heavy lepton for the electron multiplet and also for the quarks multiplets to explain the universality of weak Fermi coupling.

C. H. Albright[44] has shown that the deep inelastic inclusive neutrino reaction

$$\nu_\mu \mu^- \rightarrow M^- + \text{anything}$$

would rule out such a lepton provided that it couples to the same left-handed neutrino ν_μ with the same strength as the muon. In terms of our model (I.1.12) of multiplet, this would mean that the mixing angle is such that

$$|\cos\theta| > |\sin\theta|$$

or

$$|\tan\theta| < 1 \quad .$$

In effect, this means that if the new lepton is of this type, it has a smaller coupling to ν_μ than the muon. An interesting fact about this type of lepton is the possibility of neutral off-diagonal currents; but we will treat these kinds of currents in the similar case of (d).

(d) $L^- \equiv E^-$ has the same lepton number as e^-.

This case has been analysed in some details by Ali and Yang[45]. The most interesting characteristic of this case is the possibility of neutral currents giving rise to the decays:

$$
\begin{aligned}
L^- &\to e^- e^- e^+ \ , \\
L^- &\to e^- \mu^- \mu^+ \ , \\
L^- &\to e^- \nu_\mu \bar{\nu}_\mu \ , \\
L^- &\to e^- + \text{hadrons} \ , \\
L^- &\to e^- \nu_e \bar{\nu}_e \ .
\end{aligned}
\qquad (I.1.13)
$$

The last reaction can also occur via charged currents. Hence, the relative rates of $L^- \to \mu^- \nu_\mu \nu_e$ and $L^- \to e^- \nu_e \nu_e$ will be different, because of the neutral currents contribution to $L^- \to$ $\to e^- \nu_e \nu_e$.

However, S. Glashow and S. Weinberg call these gauge models unnatural and unaesthetic. They much prefer the so-called natural gauge models in which all the neutral currents are diagonal with respect to particle identities. This would mean for instance a neutral current of the form:

$$\bar{L}\, \gamma_\mu \{a + b\gamma_5\}\, L$$

only, where L is the heavy lepton quantum field. In these natural models, none of the decay processes above in (II.13) would occur, except for the last reaction which would be caused only by charged currents.

The SLAC data, although involving large experimental errors, is consistent with

$$\Gamma(L \rightarrow \mu\nu\nu) = \Gamma(L \rightarrow e\nu\nu)$$

which means either that the underlying gauge theory is a natural one, or if unnatural, that the neutral currents contribution to L decay is small.

One other signature of this type of lepton in an unnatural gauge model is the 4 charged lepton production mode:

$$e^+e^- \longrightarrow L^- \quad + \quad L^+$$

$$L^- \downarrow e^-e^-e^+ \qquad L^+ \rightarrow \mu^+ + \text{neutral}$$

or effectively

$$e^+e^- \rightarrow \mu\, eee + \text{neutrals}$$

This has not been seen. It is hoped that the uncertainties in the data can be eliminated or decreased soon so that definite conclusions can be reached concerning this case.

(e) Finally, the case of the sequential heavy lepton $L^- \equiv \tau^-$, in which L has a new lepton number with its own neutrino.

In a gauge theory with $SU(2)$ group, the simplest weak isospin multiplets are

$$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \qquad \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L \qquad \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L$$

In this case, the theoretical branching ratio

$$\frac{\Gamma(L \to e\nu_e\nu_L)}{\Gamma(L \to \text{all})} = .20$$

and the ratios $\dfrac{\sigma(ee)}{\sigma(e\mu)} = \dfrac{\sigma(\mu\mu)}{\sigma(e\mu)} = \dfrac{1}{2}$ exactly.

These are perfectly consistent with the experimental numbers.

I.2 Angular distribution of μe events

The distributions of the collinearity angle and the momenta have been calculated by many people,[32-35] but they can reveal little of the structure of weak interactions. Being interested in parity-violation effects in heavy-lepton decays, we calculate here the asymmetry in the cos θ distribution of muons and electrons in the final states of the μe events (see Fig. 2). The same asymmetry in inclusive muon (electron) production might serve our purpose, but it would be largely contaminated by the muons (electrons) due to charmed hadrons by the muons (electrons) due to charmed hadrons or other heavy hadrons stable against strong and electromagnetic decays. To make sure that final light leptons originate from heavy leptons, one should select out the μe events with no accompanying hadrons.

The interactions for heavy-lepton decay is written in the Fierz-transformed four-fermion form as

$$
\begin{aligned}
L_{int} = & (\bar{\psi}_\ell \psi_L)\left[\bar{\psi}_{\nu L}(C_S + C_S'\gamma_5)\psi_{\nu\ell}\right] \\
& + (\bar{\psi}_\ell \, {}_\lambda \psi_L)\left[\bar{\psi}_{\nu L}\gamma^\lambda(C_V + C_V'\gamma_5)\psi_{\nu\ell}\right] \\
& + \frac{1}{2}(\bar{\psi}_\ell \sigma_{\lambda\,\kappa}\psi_L)\left[\bar{\psi}_{\nu L}\sigma^{\lambda\kappa}(C_T + C_T'\gamma_5)\psi_{\nu\ell}\right]
\end{aligned}
$$

(Equation I.2.1 continued on the next page)

$$+ (\bar{\psi}_\ell \gamma_\lambda \gamma_5 \psi_L)\left[\bar{\psi}_{\nu L} \gamma^\lambda \gamma_5 (C_A + C_A' \gamma_5)\psi_{\nu\ell}\right]$$

$$+ (\bar{\psi}_\ell \gamma_5 \psi_L)\left[\bar{\psi}_{\nu L} \gamma_5 (C_P + C_P' \gamma_5)\psi_{\nu\ell}\right], \tag{I.2.1}$$

where ψ_ℓ and $\psi_{\nu\ell}$ are Dirac fields of a light lepton (μ or e) and its neutrino, ψ_L and $\psi_{\nu L}$ are Dirac fields of a heavy lepton and its neutrino, and the convention of the γ matrices and the metric tensor is the one that is found in Ref. 39. As is usual,[40] we define

$$\begin{aligned}
a &= |C_S|^2 + |C_S'|^2 + |C_P|^2 + |C_P'|^2, \\
b &= |C_V|^2 + |C_V'|^2 + |C_A|^2 + |C_A'|^2, \\
c &= |C_T|^2 + |C_T'|^2 , \\
\bar{\alpha} &= |C_S|^2 + |C_S'|^2 - |C_P|^2 - |C_P'|^2, \\
\bar{\beta} &= |C_V|^2 + |C_V'|^2 - |C_A|^2 - |C_A'|^2, \\
a' &= -(C_S C_P'^* + C_S' C_P^* + C_S^* C_P' + C_S'^* C_P) , \\
b' &= C_V C_A'^* + C_V' C_A^* + C_V^* C_A' + C_V'^* C_A,
\end{aligned} \tag{I.2.2}$$

(Equation I.2.2 continued on next page)

$$c' = -(C_T C_T'^* + C_T^* C_T'),$$

$$\bar{\alpha}' = -(C_S C_P'^* + C_S' C_P^* - C_S^* C_P' - C_S'^* C_P) ,$$

$$\bar{\beta}' = C_V C_A'^* + C_V' C_A^* - C_V^* C_A' - C_V'^* C_A ,$$

and also

$$\begin{aligned} A_0 &= a + 4b + 6c, \\ \rho &= (3b + 6c)/A_0 \\ \eta &= (\bar{\alpha} - 2\bar{\beta})/A_0 \\ \xi &= (-3a' - 4b' + 14c')/A_0 \\ \delta &= (3b' - 6c')/A_0\xi . \end{aligned} \qquad (I.2.3)$$

If the heavy-lepton current is $V - A$ prior to the Fierz transform, one finds that $\rho = \delta = \frac{3}{4}$ and $\xi = 1$. If it is $V + A$, one finds that $\rho = \delta = 0$ and $\xi = 3$. In either case, $\eta = 0$. In the present paper the mass of the neutrino associated with the heavy lepton is assumed to be zero.

Let us first calculate for the process in which the electron beam is polarized longitudinally along its own momentum (as a vector) with the positron beam transversely polarized. We consider the angular distribution of light leptons of negative charge in the final states of such a process (Fig. 3). The other cases are obtained trivially from this case. Only the one-photon annihilation is taken into account. In computing the differential cross section we follow the method of Tsai[32] modified to a manifestly covariant form.

The L^+L^- production cross section with L^+L^- polarized arbitrarily is given in the e^+e^- laboratory frame by

$$\frac{d\sigma}{d\Omega_L} = \frac{\alpha^2\beta}{(q^2)^3} W \qquad (I.2.4)$$

$$W = W_1 + \mathcal{P} W_2 + O(m_\ell^2) , \qquad (I.2.5)$$

$$\begin{aligned} W_1 = {} & 2\Big[(P \cdot k)(P' \cdot k') + (P \cdot k')(P' \cdot k) + \tfrac{1}{2} q^2 M^2\Big]\Big[1 - (s \cdot s') - \\ & -q^2(s \cdot P')(s' \cdot P) + \tfrac{1}{2}(q^2)^2(s \cdot s') - q^2 (s \cdot k)(s' \cdot k') + \\ & (s \cdot k')s' \cdot k)\Big] + 2(s \cdot P') \ (P \cdot k)(s' \cdot k') + (P \cdot k')(s' \cdot k)\Big] \\ & + 2(s' \cdot P)\Big[(P' \cdot k)(s \cdot k') + (P' \cdot k')(s \cdot k)\Big], \qquad (I.2.6) \end{aligned}$$

$$W_2 = -Mq^2\Big[(s \cdot k) + (s' \cdot k) - (s \cdot k') - (s' \cdot k')\Big] \qquad (I.2.7)$$

where

$$\begin{aligned}
\rho &= \text{magnitude of } e^+e^- \text{ beam polarizations,} \\
s_\mu(s'_\mu) &= \text{four-vector polarization of } L^-(L^+), \\
M &= \text{mass of } L^\pm \\
P_\mu(P'_\mu) &= \text{four-momentum of } L^-(L^+), \\
k_\mu(k'_\mu) &= \text{four-momentum of initial } e^-(e^+), \\
\beta &= |\vec{P}|/P_0, \text{ and } q = k + k'.
\end{aligned} \tag{I.2.8}$$

The differential decay rate of $L^\pm$ is given by

$$P_0 \frac{d\Gamma^-}{d^3p} = A^- + (s \cdot p)B^- \tag{I.2.9}$$

for L^- and by a similar expression for L^+ with A^- and B^- replaced by A^+ and B^+. Here p is the four-momentum of the light lepton. In the approximation of ignoring the square of the light-lepton mass, the covariant decay amplitudes $A^\pm$ and $B^\pm$ are given by

$$A^\pm = \frac{A_0 M}{96\pi^4 P_0} y \left[3\left(\frac{M}{2} - \frac{y}{M}\right) + 2\rho\left(\frac{4y}{3M} - \frac{M}{2}\right) + \frac{3mM}{y}\eta\left(\frac{M}{2} - \frac{y}{M}\right)\right], \tag{I.2.10}$$

$$B^{\pm} = \mp \frac{A_0 M^2}{96\,\pi^4 P_0}\,\xi\left[\left(\frac{M}{2} - \frac{y}{M}\right) + 2\delta\left(\frac{4y}{3M} - \frac{M}{2}\right)\right],$$

where $y = (p \cdot P)$.

The final phase volume is fully integrated over L^+, so that the s'-dependent terms in (2.4)-(2.7) do not contribute to the differential production cross section $p_o(d\sigma/d^3p)$. The spin dependence of L^- in $d\sigma/d\Omega_L$ and $p_o(d\Gamma^-/d^3p)$ is correctly included by putting them together as

$$p_o \frac{d\sigma}{d\Omega_L\, d^3p} = \frac{\alpha^2\beta r^2}{(q^2)^3\Gamma}\left[\bar{W}_1 A^- - (\bar{p} \cdot \bar{W}_2)B^-\right], \qquad \text{(I.2.12)}$$

where $\bar{W}_1$ and $\bar{W}_2^{\mu}$ are defined as

$$\frac{d\sigma}{d\Omega_L} = \frac{\alpha^2\beta}{(q^2)^3}\left[\bar{W}_1 + (s_\mu \bar{W}_2^{\mu}) + (s'_\mu \bar{W}_3^{\mu}) + (s_\mu s'_\nu \bar{W}_4^{\mu\nu})\right]$$

where r is the branching ratio of $L^{\pm}$ into $\ell^{\pm} + \nu_\ell(\bar{\nu}_\ell) + \nu_L(\nu_L)$, Γ is the decay rate in the e^+e^- laboratory frame of $L \to \ell + \nu_\ell + \nu_L$, integrated over the entire phase volume, and $\tilde{p}$ is defined by

$$\tilde{p}_\mu = p_\mu - P_\mu(p \cdot P)/M^2. \qquad \text{(I.2.14)}$$

Equation (2.12) is integrated over $d\Omega_L$ with $\vec{p}$ kept fixed. Measuring the solid angle of $\vec{P}$ with respect to $\vec{p}$, one can integrate over the azimuthal angle Φ. The polar angle Θ should be replaced by a new variable $y = (p \cdot P) = pP_0(1 - \beta \cos \Theta) + O(m^2)$. The region of integration over y is restricted by $|\cos \Theta| \leqslant 1$ and $0 \leqslant (P - p)^2 \leqslant (M - m)^2$ to

$$\mathrm{Max}\left(mM, \frac{xM^2}{1+\beta}\right) \leqslant y \leqslant \mathrm{Min}\left(\frac{1}{2}M^2, \frac{xM^2}{1-\beta}\right) \qquad (I.2.15)$$

up to $O(m^2)$. We thus obtain

$$\frac{d\sigma}{dx\, d\cos\theta} = \frac{16\alpha^2\pi r^2}{q^2(1 + 4m\eta/M)}\Big[(1 + \cos^2\theta)G_1(x) + (\tfrac{3}{2}\cos^2\theta - \tfrac{1}{2})G_2(x) - 2\xi\rho\cos\theta\, G_3(x)\Big] \qquad (I.2.16)$$

where $x = |\vec{p}|/k_0$, $\gamma = P_0/M$, θ is the emission angle of the light lepton ℓ^- measured from the direction of the electron beam momentum $\vec{k}$, and

$$G_1(x) = \left(1 + \frac{1}{2\gamma^2}\right)\Big\{\Big[\frac{3}{16} f_2(x) - \frac{1}{8} f_3(x)\Big] + \rho\Big[-\frac{1}{8} f_2(x) + \frac{1}{9} f_3(x)\Big] + \frac{m}{M}\eta\Big[\frac{3}{4} f_1(x) - \frac{3}{8} f_2(x)\Big]\Big\} \qquad (I.2.17)$$

$$G_2(x) = \frac{1}{x\gamma^2}\left\{\left[-\frac{1}{8}f_3(x) + \frac{3}{32}f_4(x)\right] + \frac{3}{x\gamma^2}\left[\frac{1}{128}f_4(x) - \frac{1}{160}F_5(x)\right] + \right.$$

$$+ \rho\left[\frac{1}{12}f_3(x) - \frac{1}{12}f_4(x)\right] + \frac{\rho}{x\gamma^2}\left[-\frac{1}{64}f_4(x) + \frac{1}{60}f_5(x)\right] +$$

$$\left. + \frac{m\eta}{M}\left[-\frac{3}{8}f_2(x) + \frac{1}{4}f_3(x)\right] + \frac{m\eta}{Mx\gamma^2}\left[\frac{1}{16}f_3(x) - \frac{3}{64}f_4(x)\right]\right\}, \tag{I.2.18}$$

$$G_3(x) = x\left[\frac{1}{4}f_1(x) - \frac{1}{8}f_2(x)\right] + \delta x\left[-\frac{1}{2}f_1(x) + \frac{1}{3}f_2(x)\right] +$$

$$+ \left[-\frac{1}{16}f_2(x) + \frac{1}{24}f_3(x)\right] + \frac{1}{x\gamma^2}\left[\frac{1}{48}f_3(x) - \frac{1}{64}f_4(x)\right] +$$

$$+ \delta\left[\frac{1}{8}f_2(x) - \frac{1}{9}f_3(x)\right] + \frac{\delta}{x\gamma^2}\left[-\frac{1}{24}f_3(x) + \frac{1}{24}f_4(x)\right], \tag{I.2.19}$$

$$f_n(x) = \left(\frac{2}{M^2}\right)^n \left[y_{max}{}^n(x) - y_{min}{}^n(x)\right]. \tag{I.2.20}$$

In Eq. (2.20), $y_{max}(x)$ and $y_{min}(x)$ stand for the upper and lower limits of integration as given in (2.15).

We write the production cross section of light leptons in the form of

$$\frac{d\sigma}{dx\, d\cos\theta} = r^2\sigma_{tot}(e^+e^- \to \mu^+\mu^-)\; R(x;\cos^2\theta)$$

$$\left[1 + \alpha(x;\cos^2\theta)\mathcal{P}\cos\theta\right].$$

The integral of $R(x;\cos^2\theta)$ in $\cos\theta$ and x over the entire physical region tends to unity as $\gamma \to \infty$, thus leading to $\sigma_{tot}(e^+e^- \to L^+L^-) \to \sigma_{tot}(e^+e^- \to \mu^+\mu^-)$.

If one measures the emission angle of ℓ^+ instead of ℓ^- in the process under consideration, one should replace the polarization $\mathcal{P}$ by $-\mathcal{P}$. The same switch of sign should be done when ℓ^- is measured in the annihilation of a transversely polarized electron beam with a positron beam polarized longitudinally along the positron beam momentum (opposite to the electron beam momentum). The formula (I.2.21) is valid in the case that ℓ^+ is measured in the annihilation of a transversely polarized electron beam and a positron beam polarized along its own momentum.

I.3 Numerical Estimate

A numerical estimate has been made for the magnitude of the cross section in terms of $R(s;\cos^2\theta)$ and for the asymmetry $\alpha(x;\cos^2\theta)$. Both of these quantities are dependent on the velocity of L. We have plotted in Figs. 4 and 5 the results for the $V - A$ and $V + A$ currents provided that the $(\ell\nu_1)$ current is always of $V - A$. For $M \gtrsim 2.0$ GeV, $R(x;\cos^2\theta)$ and $\alpha(x;\cos^2\theta)$ are insensitive to the light-lepton mass even near $x = 0$, so we have drawn the curves for $m = m_e \simeq 0$. General behaviors of R and α are rather insensitive to values of γ as long as γ is larger than 3 or so. Unlike the collinearity angle or the momentum

spectra in the annihilation of unpolarized or transversely polarized beams, the asymmetry α will distinguish $V - A$ and $V + A$ clearly. The magnitude of the asymmetry tends to unity at $\cos\theta = \pm 1$ and $x = 1$ as $\gamma \to \infty$, as is expected from helicity conservation. The magnitude of the cross section approaches zero as $x \to 1$ since $L^{\pm}$ decay into three bodies. To obtain the maximum efficiency, one should choose the region

$$0.4 \lesssim x \lesssim 0.6 \tag{I.3.1}$$

or possibly

$$x \lesssim 0.05, \tag{I.3.2}$$

and stay in the forward-backward cones of, say

$$|\cos\theta| \lesssim 0.5. \tag{I.3.3}$$

In Fig. 4 are plotted the cross section in terms of quantity $R(x; \cos^2\theta)$ and the asymmetry $\alpha(x; \cos^2\theta)$, both defined in (I.2.21), as functions of $x = |\vec{p}|/k_0$ at $\gamma = 2$, 6, and 10. The $\cos\theta$ dependence of α has been shown in Fig. 5. Finally, fixing values of x at 0.5 and 0.05, we have plotted the γ dependence of α between $\gamma = 2$ and $\gamma = 10$ in Fig. 6.

I.4 Comment

Semileptonic decays of charmed hadrons or other heavy hadrons stable agains strong and electromagnetic interactions can produce μe events, but the criterion of selecting the μe events with no accompanying hadron, charged or neutral (except for K_L^0), in final states eliminates all but $e^+e^- \to D^+D^-(F^+F^-) \to \mu + e + \nu_\mu + \nu_e$. If stable D^+ and F^+ are of $J^P = 0^-$, as the ψ spectroscopy suggests, the branching ratio into $(\ell\nu_l)$ is negligibly small in the $V \pm A$ interactions. If D^+ and F^+ of $J^P = 1^-$ should decay through weak interactions, one might have to separate them out. But strong damping of form factors still suppress two-body hadronic channels so severely that we would hardly worry about this possibility. Because of this we have proposed here to examine only the μe events rather than inclusive μ (or e) production.

CHAPTER II: A TEST FOR THE CHIRALITY OF HEAVY LEPTON CURRENT IN INCLUSIVE PION PRODUCTION

In the previous chapter, we have used the weak leptonic current of the ordinary leptons (electrons and muons) to probe, so to speak, the heavy lepton current or its chirality. In this chapter, we will use instead the standard weak hadronic current in its V - A form as a probe.

II.1 Beams Polarisation

The test will make use once again of longitudinal polarisation of the incident beams. We have learned recently that large beam polarisations will be much harder to obtain than was anticipated. Chao and Schwitters[49] have shown that a complete treatment of depolarisation effects including quantum fluctuations leads generally to very small polarisations. One has to choose carefully beam energies and duration of beam storage in order to achieve large polarisations. Nevertheless, hoping that experimentalists will attempt this difficult task and obtain large longitudinal polarisations at the PEP-colliding beam machine, we present here a clear test to determine the chiral property of the weak heavy-lepton current. It consists of measuring charged pions (and kaons) of very high energy produced through $L \to \nu_L + \pi(K)$.

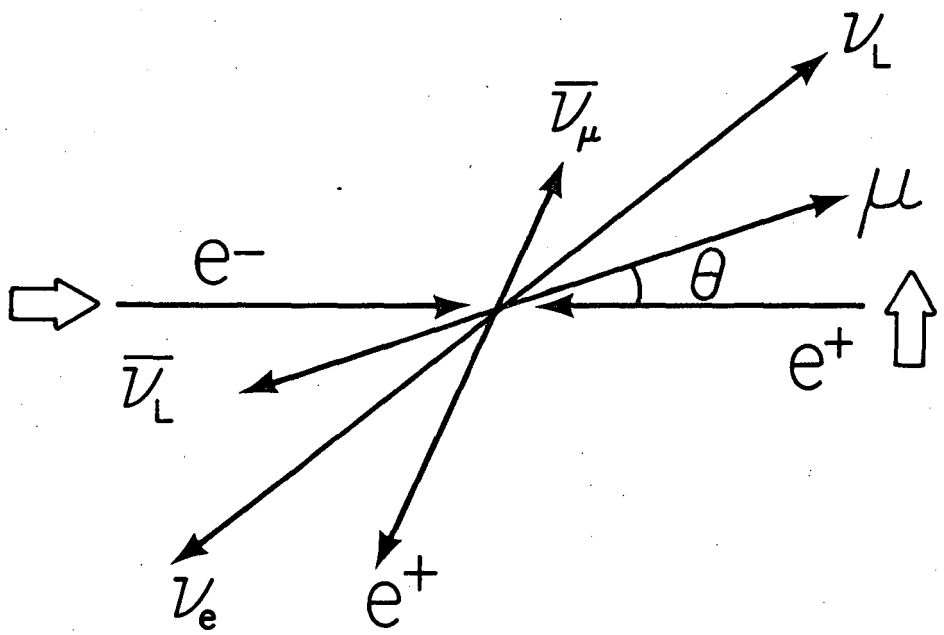

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Fig. 2

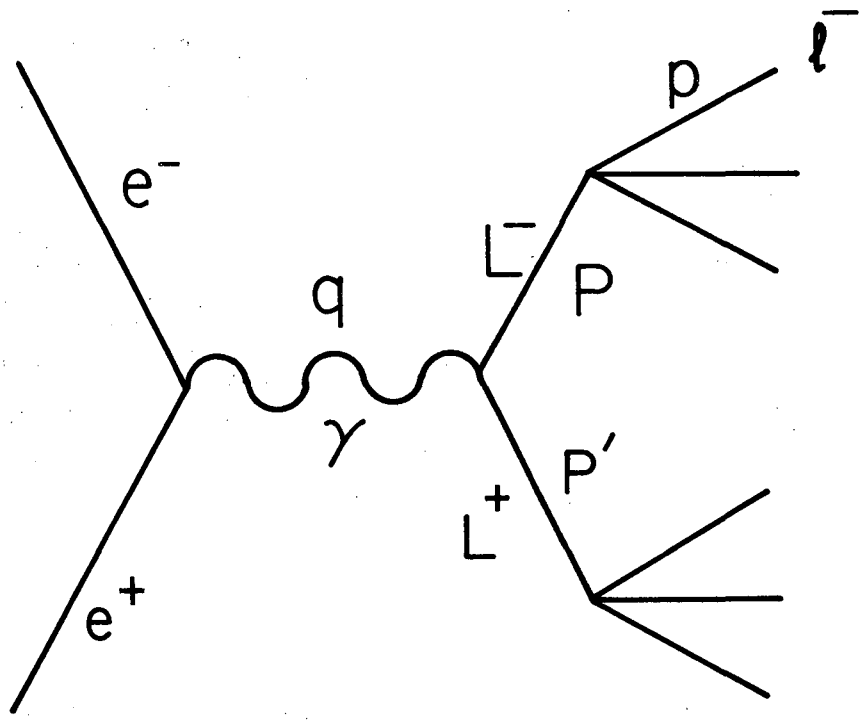

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Fig. 3

II.2 Angular Distribution

The angular distribution of pions in the one-photon annihilation

$$e^- + e^+ \xrightarrow[\text{"}\gamma\text{"}]{} L^- + L^+ \qquad \text{(II.2.1)}$$

$$L^+ \searrow \text{anything}$$

$$L^- \downarrow \nu_L + \pi^-$$

and a similar process is calculated in the center-of-mass frame of e^+e^- when one of the e^+e^- beams is longitudinally polarised. The reasons why this particular decay mode is chosen are that the decay matrix element is calculable accurately when universality holds, and that because the decay is a two-body decay the pion spectrum extends to the high-energy end of the x distribution. Let the semileptonic weak interaction of L be

$$H_{int} = \left(\frac{1}{2}\right)^{\frac{1}{2}} G\left[\nu_L \gamma_\mu (1 + g_a \gamma_5) L\right] J^\mu + \text{H.c.}, \qquad \text{(II.2.2)}$$

$$J^\mu = \bar{d}\, \gamma^\mu (1 - \gamma_5) u \cos\theta_c + \bar{s}\, \gamma^\mu (1 - \gamma_5) u \sin\theta_c, \qquad \text{(II.2.3)}$$

where θ_c is the Cabibbo angle. Above several GeV in the center-of-mass energy, electron-positron beams are highly polarised transversely, and it is possible to rotate one or both of the polarisations

into the longitudinal direction at the PEP and the PETRA.[38] For the purpose of producing an asymmetry in the angular distribution we consider first the configuration in which the electron beam is polarised longitudinally along the direction of its owm momentum and the positron beam remains polarised transversely. The differential cross section of the pions emerging from the $L^- \to \nu_L + \pi^-$ decay is given in the center-of-mass frame of e^+e^- as

$$\left(\frac{d\sigma}{dx\, d\cos\theta}\right)(e^-;\pi^-) = \frac{\pi\alpha^2}{2s} r \left[1 + \cos^2\theta - \frac{4g_a\mathcal{P}}{1 + g_a^2}(2x - 1)\cos\theta\right], \tag{II.2.4}$$

where s is the c.m. energy squared; x, defined by $2E_\pi/\sqrt{s}$, takes values in the region of

$$\frac{1}{2}\left[1 - (1 - 4m_L^2/s)^{\frac{1}{2}}\right] \leq x \leq \frac{1}{2}\left[1 + (1 - 4m_L^2/s)^{\frac{1}{2}}\right] ;$$

r is the branching ratio $\Gamma(L^- \to \nu_L + \pi^-)/\Gamma(L^- \to \text{all})$; and $\mathcal{P}$ is the magnitude of the beam polarisation. In Eq. (II.2.4) we have ignored m_N^2/s and m_L^2/s (≤ 0.01 for $m_L = 1.8$ GeV above $\sqrt{s} = 20$ GeV). The bracket $(e^-;\pi^-)$ in the left-hand side indicates that e^- is polarised longitudinally and π^- is measured inclusively. By the same calculation, we find

$$\left(\frac{d\sigma}{dx\ d\cos}\right)(e^+;\pi^-) = \frac{\pi\alpha^2}{2s}\ r\left[1 + \cos^2\theta + \frac{4g_a\mathcal{P}}{1 + g_a^2}(2x - 1)\cos\theta\right]$$

(II.2.5)

when e^+ is longitudinally polarized. Here the longitudinal polarization of e^+ is along the direction of its own momentum (opposite to that of e^-). For the processes in which positively charged pions from the two-body L^+ decay are measured inclusively, we obtain through CP invariance

$$\left(\frac{d\sigma}{dx\ d\cos\theta}\right)(e^-;\pi^+) = \left(\frac{d\sigma}{dx\ d\cos\theta}\right)(e^+;\pi^-)$$

$$\left(\frac{d}{dx\ d\cos\theta}\right)(e^+;\pi^-) = \left(\frac{d}{dx\ d\cos\theta}\right)(e^-;\pi^-)$$

where the angle θ for π^+ is measured from the direction of the e^- beam momentum just as for the π^- emission angle. It is our purpose to investigate the terms linear in $\cos\theta$ in Eqs. (II.2.4) and (II.2.5).

II.3 Numerical Estimate

From $\sqrt{s} = 3.8$ to 7.4 GeV, $d\sigma/dx$ scales reasonably well above $x \doteq 0.5$.[46] By assuming the scaling, therefore, we can estimate the hadrons near the high-x end at $\sqrt{s} \geq 20$ GeV that do

not come from the heavy leptons. The asymmetry to be measured is[47]

$$A(e^-;\pi^-) = \frac{1}{2}\left[s \frac{d\sigma}{dx\, d\cos\theta} (e^-;\pi^-)_{0^\circ} - s \frac{d\sigma}{dx\, d\cos\theta} (e^-;\pi^-)_{180^\circ} \right]$$

$$= -\pi\alpha^2 r \mathcal{P} \frac{2g_a}{1 + g_a^2} (2x - 1). \qquad \text{(II.3.1)}$$

This asymmetry may be detected in three other measurements: $A(e^-;\pi^+)$, $A(e^+;\pi^-)$, and $A(e^+;\pi^+)$. One can enhance the effect by taking the combination

$$A = \left[A(e^-;\pi^-) + A(e^+;\pi^+)\right] - \left[A(e^-;\pi^+) + A(e^+;\pi^-)\right] . \qquad \text{(II.3.2)}$$

This quantity A should be compared with $s(d\sigma/ds\, d\cos\theta)$ at $\theta = 0^\circ$ and 180° of the pions that are emitted symmetrically in forward and backward directions. Let us call those pions background pions. By extrapolating the data from lower energies by scaling, we find that in $0.9 < x < 1.0$

$$s \left(\frac{d\sigma}{dx\, d\cos\theta}\right)_{\text{background},\ 0^\circ} \simeq 50 \text{ nb GeV}^2 \qquad \text{(II.3.3)}$$

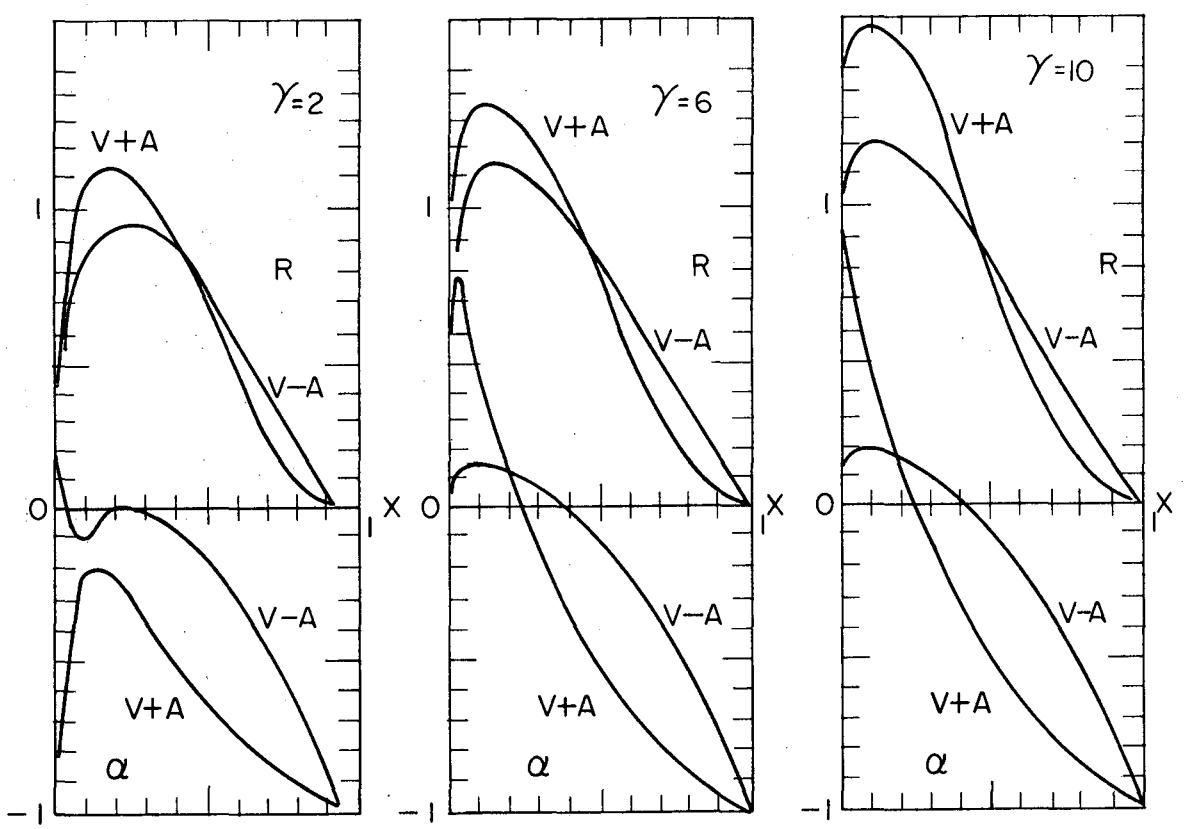

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Fig. 4

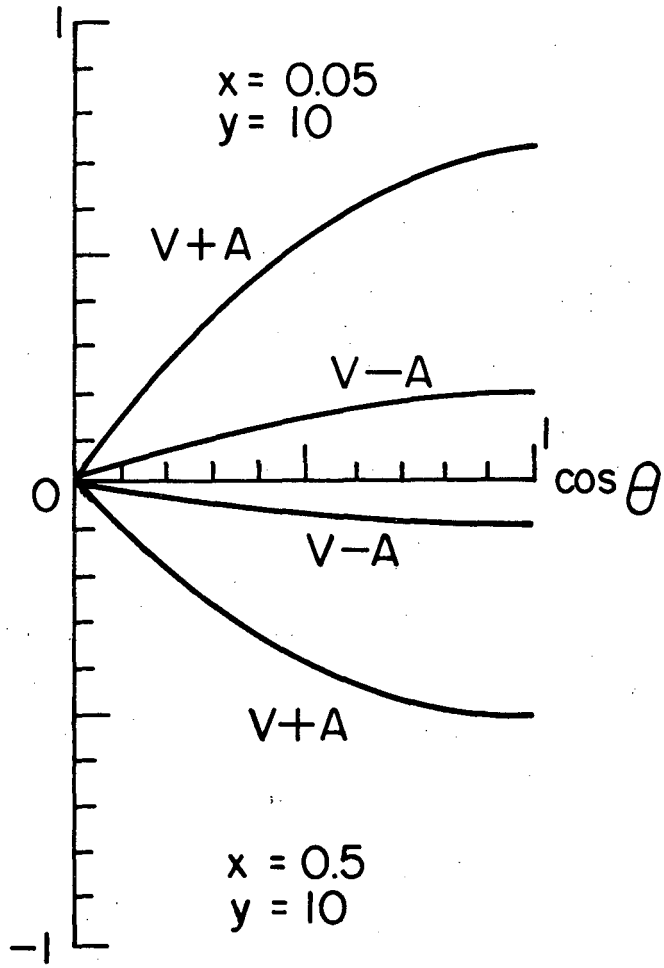

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Fig. 5

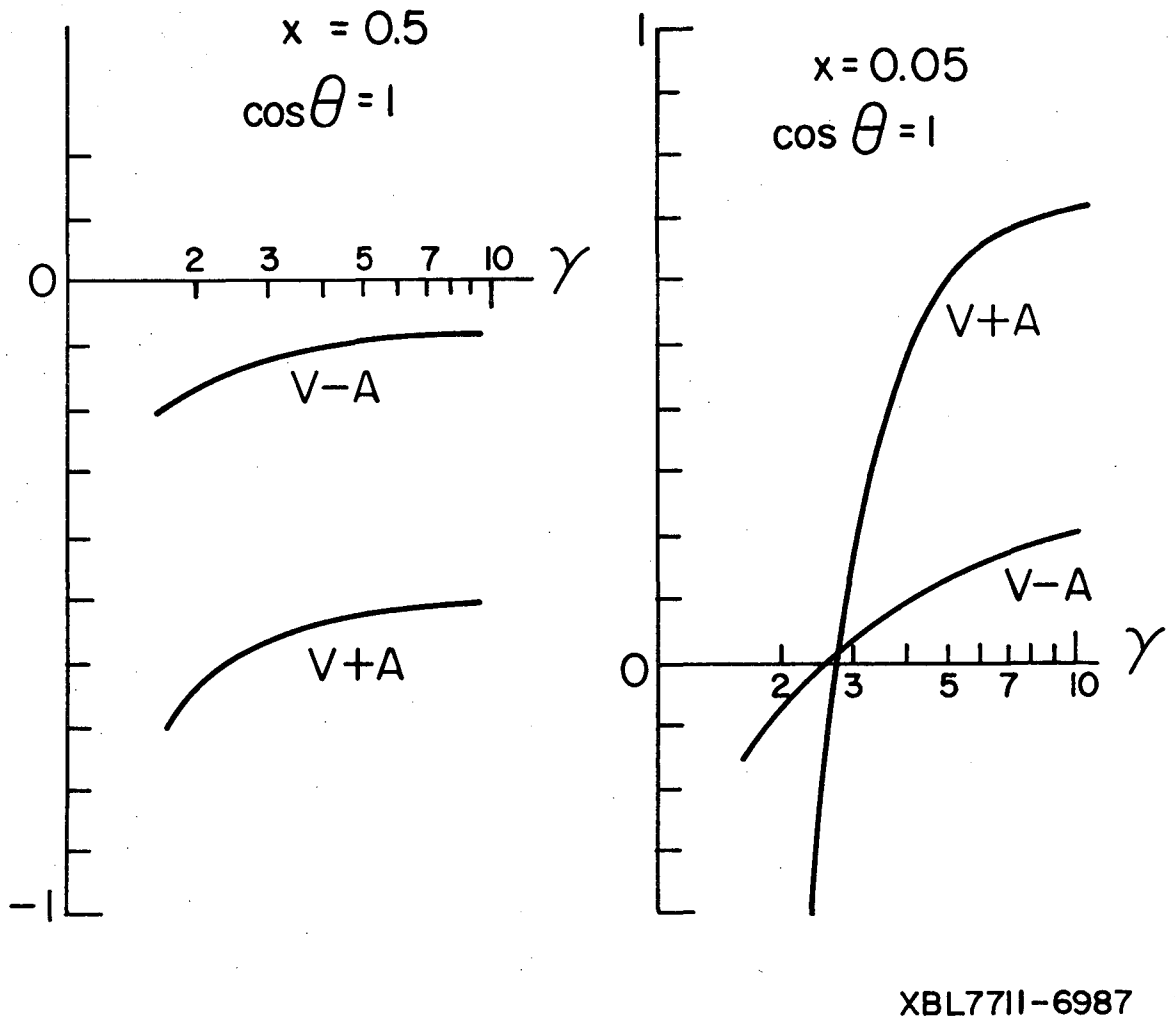
x = 0.5
cos θ = 1
0
2
3
5
7
10
γ
V−A
V+A
−1
1
x = 0.05
cos θ = 1
V+A
V−A
0
2
3
5
7
10
γ
XBL7711-6987

Fig. 6

for the sum of π^- and π^+ (and possibly $K^\pm$), provided that the $\cos\theta$ distribution is $\propto 1 + \cos^2\theta$ for the background pions.[46] On the other hand, with the decay branching ratio $r = 0.10$,[22] $g_a = \pm 1$, namely $V \pm A$, and $\mathcal{P} = 0.90$, the asymmetry in cross sections turns out to be

$$\begin{aligned} A &= \left[A(e^-;\pi^-) + A(e^+;\pi^+)\right] - \left[A(e^-;\pi^+) + A(e^+;\pi^-)\right] \\ &= \mp 2.6(2x - 1)r \times 10^2 \text{ nb GeV}^2. \end{aligned} \tag{II.3.4}$$

In the region of $0.9 < x < 1.0$, this is as large as

$$A = \mp 21 \text{ nb GeV}^2. \tag{II.3.5}$$

The asymmetry of this size is easily distinguished from the asymmetry of other origins, which we will discuss in the following section.

II.4 Asymmetry of other origins

If there are charmed hadrons that are stable against strong and electromagnetic interactions, they may emit pions through weak decay to produce the asymmetry. However, if such stable charmed hadrons are pseudoscalar mesons, simple kinematics shows that there will be no $\cos\theta$ asymmetry in one-particle inclusive spectra. When there are such hadrons of nonzero spin, they can, in principle,

cause an asymmetry. But strong damping of hadronic form factors suppresses severely two-particle channels, and when charmed hadrons are produced together with other hadrons, the chance that a secondary decay pion reaches the high-x end of the spectrum is negligibly small because the x distribution falls of very fast towards $x = 1$ in the case of multibody states. Therefore, we can safely ignore the effect due to decays of metastable hadrons.

The other source of the asymmetry is the neutral-W-boson process that competes with the one-photon annihilation. It depends on specific models of weak interactions. If we choose, for instance, the four-quark model of Salam and Weinberg, we find the asymmetry

$$\frac{A(e^-;\pi^-)}{s(d\sigma/dx\, d\cos\theta)_{\text{background},\,0^\circ}} = 7.7 \times 10^{-5} \frac{s/m_N^2}{1 - (s/M_W^2)\cos^2\theta_W} \times \left[1 + (4\sin^2\theta_W - 1)\mathcal{P}\right]\rho_- , \tag{II.4.1}$$

$$\rho_- = \nu\overline{W}_3^{(ew)}(\pi^-)/m_\pi\overline{W}_1^{(e)}(\pi^-) , \tag{II.4.2}$$

where θ_W is Weinberg's angle, M_W is the charged-W-boson mass, $\overline{W}_1^{(e)}(\pi^-)$ is the inclusive structure function $\overline{W}_1$ for π^- with the electromagnetic currents, and $\overline{W}_3^{(ew)}(\pi^-)$ is the inclusive structure

function $\overline{W}_3$ for π^- with one electromagnetic current and one weak current. Similarly,

$$\frac{A(e^-;\pi^+)}{s(d\sigma/dx\, d\cos\theta)_{\text{background},\,0^\circ}} \simeq 7.7 \times 10^{-5} \frac{s/m_N^2}{1 - (s/M_W^2)\cos^2\theta_W} \times \{1 + (4\sin^2\theta_W - 1)\mathcal{P}\}\rho_+, \tag{II.4.3}$$

with ρ_+ being a ratio of structure functions for π^+ in parallel to Eq. (II.4.2). Charge-conjugation invariance of strong interactions requires that $\rho_- = -\rho_+$. Further, CP invariance of overall processes leads us to

$$A(e^+;\pi^+) = -A(e^-;\pi^-)_{\mathcal{P}\to-\mathcal{P}} \tag{II.4.4}$$

$$A(e^+;\pi^-) = -A(e^-;\pi^+)_{\mathcal{P}\to-\mathcal{P}} \tag{II.4.5}$$

for processes in which the e^+e^- pair annihilates through one photon or one neutral W boson. On the right-hand sides of Eqs. (II.4.4) and II.4.5), the subscripts indicate that the polarization $\mathcal{P}$ is to be reversed in sign. If one estimates ρ_- and ρ_+ in the simple parton model, an upper bound is set on $|\rho_- - \rho_+|$;

$$|\rho_- - \rho_+| = \frac{6}{5} \frac{|u(\pi^+) - u(\pi^-)|}{u(\pi^+) + d(\pi^+) + s(\pi^+) + c(\pi^+)} ,$$
$$\leq \frac{6}{5} . \qquad \text{(II.4.6)}$$

Because of Eqs. (II.4.4) and (II.4.5), the $\mathcal{P}$-independent terms cancel each other in the sums $A(e^-;\pi^-) + A(e^+;\pi^+)$ and $A(e^-;\pi^+) + A(e^+;\pi^-)$. A net contribution to

$$A = A(e^-;\pi^-) + A(e^+;\pi^+) - A(e^-;\pi^+) + A(e^+;\pi^-)$$

is proportional to $4 \sin^2 \theta_W - 1$, which is a small number. One thus obtains with $\sin^2 \theta = 0.3$ the upper bound on the asymmetry due to the neutral W boson as

$$\left| \frac{A}{s(d\sigma/dx\, d\cos\theta)_{\text{background},\, 0^\circ}} \right| < 3.5 \times 10^{-2} \qquad \text{(II.4.7)}$$

at $\sqrt{s} = 20$ GeV for $M_W = 50$ GeV and $\mathcal{P} = 0.90$. This is much smaller than $\sim 40\%$, the value that is expected for the asymmetry due to pions from the heavy leptons. To make sure that the neutral-W-boson annihilation is really small, one should measure the asymmetry at different energies to see its energy dependence. The asymmetry due to the heavy leptons scales precisely, while the asymmetry due to the neutral W boson increases linearly in s or even faster if the neutral-W-boson mass is not too heavy.

Finally, one might suspect that the interference between one-photon and two-photon processes can cause the asymmetry in cos θ. But even at highest PEP and PETRA energies this is expected to be no more than several percent.[48] Moreover, charge-conjugation invariance of overall processes requires that $A(e^-;\pi^-) = -A(e^+;\pi^+)$ and $A(e^-;\pi^+) = -A(e^+;\pi^-)$, and therefore no effect remains in the combination

$$A = \left[A(e^-;\pi^-) + A(e^+;\pi^+)\right] - \left[A(e^-;\pi^+) \text{ A } A(e^+;\pi^-)\right].$$

II.5 Comment

With longitudinally polarized e^+e^- beams, one can test clearly whether the weak current of the heavy leptons is V - A or V + A. The asymmetry in cos θ of the inclusive pion production is quite large near the high-energy end of the x distribution, and its sign depends on whether the current is V - A or V + A. Other sources of the asymmetry that may compete turn out to be either sufficiently small or, if not, unambiguously separable by the center-of-mass energy dependence. We have assumed in this chapter that the heavy-lepton mass is about 1.8 GeV and that the decay obeys the universality of weak interactions. The branching ratio $\Gamma(L \to \nu_L \pi)/\Gamma(L \to \text{all})$ is uniquely determined for such leptons. If there are heavier leptons, say, of mass $\sim$ 6 GeV, they would not

contribute significantly to the asymmetry near $x = 1$, since their branching ratio into the two-body ($\nu\pi$) channel is negligibly small ($\sim 1.3\%$).[32] However, even universality should not be taken for granted.

CHAPTER III: RIGHT-LEFT ASYMMETRY IN HADRON PRODUCTION

In this chapter, we leave the domain of heavy lepton physics and consider the neutral weak current and, presumably, if one believes in gauge theories of the weak interactions, its carrier, the neutral weak intermediate vector boson Z. Ever since the first discovery of neutral currents at CERN in 1973, in neutrino-indiced reactions, one has had to take these weak neutral currents seriously. We will be concerned here, not with the neutrino or antineutrino-indiced currents but with the elctronic weak neutral currents in e^+e^- annihilation.

Various tests of weak interactions in e^+e^- annihilation have been proposed[50], and one of the few promising tests is to measure a forward-backward asymmetry in $\cos\theta$ distribution of final hadrons. Unlike the $e^+e^- \to \mu^+\mu^-$ process, strong interaction dynamics prevents us from estimating accurately hadrons due to two-photon-annihilation (without e^+e^- in final states) interfering with one-photon-annihilation. We propose here an alternative method to detect an interference of weak and electromagnetic interactions by making use of high transverse polarisation of e^+e^- beams.

III.1 Angular Distribution

Let us write a weak neutral current interaction as

$$L^{z}_{int} = g_{\ell}\bar{e}\,\gamma_{\mu}(1 + g_a\gamma_5)e\,Z^{\mu} - g_h J^{(z)}_{\mu} Z^{\mu} \qquad \text{(III.1.1)}$$

where $J^{(z)}_{\mu}$ is a hadronic weak neutral current, and g_{ℓ}, g_a, g_h are model-dependent weak coupling constants. We consider only the lowest order contribution of one photon-one Z-boson exchange (Fig. 7). With the usual electromagnetic interaction

$$L^{e}_{int} = e\,\bar{e}\,\gamma_{\mu} e\,A^{\mu} - eJ^{(em)}_{\mu} A^{\mu} \qquad \text{(III.1.2)}$$

the Lorentz-invariant S-matrix element will be:

$$\mathcal{M} = -i\frac{e^2}{s}\,\bar{v}\,\gamma_{\mu}\,u\,\langle\, nh|J^{\mu(em)}|0\,\rangle + \frac{ig_{\ell}g_h}{(s - M_z^2)}\,\bar{v}\,\gamma_{\mu}\,\cdot$$

$$\cdot\,(1 + g_a\gamma_5)\,u\,\langle\, hn|J^{\mu(z)}|0\,\rangle \qquad \text{(III.1.3)}$$

where v is the position spinor,

u is the electron spinor,

R is the detected hadron,

n are all other hadrons which are summed over.

For $s \ll M_z^2$, the dominant contribution in the cross-section will be the square of the electromagnetic amplitude; the asymmetry, that is the subject of this chapter, will come from the interference between the electromagnetic and weak neutral amplitude and hopefully will be big enough to be observable. The square of the weak amplitude will be negligible and will be discarded in this calculation. The cross-section is calculated to be

$$E_p \frac{d\sigma}{d^3p} = \frac{1}{s^2} \frac{e^4}{16\pi^2} \left\{ W_1 + W_2 \frac{\vec{p}^2}{M^2} \sin^2\theta \left[\frac{1}{2} - \frac{1}{2}\mathcal{P}^2 + \mathcal{P}^2 \cos^2\phi \right] \right\}$$

$$- \frac{e^2 g_\ell g}{16\pi^2} \frac{1}{s(s - M_z^2)} \left\{ 2\, \mathrm{Re}U_1 + 2\, \mathrm{Re}U_2 \frac{\vec{p}^2}{M^2} \sin^2\theta \left[\frac{1}{2} - \frac{\mathcal{P}^2}{2} + \mathcal{P}^2 \cos^2\phi \right] \right.$$

$$\left. + \mathrm{Re}U_3 \; g|\vec{p}| \frac{\sqrt{s}}{M^2} \cos\theta - 2\, \mathrm{Im}U_2 \; g\mathcal{P}^2 \frac{\vec{p}^2}{2M^2} \sin^2\theta \sin 2\phi \right\}$$

(III,1.4)

where $\vec{p}$ is the detected hadron momentum,

s is the square of the total center-of-mass energy,

$\mathcal{P}$ is the magnitude of the transverse polarisation of e^+ and e^-,

ϕ is the azimuthal angle of p measured from the direction of the polarisation vector of e^-,

the W's and U's are structure functions that are defined below,

and M is the detected hadron mass.

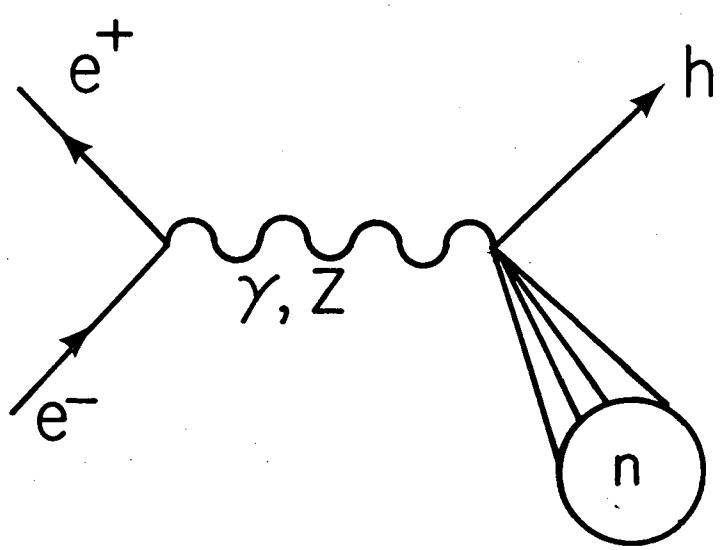

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Fig. 7

We have chosen the **z-axis along** the e^- beam and the **x-axis along the e^- polarization vector.** (θ, ϕ) are the usual spherical angles defined from these 2 axes.

Consider the hadronic tensors $H_{\mu\nu}$ in its most general form allowed by Lorentz invariance,

$$H_{\mu\nu} \equiv (2\pi)^3 (2 E_p) \sum_{\substack{\text{all } n \text{ ex-} \\ \text{cept } \vec{p}}} \langle 0 | J_\mu | \vec{p}\, n \rangle \langle \vec{p}\, n | J_\nu | 0 \rangle \delta^4(q - p - n)$$

$$= -W_1 g_{\mu\nu} + W_2 \frac{p_\mu p_\nu}{M^2} - i \epsilon_{\mu\nu\alpha\beta} \frac{p^\alpha q^\beta}{2M^2} W_3 +$$

$$\frac{W_4}{2M^2} (p_\mu q_\nu + p_\nu q_\mu) + i \frac{W_5}{2M^2} (p_\mu q_\nu - p_\nu q_\mu) + W_6 \frac{q_\mu q_\nu}{M^2} .$$

(III.1.5)

If both J_μ in Eq. III.1.5 are the hadronic electromagnetic current, then hermiticity of the current, time-reversal invariance, and electric gauge invariance restrict $H_{\mu\nu}$ to the form

$$H^{(e)}_{\mu\nu} = -W_1 \left\{ g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2} \right\} + \frac{W_2}{M^2} \left(p - q_\mu \frac{p \cdot q}{q^2} \right) \left(p_\nu - q_\nu \frac{p \cdot q}{q^2} \right)$$

(III.1.6)

where W_1 and W_2 are real functions of the Lorentz invariants q^2 and $p \cdot q$.

On the other hand, the interference term becomes, using the same arguments,

$$H^{(ew)}_{\mu\nu} = (2\pi)^3 2E \sum_{\substack{\text{all } n \text{ ex-} \\ \text{cept } p}} \langle 0|J^{(e)}|\vec{p}n\rangle \langle \vec{p}n|J^{(z)}_{v}|0\rangle \; \delta^4(q - p - n)$$

$$= U_1 \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}\right) + \frac{U_2}{M^2}\left(p_\mu - q_\mu \frac{p\cdot q}{q^2}\right)\left(p_\nu - q_\nu \frac{p\cdot q}{q^2}\right)$$

$$-i\,\varepsilon_{\mu\nu\alpha\beta}\, p^\alpha q^\beta \frac{U_3}{2M^2} \qquad \text{(III.1.7)}$$

where $U_i = U_i(q^2, p\cdot q)$ are Lorentz invariant functions, but this time not restricted to be real by time-reversal invariance. Interference of final state interactions in different eigenchannels ($I = 0$ and 1) gives rise to a nonzero imaginary part since

$$\sum_n |n\,\vec{p}\rangle^{out}\,{}^{out}\langle n\,\vec{p}| \neq \sum_n |n\,\vec{p}\rangle^{in}\,{}^{in}\langle n\vec{p}|$$

when $\vec{p}$ is not summed over. The "in" and "out" refers to the usual incoming and outgoing boundary conditions. Looking back at Eq. (III.1.4), we see that the term $\mathrm{Im}\, U_2$ will give rise to a left-right asymmetry. Notice also that the terms $\mathrm{Re}\, U_1$

and $\mathrm{Re}U_2$ have the same angular distribution as the dominant electromagnetic terms and so are insignificant in the calculation of the angular asymmetry that interests us and so we drop them.

We now make use of the successful Bjorken scaling assumption. It is assumed that at these high energies, there are no important parameters with mass dimension; thus the structure functions will depend on the kinematical variables only through the dimensionless $x = \frac{2p \cdot q}{q^2}$. So

$$\frac{p \cdot q}{U^2} \; U_2 (p \cdot q, q^2) \rightarrow F_2^{(ew)}(x)$$

$$\frac{p \cdot q}{M^2} \; U_3 (p \cdot q, q^2) \rightarrow F_3^{(ew)}(x)$$

$$W_1 (p \cdot q, q^2) \rightarrow F_1^{(e)}(x)$$

$$\frac{p \cdot q}{M^2} \; W_2 (p \cdot q, q^2) \rightarrow F_2^{(e)}(x)$$

The cross-section (III.1.4) then becomes

$$E_p \frac{d^3\sigma}{dp^3} \approx \frac{1}{s^2} \left(\frac{e^2}{4\pi}\right)^2 \{F_1^{(e)} + \frac{1}{4}\beta^2 x^{(e)} \sin^2\theta (1 - \mathcal{P}^2 \cos 2\phi)\}$$

(Equation III.1.9 continued on the next page)

$$+ \frac{1}{s(s - M_z^2)} \left(\frac{e^2 g_a g_\ell g_h}{16^2}\right) \{\frac{1}{2} \beta^2 x \mathcal{P}^2 \mathrm{Im}\, F_2^{(ew)} \sin^2\theta \ \sin 2\phi \ -$$

$$\beta \,\mathrm{Re}\, F_3^{(ew)} \cos\theta\}. \qquad \text{(III.1.9)}$$

III.2 The Left-Right Asymmetry

Let us define a right-left asymmetry parameter at $\theta = 90^\circ$ as

$$\Delta(\phi) = \{N(\phi) - N(-\phi)\}/\{N(\phi) + N(-\phi)\} , \qquad \text{(III.2.1)}$$

where $N(\phi)$ is the number of hadrons coming out in the direction of $(\theta = 90^\circ, \phi)$. Then

$$\Delta(\phi) - \Delta(\pi - \phi) = \left(\frac{g_a g_\ell g_h}{e^2}\right) \left(\frac{s}{s - M_z^2}\right) \frac{\mathcal{P}^2 \sin 2\phi}{2 - \alpha(x)(1 + \mathcal{P}^2 \cos 2\phi)} \cdot$$

$$\cdot \frac{x \,\mathrm{Im}\, F_2^{(ew)}}{F_1^{(e)}} .$$

At $\sqrt{s} = 7.4$ GeV, $\alpha(x)$ is close to unity for $s \gtrsim 0.5$, and $\mathcal{P}$ is about 70%.[51] It is expected that at higher energies $\alpha(x)$ will be close to unity above even smaller values of x. Hadrons produced by the electromagnetic current are depleted at $\theta = 90^\circ$ with $\phi = 0^\circ$ and 180°, while the right-left asymmetry shows up maximally at $\phi = \frac{1}{2} \cos^{-1}\{\alpha \mathcal{P}^2/(2 - \alpha)$.

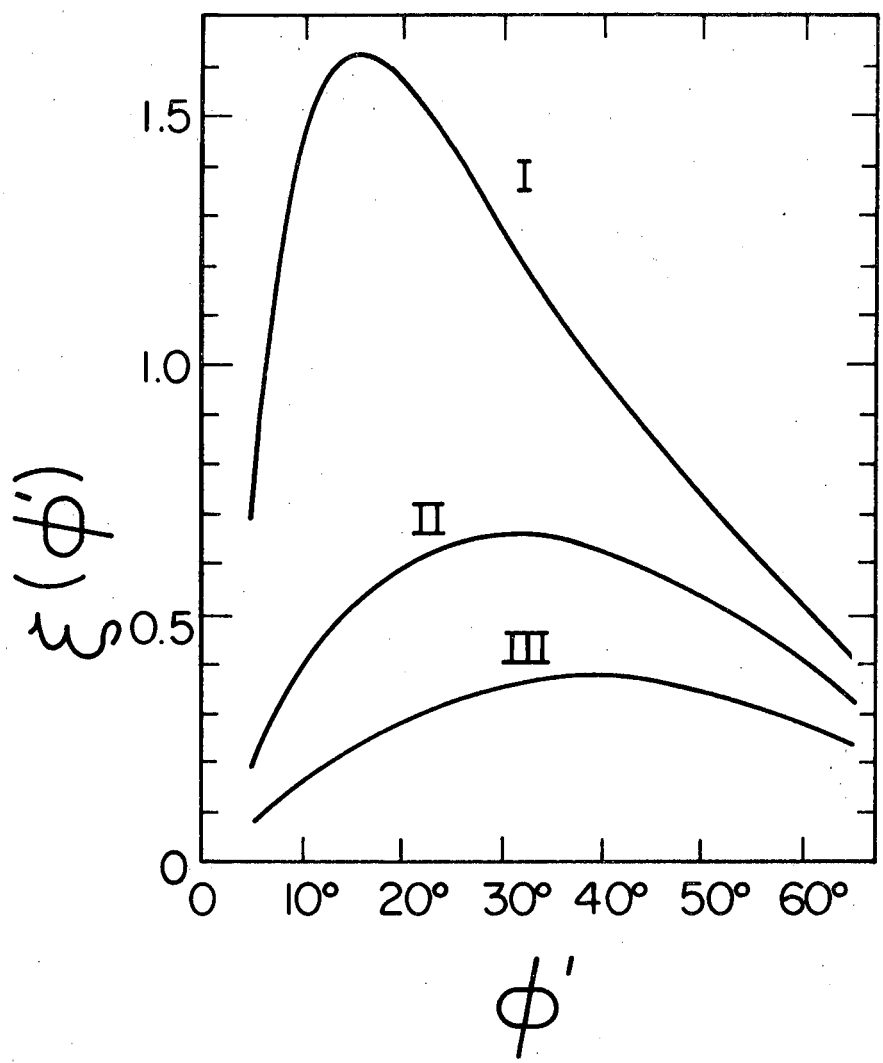

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Fig. 8

In the standard Salam-Weinberg model, the weak couplings are given by $g_a g_\ell g_h = \frac{1}{2} e^2/\sin^2 2\theta_w$, where $\sin^2 \theta_w \approx 0.30$ to fit experimental data. Equation II.2.2 is written in this model as

$$\Delta(\phi) - \Delta(\pi - \phi) \approx -0.29(\sqrt{s}/37 \text{ GeV})^2 \; \xi(\phi) \; \frac{x \text{ Im } F_2^{(ew)}}{F_1^{(e)}} \; , \tag{III.2.3}$$

$$\xi(\phi) = \mathcal{P}^2 \sin 2\phi / \{2 - \alpha(x) - \alpha(x)\mathcal{P}^2 \cos 2\phi\} \; . \tag{II.2.4}$$

The function $\xi(\phi)$ is of the order of unity or less. It is plotted in Fig. 8 for a few different values of $\mathcal{P}^2$ and $\alpha(x)$. A crucial question is how large $x \text{ Im } F_2^{(ew)}/F_1^{(e)}$ would be. If all parton distribution functions are chosen equal, the parton model predicts that in the four-quark model

$$|x F_2^{(ew)} / F_1^{(e)}| \approx \frac{9}{5} - 4 \sin^2 \theta_w \approx 0.60. \tag{III.2.5}$$

III.3. Comment

Only a product of currents antisymmetric in isovector and isoscalar currents can contribute to $\text{Im } F_2^{(ew)}$. Estimate of its imaginary part involves details of strong interaction dynamics at low energies, which forbids anything beyond a mere guess. We cannot even exclude a pessimistic guess that the phase difference of $I = 0$ and $I = 1$

final state interactions may come out to be negligibly small. We find from the isospin structure of currents that the ratio $\mathrm{Im}\, F_2^{(ew)}/F_2^{(ew)}$ is about 0.50 at maximum interference in the Salam-Weinberg model. It should be pointed out that if final states consist only of pions, G parity forbids the interference. Only the events that contain $K\bar{K}$, $N\bar{N}$, and so forth contribute to the asymmetry. At higher energies such events occupy a portion of cross section large enough to produce a substantial right-left asymmetry. Although this test seems to be rather difficult at SPEAR/DORRIS energies, it will be quite realistic at PEP/PETRA energies ($\sqrt{s} = 20 \sim 36$ GeV).

To conclude, we emphasize that unlike the forward-backward asymmetry the right-left asymmetry is a clean direct test of parity nonconservation free from backgrounds due to higher order electromagnetic interactions or to weak decays of heavy leptons or charmed hadrons produced electromagnetically. If a hadronic weak neutral current is parity-conserving but a leptonic current contains both vector and axial vector currents, there would be no $\cos\theta$ asymmetry besides a possible secondary effect due to weak decays in final states and therefore the present method would be useful to supplement information obtainable in the $e^+e^- \to \mu^+\mu^-$ process.

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REFERENCES

1. J. Schwinger, Ann. Phys. (N. Y.) 2, 407 (1957).
2. S. L. Glashow, Nucl. Phys. 22, 579 (1961).
3. A. Salam and J. Ward, Phys. Lett. 13, 168 (1964).
4. R. Feynmann and M. GellMann, Phys. Rev. 109, 193 (1958), R. E. Marshak and E. C. G. Sudarshan, Phys. Rev. 109, 1860 (1958).
5. J. Goldstone, Nuovo Cimento 19, 154 (1961).
6. J. Goldstone, A. Salam and S. Weinberg, Phys. Rev. 127, 965 (1962); S. Bludman and A. Klein, Phys. Rev. 131, 2363 (1963); G. Jona-Lasinio, Nuovo Cimento, 34, 1790 (1964).
7. P. Higgs, Phys.Rev. Lett.12, 132 (1964).
8. T. W. Kibble, Phys. Rev. 155, 1554 (1967); G. Guralnik, C. R. Hagen, and T. W. Kibble, Phys. Rev. Lett. 13, 585 (1964); F. Englert and R. Brout, Phys. Rev. Lett. 13, 321 (1964).
9. S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967).
10. A. Salam, Elementary Particle Theory, N. Svartholm, Editor (Almquist and Forlag, A. B. Stockholm, 1968).
11. S. L. Glashow, J. Iliopoulos and L. Maiani, Phys. Rev. D2, 1285 (1970).
12. G. t'Hooft, Nucl. Phys. B33, 173 (1971); G. t'Hooft, Nucl. Phys. B35, 167 (1971).
13. J. D. Bjorken, Phys. Rev. 179, 1547 (1969).

14. D. J. Gross and F. Wilczek, Phys. Rev. Lett. 30, 1343 (1973).

15. H. D. Politzer, Phys. Rev. Lett. 30, 1346 (1973).

16. S. Coleman and D. J. Gross, Phys. Rev. D8, 4383 (1973).

17. B. J. Bjorken and S. Glashow, Phys. Lett. 11, 255 (1964).

18. S. Glashow, J. Iliopoulos, L. Maiani, Phys. Rev. D2, 1285 (1970).

19. G. Golhaber, Lecture at the "International Scholl of Physics Enrico Fermi", July 1977.

20. P. Musset, Report to 18th International Conference on High Energy Physics, Tblisi, USSR (1976).

21. H. Faissner, Report to 18th International Conference on High Energy Physics, Tblisi, USSR (1976).

22. F. Reines, H. Gurr and H. Sobel, Phys. Rev. Lett. 37, 315 (1976).

23. A. Benvenuti et. al., Phys. Rev. Lett. 37, 1039 (1976).

24. D. Buchholz, Report to Aachen Neutrino conference (1976).

25. W. Lee et. al., Phys. Rev. Lett. 37, 186 (1976).

26. D. Cline et. al., Phys. Rev. Lett. 37, 252 (1976); 37, 648 (1976).

27. L. Lederman, European Conference on Particle Physics, Budapest, Hungary, July 1977.

28. P. E. G. Baird et. al., Nature 264, 528 (1976).

29. E. Henley and L. Wilets, Phys. Rev. (to be published).

30. M. Brinicombe, C. Loving and P. Sanders, Oxford University preprint.

31. M. L. Perl et al., Phys. Rev. Lett. 35, 1489 (1975).

32. Y.-S. Tsai, Phys. Rev. D 4, 2821 (1971).

33. K. Fujikawa and N. Kawamoto, Phys. Rev. Lett. 35, 1560 (1975); Phys. Rev. D 13, 2534 (1976).

34. S.-Y. Pi and A. I. Sanda, Phys. Rev. Lett. 36, 1 (1976).

35. A. Pais and S. B. Treiman, Phys. Rev. D 14, 293 (1976); K. J. F. Gaemers and R. Raitio, Phys. Rev. D 14, 1262 (1976).

36. R. F. Schwitters et al., Phys. Rev. Lett. 35, 1320 (1975).

37. The only realistic test utilizing transverse polarizations is, to our knowledge, the right-left asymmetry in inclusive hadron production at PEP and PETRA energies. R. Simard and M. Suzuki, Phys. Lett. 63B, 304 (1976).

38. R. F. Schwitters, Nucl. Instrum. Methods 117, 331 (1974); R. F. Schwitters and B. Richter, SLAC-LBL Note No. SPEAR-175, PEP-87; A. Garren and J. Kadyk, SLAC-LBL Note No. PEP-184 (unpublished). See also A. W. Chao and R. F. Schwitters, SLAC-LBL Note No. SPEAR-194, PEP-217 (unpublished). This points out a serious difficulty in obtaining a large polarization.

39. J. D. Bjorken and S. D. Drell, Relativistic Quantum Mechanics (McGraw-Hill, New York, 1964).

40. R. E. Marshak, Riazuddin, and C. P. Ryan, Theory of Weak Interactions in Particle Physics (Wiley-Interscience, New York, 1969), p. 185.

41. B. C. Barish et. al., Phys. Rev. Lett. 32, 1387 (1974).

42. M. L. Perl, SLAC-PUB-1923, April 1977.

43. G. J. Feldman, and M. L. Perl, SLAC-PUB-1972, July 1977.

44. C. H. Albright, Phys. Rev. D 13, 2508 (1976).

45. A. Ali and T. C. Yang, Phys. Rev. D 14 , 3052 (1976).

46. See for a summary R. F. Schwitters, in Proceedings of the 1975 International Symposium on Lepton and Photon Interactions at High Energies, Stanford, California, edited by W. T. Kirk (SLAC, Stanford 1976) p. 5.

47. It is easier to measure an asymmetry between the forward and backward hemispheres. With the integral over cos θ, this asymmetry is given by the right-hand side of Eq. (3.1), multiplied by $\frac{1}{2}$.

48. R. Gatto and G. Preparata, Riv. Nuovo Cimento 4, 445 (1974).

49. A. W. Chao and R. F. Schwitters, SLAC-LBL Note No. Spear-194, PEP-217, 1976 (unpublished).

50. N. Cabibbo and R. Gatto, Phys. Rev. 124, 1577 (1961); R. Gatto, Springer Tracts in Modern Physics 39, 106 (1965); R. Gatto and G. Preparata, Riv. Nuovo Cim. 4, 445 (1974); A. McDonald, Nucl. Phys. B 75, 343 (1974); J. Boldine and A. Hankey, Phys. Rev. D 6, 3301 (1972); A. Love, Lett. Nuovo Cim. 5, 113 (1972); R. Bundy, Phys. Lett. B 45, 340 (1973); R. Lendvai and G. Posik, Phys. Lett. B 56, 462 (1975).

51. R. Schwitters, Proc. 1975 Internation Symp. on Lepton and photon interactions at high energies (ed. W. J. Kirk), Stanford University (1975), p. 5.

FIGURE CAPTIONS

Fig. 1: A forbidden configuration ($\nu_L \equiv \nu_e$).

Fig. 2: A μe event with an electron beam polarised longitudinally.

Fig. 3: Kinematics of μe events in the e^+e^- laboratory frame.

Fig. 4: $R(x;\ \cos^2\theta)$ and asymmetry parameter $\alpha(x;\ \cos^2\theta)$ as functions of $x = \frac{|\vec{p}|}{k_0}$ at $\cos\theta = 1$. See Eq. (I.2.21) in text for the definitions of $R(x;\ \cos^2\theta)$ and $\alpha(x;\ \cos^2\theta)$.

Fig. 5: Asymmetry parameter $\alpha(x;\ \cos^2\theta)$ as a function of $\cos\theta$ at $x = 0.5$ and $x = 0.05$. The value of γ is equal to 10 in both cases.

Fig. 6: Dependence on γ of asymmetry parameter $\alpha(x;\ \cos^2\theta)$ at $x = 0.5$ and $x = 0.05$. The value of θ is equal to 0° in both cases.

Fig. 7: The one boson exchange diagram.

Fig. 8: $\xi(\phi)$ for three sets of values of $\mathcal{P}^2$ and $\alpha(x)$.

I. for $\mathcal{P}^2 = 0.85$ and $\alpha(x) = 1.0$.

II. for $\mathcal{P}^2 = 0.70$ and $\alpha(x) = 0.80$.

III. for $\mathcal{P}^2 = 0.50$ and $\alpha(x) = 0.60$.

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