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THE STRUCTURE OF NEUTRAL CURRENTS[*]

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CHAPTER HEADINGS

CHAPTER HEADINGS   (continued)

I. INTRODUCTION

One of the major advances in particle physics of the last decade
has been the experimental discovery of neutral currents by the
Gargamelle collaboration (Hasert et al., 1973) which was quickly
confirmed by the Harvard-Pennsylvania-Wisconsin collaboration
(Benvenuti et al., 1974). This far-reaching event has generated a
tremendous effort on both the experimental and theoretical fronts in
the attempt to clarify the structure of neutral currents. The purpose
of this review is the assessment of where we stand in the determi-
nation of neutral-current coupling parameters. As we will show in
the subsequent chapters, what comes out from all the various
phenomenological analyses is the confirmation of the "standard"
SU(2) ⊗ U(1) Glashow, Weinberg, Salam, Ward gauge model as the most
viable theory describing low-energy electroweak phenomena.

Historically speaking, the existence of neutral currents other
than the familiar electromagnetic currents was predicted as early as
1958 by Bludman (Bludman, 1958) who constructed a model based on a
local SU(2) gauge symmetry. The model incorporated both charged
(responsible for β- decay) and neutral currents. The space time
structure of the neutral currents in his model is pure V-A (vector
minus axial-vector) and thus they cannot be identified with the
electromagnetic currents which are (parity-conserving) purely
vectorial. There is no unification with electromagnetism in Bludman's
model. A model truly unifying weak and electromagnetic interactions
incorporating two kinds of neutral currents (electromagnetic and weak)

was invented by Glashow (1961) and by Salam and Ward (1964). It is called the SU(2) ⊗ U(1) model. As the model stands, there is no mechanism for mass generation of the intermediate vector bosons and thus the relative strength of weak neutral-current interactions to that of charged-current interactions is a completely free parameter. This problem was settled in 1967 by Weinberg (Weinberg, 1967) who incorporated the idea of spontaneous breakdown of local gauge symmetry (Higgs, 1964; Englert and Brout, 1964) into the SU(2) ⊗ U(1) model. A somewhat similar mechanism was proposed in 1968 by Salam (Salam, 1968). The mass of the intermediate boson which mediates the neutral current interactions, called the Z-boson, is related in a definite way to that of its charged counterparts, called the W-boson. The relative strength is therefore fixed once and for all, in this minimal version of the SU(2) ⊗ U(1) model. There was thus a definite prediction of the structure of weak neutral currents (a mixture of vector and axial vector currents) and its strength of interactions. As it stands, it is a <u>single</u> parameter ($\sin^2\theta_W$) theory.

With the discovery, in 1973, of neutral currents, the SU(2) ⊗ U(1) model stood out as a strong candidate for a theory of electroweak interactions. With more data pouring in, it was customarily in those days to check whether or not one set of data gives a value of $\sin^2\theta_W$ consistent with another set of data. There is certainly a danger in this because in so doing we are already assuming that half of the theory is right. Even in the SU(2) ⊗ U(1) model itself, in principle there are more than just one single parameter $\sin^2\theta_W$. We have decided to take a completely phenomenological approach to the problem.

We will show in the next chapter that neutral-current phenomena involving neutrino-quark, neutrino-electron, electron-quark and electron-muon reactions can be characterized by thirteen independent parameters. The fundamental goal of neutral-current physics is therefore the complete determination of all these thirteen independent parameters. Once the goal is accomplished, we can then compare these parameters to one's favorite gauge model. It is then that one can decide whether or not the minimal $SU(2) \otimes U(1)$ model is correct or some other alternative is more desirable.

Fortunately, we did not have to have all thirteen parameters to be able to make that decision. It turns out that data from neutrino-hadron and (polarized) electron hadron are sufficient to determine the low-energy structure of weak neutral currents and its strength of interactions. If the neutral currents were mediated by a single massive neutral boson then there should be a set of factorization relations among various neutral-current parameters. In fact, we will show that some of these relations are indeed consistent with current data, while the others are not yet checked due to the lack of available data. Besides the objectivity of the phenomenological analysis, there is also a considerable subjective interest in a precise determination of $\sin^2\theta_W$ once we assume that the $SU(2) \otimes U(1)$ model is the correct electroweak theory. A precise value of $\sin^2\theta_W$ has important implications on grand unified models which try to unify the electroweak and strong interactions.

The organization of our review is as follows. In Chapter II, we define the neutral-current parameters which enter on various processes. We also show the forms that these parameters take in different models. Chapter III deals with the analysis of neutrino-electron (νe) reactions. Chapter IV gives a detailed analysis of neutrino-hadron (νH) interactions and shows how a unique solution for the neutrino-quark parameters is chosen. This unique solution is in excellent agreement with the standard $SU(2) \otimes U(1)$ model. In Chapter V, the discussion centers on the neutral-current couplings between electrons and hadrons. The relevant experiments are the SLAC polarized electron-deuteron $(e_{L,R}D)$ and the atomic physics experiments (Seattle, Oxford, Berkeley, Novosibirsk...). It is shown how the SLAC experiment decisively selected the minimal standard $SU(2) \otimes U(1)$ model over other gauge alternatives while most of the atomic physics experiments finally came into agreement with the standard model after a tumultuous period of controversial results. In Chapter VI, high-energy e^+e^- reactions are discussed with the emphasis on recent data from PETRA. We discuss the final muon forward-backward asymmetry as well as possible effects from a low mass Z-boson or from a multi-Z boson model. Factorization relations to test the single Z-boson hypothesis are detailed in Chapter VII. The future outlook is discussed in Chapter VIII.

II. NOTATIONS AND MODELS

2.1 Notations

Neutral-current coupling occurring in various reactions can be parameterized under the following (rather general) assumptions:

(a) Low-energy neutral-current reactions can be described by an effective four-fermi (or current-current) interactions in the same manner as the low-energy charged-current reactions. This assumption is reinforced by the apparent absence of any low mass intermediate boson (no low mass Z-resonance in e^+e^- annihilation, no effects on scaling violation due to the propagator effect). If the interactions are mediated by a neutral Z-boson, this assumption means that the momentum transfer $Q^2 \ll m_Z^2$ (mass of Z-boson).

(b) The weak neutral currents are linear combinations of V(vector) and A(axial-vector) currents. This is a characteristic of all electroweak gauge models. Furthermore, the y-distribution in deep-inelastic scattering rules out any combination of pure S(scalar) and P(pseudoscalar). The same conclusion comes from the SLAC $e_{L,R}D$ experiment since there, the electromagnetic current which conserves helicity, cannot interfere with S, P and T (helicity flip) while experimentally weak and electromagnetic interference is observed. It is still possible that effective interactions involving quarks and leptons (νq and eq e.g.) could contain scalar-scalar (S-S) and pseudoscalar-pseudoscalar (P-P) in addition to the parity

conserving VV + AA. This could arise from a Fierz transformation of $\bar{\ell}\gamma_\mu q \bar{q} \gamma_\mu \ell$, where ℓ stands for leptons and q stands for quarks, which is brought about by lepto-quark gauge boson exchanges. The strenght of these types of parity-conserving effective interaction would be expected to be extremely feeble and could be ignored as far as available experiments are concerned.

(c) The absence of <u>second-class</u> neutral currents. In this connection, we remind the reader that <u>second-class</u> charged currents have been ruled out by recent data on beta-decay processes. See the review article by C. S. Wu (1978) for more details.

(<u>d</u>) Weak neutral currents are flavor-conserving. Experimentally there is a very strong suggestion pointing towards the absence strangeness changing neutral-currents and to a certain extent, also charm-changing neutral currents. Theoretically, the absence of flavor-changing neutral currents is guaranteed by the Glashow-Iliopoulos-Maiani (GIM) mechanism (Glashow, Iliopoulos, Maiani, 1970) or its generalized version, the Kobayashi-Maskawa mechanism (Kobayashi and Maskawa, 1973).

(<u>e</u>) The hadronic (or quark) neutral-currents are a combination of strong isovector (I = 1) and (strong) isoscalar (I = 0) currents. Experimentally there is no evidence for a component having isospin two or higher in electromagnetic and charged-current interactions. Furthermore, in the conventional quark model, it is impossible to build an (quark bilinear) object with (strong) isospin two or higher.

Having listed all the assumptions needed, we can write down the following effective Lagrangians (Hung and Sakurai; 1976, 1977a).

NEUTRINO-NUCLEON INTERACTIONS   To the extent that nucleons are made up of valence u and d quarks, the effective Lagrangian describing νN interactions and obeying assumptions (a) - (e), contains four independent parameters α, β, γ and δ, and is given by

$$\mathcal{L} = -\frac{G}{\sqrt{2}} \, \bar{\nu}\gamma_\lambda (1 + \gamma_5)\nu \left\{ \frac{1}{2}[\bar{u}\gamma_\lambda(\alpha + \beta\gamma_5)u - \bar{d}\gamma_\lambda(\alpha + \beta\gamma_5)d] \right.$$

$$\left. + \frac{1}{2} [\bar{u}\gamma_\lambda(\gamma + \delta\gamma_5)u + \bar{d}\gamma_\lambda(\gamma + \delta\gamma_5)d] \right\}. \qquad (2.1)$$

The parameters α, β, γ and δ have the following meaning: $\alpha \equiv$ isovector vector; $\beta \equiv$ isovector axial-vector; $\gamma \equiv$ isoscalar vector; $\delta \equiv$ isoscalar axial-vector. Here isovector and isoscalar mean respectively $\frac{1}{2}(\bar{u}u - \bar{d}d)$ and $\frac{1}{2}(\bar{u}u + \bar{d}d)$. In our convention, $\frac{1}{2}(1 + \gamma_5)$ and $\frac{1}{2}(1 - \gamma_5)$ are the left-handed and right-handed projection operators respectively. One can alternatively parameterize νN interactions by four "chiral coupling constants" $\varepsilon_L(u)$, $\varepsilon_R(u)$, $\varepsilon_L(d)$ and $\varepsilon_R(d)$ (Bjorken, 1976; Sehgal 1977) which are related to α, β, γ and δ by

$$\varepsilon_L(u) = \frac{1}{4}(\alpha + \beta + \gamma + \delta), \quad \varepsilon_R(u) = \frac{1}{4}(\alpha - \beta + \gamma - \delta),$$

$$\varepsilon_L(d) = \frac{1}{4}(-\alpha - \beta + \gamma + \delta), \quad \varepsilon_R(d) = \frac{1}{4}(-\alpha + \beta + \gamma - \delta),$$

$$(2.2)$$

and which are the coefficients of $\frac{1}{2}\bar{u}\gamma_\lambda(1 \pm \gamma_5)u$ and $\frac{1}{2}\bar{d}\gamma_\lambda(1 \pm \gamma_5)d$ respectively.

NEUTRINO-ELECTION REACTIONS   Assuming $\mu - e$ universality, the effective Lagrangian describing $\nu_\mu e$, $\bar{\nu}_\mu e$ and $\bar{\nu}_e e$, reactions (both charged and neutral) is given by

$$\mathcal{L} = \mathcal{L}_{CC} + \mathcal{L}_{NC} \quad, \tag{2.3}$$

where

$$\mathcal{L}_{CC} = -\frac{G}{\sqrt{2}} \, [\bar{\nu}_e \gamma_\lambda(1 + \gamma_5)e][\bar{e}\gamma_\lambda(1 + \gamma_5)\nu_e], \tag{2.4a}$$

$$\mathcal{L}_{NC} = -\frac{G}{\sqrt{2}} \, [\bar{\nu}_\mu \gamma_\lambda(1 + \gamma_5)\nu_\mu + \bar{\nu}_e \gamma_\lambda(1 + \gamma_5)\gamma_e]\bar{e}\gamma_\lambda(g_V + g_A\gamma_5)e \tag{2.4b}$$

Making a Fierz transformation in Eq. (2.4a), we obtain the following neutral-current effective Lagrangian

$$\widetilde{\mathcal{L}}_{NC} = -\frac{G}{\sqrt{2}} \left\{ \bar{\nu}_\mu \gamma_\lambda(1 + \gamma_5)\nu_\mu \bar{e}\gamma_\lambda(g_V + g_A\gamma_5)e \right.$$

$$\left. + \bar{\nu}_e \gamma_\lambda(1 + \gamma_5)\nu_e \bar{e}\gamma_\lambda(G_V + G_A\gamma_5)e \right\} , \tag{2.5}$$

where

$$G_V = 1 + g_V, \ G_A = 1 + g_A, \tag{2.6}$$

and the 1 in (2.6) comes from the charged current contribution. We have seen that with the assumption of $\mu - e$ universality, $\nu_\mu e$, $\bar{\nu}_\mu e$

and $\bar{\nu}_e$ e reactions can be described by just two independent parameters g_V and g_A. $\bar{\nu}_e$ e data comes from reactor experiments at low energy. At the moment, no good data exists yet for ν_e e.

ELECTRON-NUCLEON REACTIONS   This class of reactions involves the SLAC $e_{L,R}$ D and atomic physics experiments. The parity conserving part of the neutral-current interactions between electrons and nucleons is completely overwhelmed by the much stronger electromagnetic interactions. We, therefore, concentrate on the <u>parity non-conserving</u> effective Lagrangian which is described by four independent parameters $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}$ and $\tilde{\delta}$ in analogy with the νN case. These parameters have similar meanings to α, β, γ and δ. We have

$$\mathcal{L} = -\frac{G}{\sqrt{2}} \left\{ \bar{e}\gamma_\lambda\gamma_5 e \left[\frac{\tilde{\alpha}}{2} (\bar{u}\gamma_\lambda u - \bar{d}\gamma_\lambda d) + \frac{\tilde{\gamma}}{2} (\bar{u}\gamma_\lambda u + \bar{d}\gamma_\lambda d) \right] \right.$$

$$\left. + \bar{e}\gamma_\lambda e \left[\frac{\tilde{\beta}}{2} (\bar{u}\gamma_\lambda\gamma_5 u - \bar{d}\gamma_\lambda\gamma_5 d) + \frac{\tilde{\delta}}{2} (\bar{u}\gamma_\lambda\gamma_5 u + \bar{d}\gamma_\lambda\gamma_5 d) \right] \right\}.$$

$$(2.7)$$

This effective Lagrangian has the form: $A_{lepton}V_{quark} + V_{lepton}A_{quark}$. Both <u>terms</u> contribute to the SLAC $e_{L,R}$ D experiment and atomic physics experiments on hydrogen and deuterium. Atomic physics experiments on heavy atoms (Bismuth, Thallium, etc...) probe the $A_{lepton}V_{quark}$ term.

ELECTRON-POSITRON ANNIHILATION INTO MUON PAIRS   Assuming $\mu - e$ universality and with the center-of-mass energy $\sqrt{s}$ being much lower than any intermediate boson mass, the most general weak neutral-current effective Lagrangian for $e^+e^- \rightarrow \mu^+\mu^-$ involving V,A couplings can be written as

$$\mathcal{L} = -\frac{G}{\sqrt{2}} \left\{ h_{VV}(\bar{e}\gamma_\lambda e + \bar{\mu}\gamma_\lambda \mu)(\bar{e}\gamma_\lambda e + \bar{\mu}\gamma_\lambda \mu) \right.$$

$$+ 2h_{VA}(\bar{e}\gamma_\lambda e + \bar{\mu}\gamma_\lambda \mu)(\bar{e}\gamma_\lambda \gamma_5 e + \bar{\mu}\gamma_\lambda \gamma_5 \mu)$$

$$\left. + h_{AA}(\bar{e}\gamma_\lambda \gamma_5 e + \bar{\mu}\gamma_\lambda \gamma_5 \mu)(\bar{e}\gamma_\lambda \gamma_5 e + \bar{\mu}\gamma_\lambda \gamma_5 \mu) \right\} . \qquad (2.8)$$

The constants h_{VV}, h_{AA} and h_{VA} could, in principle, be measured by studying the magnitude and energy dependence of the cross section, the forward-backward asymmetry in the angular distribution of the final muon, and the muon logitudinal polarization, respectively. When center-of-mass energy approaches the intermediate boson mass, the propagator effect will have to be taken into account, producing a characteristic energy dependence in the forward-backward asymmetry of the muon.

We have thus seen that "low-energy" neutral-current experiments can be described by thirteen parameters; $\alpha, \beta, \gamma, \delta$, g_V, g_A, $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}, \tilde{\delta}$, h_{VV}, h_{AA} and h_{VA}. The phenomenological role of neutral-current physics is to determine all these couplings. (See Fig. 2.1).

2.2 The Standard Model

Since the concepts and construction of a renormalizable, spontaneously broken gauge theory are explained in many excellent reviews, we only briefly recall the basic features of the most popular model, the standard $SU(2) \otimes U(1)$ model and its generalization, with the purpose of establishing notations. In particular, the minimal model and its generalization will be parameterized by a certain number of parameters and we will show how these parameters enter the thirteen neutral-current couplings discussed earlier.

What is now commonly called the "standard" model is in fact the Weinberg-Salam version of the Glashow-Salam-Ward $SU(2)_L \otimes U(1)$ model. In the Weinberg-Salam version, the local $SU(2)_L \otimes U(1)$ symmetry (weak isospin + weak hypercharge gauge groups) is spontaneously broken down to $U(1)_{em}$ (the electromagnetic gauge group) by a weak doublet of scalars $\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}$ which acquires a non-vanishing vacuum expectation value $< \Phi > = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}$. In the process, three gauge bosons W^+, W^- and Z acquire masses while one remains massless, the photon A, corresponding to the unbroken gauge group $U(1)_{em}$. The masses of $W^\pm$ and Z bosons are given by (in terms of $< \Phi >$)

$$m_W = \frac{1}{2} gv,$$

$$m_Z = \frac{1}{2} (g^2 + g'^2)^{1/2} v, \qquad (2.9)$$

where g and g' are the gauge couplings of $SU(2)_L$ and U(1) respectively. With the mixing angle θ_W defined by (e being the electromagnetic coupling)

$$\tan \theta_W = \frac{g'}{g}, \quad e = g \sin\theta_W = g'\cos \theta_W. \qquad (2.10)$$

the relationship between m_W and m_Z is just

$$m_W = m_Z \cos \theta_W . \qquad (2.11)$$

In this minimal model, each helicity of fermions has the

following $SU(2)_L$ (weak isospin) assignments (doublets for left-handed

fermions and singlets for right-handed ones):

left-handed: $\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L$, $\begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$, $\ldots$

$\begin{pmatrix} u \\ d \end{pmatrix}_L$, $\begin{pmatrix} c \\ s \end{pmatrix}_L$, $\ldots$

Right-handed: e_R^-, μ_R^-, u_R, d_R, $\ldots$. The weak neutral-current $J^{N.C.}$
is coupled to the Z-boson by

$$\mathcal{L}^{N.C.} = (g/\cos\theta_W)\ J_\lambda^{N.C.} Z_\lambda \ , \tag{2.12}$$

where

$$J_\lambda^{N.C.} = J_\lambda^3 - \sin^2\theta_W J_\lambda^{em} \ . \tag{2.13}$$

The most general form of J_λ^3 can be written as

$$J_\lambda^3 = i\Sigma \bar{f}_i \gamma_\lambda \left[\frac{(1 + \gamma_5)}{2} T_{3L}^i + \frac{(1 - \gamma_5)}{2} T_{3R}^i \right] f_i \ , \tag{2.14}$$

where the sum is over all fermions. In the "standard" $SU(2) \otimes U(1)$

model, $T_{3L}^i = \pm\frac{1}{2}$ and $T_{3R}^i = 0$. In Chapter IV, we will show that, by

taking $T_{3L}^i = \pm\frac{1}{2}$ and keeping T_{3R}^i arbitrary, the νN data confirm the

assignment $T_{3R}^i = 0$.

From Eq. (2.13), the low-energy effective Lagrangian in the SU(2) $\otimes$ U(1) model can be written as

$$\mathcal{L}_{eff} = 4\,\frac{G}{\sqrt{2}}\,\rho(J_\lambda^3 - \sin^2\theta_W J_\lambda^{em})^2 \;, \tag{2.15}$$

where

$$\rho = m_W^2/m_Z^2 \cos^2\theta_W \;, \tag{2.16}$$

$$G/\sqrt{2} = g^2/8m_W^2 \;. \tag{2.17}$$

We have seen that in the minimal model (i.e. with only Higgs doublets), $m_W = m_Z \cos\theta_W$ which gives $\rho = 1$. From a phenomenological point of view, we would keep ρ arbitrary since we could not be sure that the Higgs scalars which break SU(2) $\otimes$ U(1) come in doublets. The expressions for the thirteen parameters $\alpha, \beta, \ldots, h_{VA}$ in terms of the SU(2) $\otimes$ U(1) parameters (ρ and T_{3R}^i arbitrary, $T_{3L} = \pm\frac{1}{2}$) are given in Table 2.1. Also listed are the expressions for the one-parameter "standard" SU(2) $\otimes$ U(1) model. For convenience $x \equiv \sin^2\theta_W$.

In Chapter VII, we will compare the current available data with the "standard" model as well as with some other multi-Z boson models.

III. NEUTRINO-ELECTRON REACTIONS

Because of their simplicity, we first discuss reactions involving neutrinos on electrons, namely $\nu_\mu e$, $\bar{\nu}_\mu e$ and $\bar{\nu}_e e$. The cross-sections for these reactions can be computed from the effective Langrangian, Eq. (2.5), to be ('tHooft, 1971)

$$\frac{d\sigma}{dy}(\nu_\mu e) = 2\frac{G^2 m_e E_\nu}{\pi}\left[g_L^2 + g_R^2(1-y)^2 - \frac{m_e}{E_\nu}y\,g_L g_R\right], \quad (3.1a)$$

$$\frac{d\sigma}{dy}(\bar{\nu}_\mu e) = 2\frac{G^2 m_e E_\nu}{\pi}\left[g_R^2 + g_L^2(1-y)^2 - \frac{m_e}{E_\nu}y\,g_L g_R\right], \quad (3.1b)$$

$$\frac{d\sigma}{dy}(\bar{\nu}_e e) = 2\frac{G^2 m_e E_\nu}{\pi}\left[G_R^2 + G_L^2(1-y)^2 - \frac{m_e}{E_\nu}y\,G_L G_R\right], \quad (3.1c)$$

where $y = E_e/E_\nu$ and

$$g_{L,R} = \frac{1}{2}(g_V \pm g_A), \quad (3.2a)$$

$$G_{L,R} = \frac{1}{2}(G_V \pm G_A) = \begin{Bmatrix} 1 + g_L \\ g_R \end{Bmatrix}. \quad (3.2b)$$

Since $m_e/E_\nu \ll 1$ for most practical purposes, we can ignore the last terms in Eqs. (3.1a) - (3.1c).

We can obtain total cross-sections by integrating Eqs. (3.1a) - (3.1c) from $y = 0$ to $y = 1$. It is then customary to present the experimental results on a $g_V - g_A$ plane by ellipses as one can see from Fig. (3.1)(Buesser, 1980). The $\nu_\mu e$ and $\bar{\nu}_\mu e$ ellipses have axes perpendicular to each other and intersect at 4 allowed regions.

These 4 solutions reflect the sign and vector-axial-vector ambiguities in the couplings g_V and g_A. The sign ambiguity can be resolved by including data from reactor neutrinos, i.e. $\bar{\nu}_e e$ scatterings. The physical reason for this resolution is the fact that the sign of the interference of the charged and neutral-current amplitudes is known and there is a term which is linear in g_V and g_A in the $\bar{\nu}_e e$ cross section. The importance of observing the interference term has been stressed by Kayser et al. (1979). As shown in Fig (3.1), the $\bar{\nu}_e e$ curve intersects with the $\nu_\mu e$ and $\bar{\nu}_\mu e$ ellipses at two regions leaving the vector-axial-vector ambiguity intact. The resolution of this last ambiguity is discussed in Chapter 7 in connection with factorization tests. The best fit (Dydak, 1979) to the data gives

$$g_A = -0.52 \pm 0.06, \quad g_V = 0.06 \pm 0.08 \text{ (axial dominant)}, \quad (3.3)$$

with the VA ambiguity $g_V \leftrightarrow g_A$. This axial dominant solution is in excellent agreement with the "standard" model which, for $\sin^2\theta_W = 0.23$, gives

$$g_A = -0.500, \quad g_V = -0.04 . \quad (3.4)$$

It will be shown in Chapter VII that assuming factorization, the axial dominant solution is indeed the correct one.

We show below the summary of experimental results on $\nu_\mu e$ and $\bar{\nu}_\mu e$ scatterings taken from a recent review talk of F. Dydak (1979). Some of the results are new data coming from the CHARM collaboration (Jonker et al., 1979). We also show in Table (3.1) and (3.2) earlier data from GGM (Blietschau et al., 1978; Armenise et al. 1979 a),

AP (Faissner et al., 1978); CB (Cnops et al. 1978) and BEBC-TST (Armenise et al. 1979 b).

From Table (3.1), the average cross-section from all experiments on $\nu_\mu e$ scattering is

$$\sigma(\nu_\mu e)/E_\nu = (1.6 \pm 0.4) \times 10^{-42} \text{ cm}^2/\text{GeV}. \tag{3.5}$$

In terms of the one parameter standard model, we have

$$\sigma(\nu_\mu e)/E_\nu = \frac{G^2 m_e}{2\pi} (1 - 4 \sin^2\theta_W + \frac{16}{3} \sin^4\theta_W) , \tag{3.6}$$

where

$$G^2 m_e/2\pi = 4.31 \times 10^{-42} \text{ cm}^2/\text{GeV} . \tag{3.7}$$

The comparison of Eq. (3.5) with Eq. (3.6) yields

$$\sin^2\theta_W = 0.22 \begin{array}{c} + 0.08 \\ - 0.05 \end{array} . \tag{3.8}$$

The average cross section for $\bar{\nu}_\mu e$ scatterings from Table (3.2) yields

$$\sigma(\bar{\nu}_\mu e)/E_\nu = (1.3 \pm 0.6) \times 10^{-42} \text{ cm}^2/\text{GeV} . \tag{3.9}$$

The one-parameter standard model gives

$$\sigma(\bar{\nu}_\mu e)/E_\nu = \frac{G^2 m_e}{2\pi} \left(\frac{1}{3} - \frac{4}{3} \sin^2\theta_W + \frac{16}{3} \sin^4\theta_W \right) . \qquad (3.10)$$

From Eq. (3.9) and Eq. (3.10), we get

$$\sin^2\theta_W = 0.23 \begin{array}{c} + 0.09 \\ - 0.23 \end{array} . \qquad (3.11)$$

New data for $\nu_\mu e \to \nu_\mu e$ from the VPI-Maryland-NSF-Oxford-Peking collaboration (Heisterberg et al., 1980) gives

$$\sigma(\nu_\mu e \to \nu_\mu e) = (1.40 \pm 0.30) \times 10^{-42} E_\nu \ cm^2/GeV, \qquad (3.12)$$

which implies

$$\sin^2\theta_W = 0.25 \begin{array}{c} + 0.07 \\ - 0.05 \end{array} . \qquad (3.13)$$

IV. NEUTRINO-HADRON REACTIONS

This set of reactions has been among the determining factors in the model-independent analysis of neutral-current data. The inclusive $\overset{(-)}{\nu}N$ reactions yield very precise measurements and thus can be used for an accurate determination of $\sin^2\theta_W$. On the model-independent side, the combined inclusive, semi-inclusive and exclusive $\overset{(-)}{\nu}N$ reactions were the firsts to give an unambiguous and complete determination of neutrino-hadron neutral-current couplings.

4.1   Deep-inelastic Scattering on Isoscalar Targets

We first discuss the deep-inelastic scattering on isoscalar targets (equal number of protons and neutrons, though in principle it is only approximate). To make the discussion simple we use the language of the quark-parton model and ignore contributions from s, $\bar{s}$, ... Actually the two "master formulae" derived below, are largely independent of the details of the quark-parton model. For example, the effects of scaling violations tend to cancel since these effects occur equally in charged and neutral currents. The cross-sections for both charged and neutral-current interactions are given by (in units of $G^2 M_N E_\nu/\pi$), with $N = \frac{1}{2}(n + p)$,

$$\frac{d\sigma}{dy}^{CC}(\nu N) = Q + \bar{Q}(1 - y)^2,$$

(4.1a)

$$\frac{d\sigma}{dy}^{CC}(\bar{\nu}N) = \bar{Q} + Q(1 - y)^2,$$

(4.1b)

$$\frac{d\sigma}{dy}^{NC}(\nu N) = \left[|\varepsilon_L(u)|^2 + |\varepsilon_L(d)|^2\right]\left[Q + \bar{Q}(1-y)^2\right]$$

$$+ \left[|\varepsilon_R(u)|^2 + |\varepsilon_R(d)|^2\right]\left[\bar{Q} + Q(1-y)^2\right], \qquad (4.1c)$$

$$\frac{d\sigma}{dy}^{NC}(\bar{\nu}N) = \left[|\varepsilon_L(u)|^2 + |\varepsilon_L(d)|^2\right]\left[\bar{Q} + Q(1-y)^2\right]$$

$$+ \left[|\varepsilon_R(u)|^2 + |\varepsilon_R(d)|^2\right]\left[Q + \bar{Q}(1-y)^2\right], \qquad (4.1d)$$

where $y = E_{had}/E_\nu$, $x = Q^2/2P \cdot Q$ (the Bjorken scaling variable) and

$$Q = \int x[u(x) + d(x)]dx, \qquad (4.2a)$$

$$\bar{Q} = \int x[\bar{u}(x) + \bar{d}(x)]dx, \qquad (4.2b)$$

with $\overset{(-)}{u}(x)$ and $\overset{(-)}{d}(x)$ being the up and down quark (antiquark) distribution functions. The reader is referred to some excellent reviews of the basics of the quark-parton model for the derivation of Eqs. (4.1a)-(4.1d). From Eqs. (4.1a) – (4.1d) we obtain two "master formulae", interms of $\alpha,...,\delta$ or $\varepsilon_L(u),..., \varepsilon_R(d)$ (ignoring strange quark corrections) (Rajasekaran and Sarma, 1974; Hung and Sakurai, 1976):

$$\frac{\sigma_{NC}(\nu N) + \sigma_{NC}(\bar{\nu}N)}{\sigma_{CC}(\nu N) + \sigma_{CC}(\bar{\nu}N)} = \frac{1}{4}(\alpha^2 + \beta^2 + \gamma^2 + \delta^2)$$

$$= |\varepsilon_L(u)|^2 + |\varepsilon_L(d)|^2 + |\varepsilon_R(u)|^2 + |\varepsilon_R(d)|^2,$$

$$(4.3a)$$

$$\frac{\sigma_{NC}(\nu N) - \sigma_{NC}(\bar{\nu}N)}{\sigma_{CC}(\nu N) - \sigma_{CC}(\bar{\nu}N)} = \frac{1}{2}(\alpha\beta + \gamma\delta)$$

$$= \left|\epsilon_L(u)\right|^2 + \left|\epsilon_L(d)\right|^2 - \left|\epsilon_R(u)\right|^2 - \left|\epsilon_R(d)\right|^2$$

$$(4.3b)$$

Equation (4.3a) measures the <u>overall strength</u> of the neutral-current interactions relative to that of the charged-current counterparts. Equation (4.3b) measures the VA interference of neutral-current interactions. Notice that Eqs. (4.3a) and (4.3b) remain invariant under the following interchanges: $\alpha \leftrightarrow \beta, \gamma \leftrightarrow \delta$, reflecting the VA ambiguity mentioned earlier. We also get an invariance under a sign reflection of all four couplings. Obviously with only two equations for four independent couplings, we cannot go very far. Nonetheless, we can still use Eqs. (4.3a) and (4.3b) to determine whether or not the neutral-current interactions are pure V-A or V+A or a combination of both (in addition to the y-measurement). With our definition of $1 + \gamma_5$ for V-A and $1 - \gamma_5$ for V+A, we can see from Eqs. (4.3a) - (4.3b) that the neutral-current interactions are <u>pure</u> V-A when $\epsilon_R(u) = \epsilon_R(d) = 0$, i.e. $\alpha = \beta, \gamma = \delta$; <u>pure</u> V+A when $\epsilon_L(u) = \epsilon_L(d) = 0$, i.e. $\alpha = -\beta, \gamma = -\delta$; and pure V or A when $\left|\epsilon_L(u)\right| = \left|\epsilon_R(u)\right|, \left|\epsilon_L(d)\right| = \left|\epsilon_R(d)\right|$.

The most recent and accurate data on deep-inelastic scattering on isoscalar targets are from the CDHS (Geweniger, 1979) and CHARM (Jonker et al., 1981) groups which report

$$R_\nu = \left(\frac{NC}{CC}\right)_\nu = 0.307 \pm 0.008, \quad R_{\bar{\nu}} \equiv \left(\frac{NC}{CC}\right)_{\bar{\nu}} = 0.373 \pm 0.025, \quad (4.4)$$

for CDHS,

$$R_\nu = 0.320 \pm 0.010, \quad R_{\bar\nu} \equiv 0.377 \pm 0.020 \text{ for CHARM} \qquad (4.5)$$

$$r \equiv \frac{\sigma_{\bar\nu}^{CC}}{\sigma_\nu^{CC}} = 0.491 \pm 0.019 . \qquad (4.6)$$

From the master formulae, we can derive

$$\left|\varepsilon_L(u)\right|^2 + \left|\varepsilon_L(d)\right|^2 = \frac{R_\nu - r^2 R_{\bar\nu}}{1 - r^2} + \text{(corrections)}, \qquad (4.7a)$$

$$\left|\varepsilon_R(u)\right|^2 + \left|\varepsilon_L(d)\right|^2 = \frac{r\left(R_\nu - R_{\bar\nu}\right)}{1 - r^2} + \text{(corrections)}, \qquad (4.7b)$$

where the corrections are small arising from theoretical uncertainties
in the strange quark effects (and to a much smaller extent, the charmed
quark effects). Using e.g. the CDHS data, taking into account correc-
tion for neutron excess in a Fe target, one obtains

$$\left|\varepsilon_L(u)\right|^2 + \left|\varepsilon_L(d)\right|^2 = 0.300 \pm 0.015, \qquad (4.8a)$$

$$\left|\varepsilon_R(u)\right|^2 + \left|\varepsilon_R(d)\right|^2 = 0.024 \pm 0.008, \qquad (4.8b)$$

where the errors include the corrections mentioned above. Similar
numbers can be obtained using data of other groups: Gargamelle, HPWF,
Caltech-Fermilab, BEBC, ... We also show in Fig. 4.1 the famous
"Weinberg nose" which is a function of $\sin^2\theta_W$ in the R_ν and $R_{\bar\nu}$ plane.

Notice that the master formula, Eq. (4.3b), reduces to

$$\frac{\sigma_{NC}(\nu N) - \sigma_{NC}(\bar{\nu}N)}{\sigma_{CC}(\nu N) - \sigma_{CC}(\bar{\nu}N)} = \frac{1}{2}(1 - 2\sin^2\theta_W), \qquad (4.3c)$$

in the standard model ($\delta = 0$). Eq. (4.3c) is known as the Paschos-Wolfenstein (1973) relation. The cancellation of scaling violation effects is especially true for Eq. (4.3c) and the use of partial cross-sections on the left-hand side of the equation may be justified. The CHARM data yield, with the help of Eq. (4.3c), the following value of $\sin^2\theta_W$:

$$\sin^2\theta_W = 0.230 \pm 0.023 . \qquad (4.8c)$$

It is clear from Eqs. (4.8a) and (4.8b) that the neutral-current interactions are predominantly V-A but not pure V-A since the possibility that $\epsilon_R(u)$ and $\epsilon_R(d)$ are both zero is ruled out. Futhermore, the possibility that hadronic neutral-current interactions are either pure V or A or V+A is ruled out by many standard deviations. Within the framework of the standard SU(2) $\otimes$ U(1) model this feature can easily be seen by recalling that here the neutral current is $J_\lambda^{N.C.} = J_\lambda^3 - \sin^2\theta_W J_\lambda^{e.m.}$, with J_λ^3 being pure V - A. Being predominantly V - A means that $\sin^2\theta_W$ is small. This is one of the first successful tests of the standard model. In summary, the hadronic neutral current is neither parity pure nor chirality pure.

4.2 Isospin Structure of Hadronic Neutral Currents

So far we have been able to obtain only $|\varepsilon_L(u)|^2 + |\varepsilon_L(d)|^2$ and $|\varepsilon_R(u)|^2 + |\varepsilon_R(d)|^2$ by using deep-inelastic scattering data on _isoscalar_ targets. To obtain $|\varepsilon_L(u)|$, $|\varepsilon_L(d)|$, $|\varepsilon_L(u)|$ and $|\varepsilon_R(d)|$ separately, we need two more equations in addition to Equations (4.3a) and (4.3b). The two extra equations are needed in order to determine the relative amounts of isoscalar and isovector, or the strength of the d quark interaction compared to that of the u quark. Is the hadronic neutral current pure isovector or pure isoscalar? We can obtain a definite answer to that question by considering

$$\frac{\sigma(\nu p \to \nu \pi^o p) - \sigma(\nu n \to \nu \pi^o n)}{\sigma(\nu p \to \nu \pi^o p)} = 0.40 \pm 0.20 \quad \text{(GGM)}, \qquad (4.9\text{a})$$

$$\frac{\sigma(\nu n \to \nu n \pi^+ \pi^-)}{\sigma(\nu p \to \nu p \pi^+ \pi^-)} = 0.51 \pm 0.10 \quad \text{(BNL 7-ft. chamber)}. \qquad (4.9\text{b})$$

The ratios (4.9a) (Krenz et al., 1978) and (4.9b) (Samios, private communication) should be zero and one respectively if the current is isospin pure. So there is a definite (greater than two standard deviations) evidence for isoscalar-isovector interference from these exclusive reactions, but it is difficult to constrain the coupling parameters short of reliable models of single and double pion production.

4.3   Deep Inelastic Scattering on Proton and Neutron Targets

To study the isospin dependence of the hadronic neutral current, we have to resort to methods other than deep-inelastic scattering on isoscalar targets. One such method is deep-inelastic scattering on proton and neutron targets separately (Hung and Sakurai, 1976). It can be shown that the two extra equations which measure isovector-isoscalar interferences are given by

$$\frac{(\sigma_{NC}^{\nu p} + \sigma_{NC}^{\bar{\nu}p} - \sigma_{NC}^{\nu n} - \sigma_{NC}^{\bar{\nu}n})}{(\sigma_{CC}^{\nu p} + \sigma_{CC}^{\bar{\nu}n} - \sigma_{CC}^{\nu n} - \sigma_{CC}^{\bar{\nu}p})} = -\frac{1}{2}(\alpha\gamma + \beta\delta)$$

$$= -[|\varepsilon_L(u)|^2 - |\varepsilon_L(d)|^2 + |\varepsilon_R(u)|^2$$

$$- |\varepsilon_R(d)|^2], \qquad (4.10a)$$

$$\frac{(\sigma_{NC}^{\nu p} - \sigma_{NC}^{\bar{\nu}p}) - (\sigma_{NC}^{\nu n} - \sigma_{NC}^{\bar{\nu}n})}{(\sigma_{CC}^{\nu p} - \sigma_{CC}^{\bar{\nu}n}) - (\sigma_{CC}^{\nu n} - \sigma_{CC}^{\bar{\nu}p})} = -\frac{1}{2}(\alpha\delta + \beta\gamma)$$

$$= -[|\varepsilon_L(u)|^2 - |\varepsilon_L(d)|^2 - |\varepsilon_R(u)|^2$$

$$+ |\varepsilon_R(d)|^2], \qquad (4.10b)$$

recalling that α, β, γ and δ stand for isovector-vector, isovector axial-vector, isoscalar-vector and isoscalar axial-vector respectively. Equation (4.10a) measures isovector-isoscalar interference and Eq. (4.10b) measures both isoscalar-isovector and V-A interferences.

The most recent data are expressed in terms of

$$R_{\nu p} \equiv \frac{\sigma_{NC}^{\nu p}}{\sigma_{CC}^{\nu p}} \;, \quad R_{\bar{\nu} p} \equiv \frac{\sigma_{NC}^{\bar{\nu} p}}{\sigma_{CC}^{\bar{\nu} p}} \;, \quad R_{\nu}^{n/p} \equiv \frac{\sigma_{NC}^{\nu n}}{\sigma_{NC}^{\nu p}} \;, \quad R_{\bar{\nu}}^{n/p} \equiv \frac{\sigma_{NC}^{\bar{\nu} n}}{\sigma_{NC}^{\bar{\nu} p}} \;.$$ The experimental

values are

$$R_{\nu p} = \begin{cases} 0.52 \pm 0.04 & \text{(Blietschau et al., 1979)} \\ 0.48 \pm 0.17 & \text{(Harris et al., 1977)} \end{cases} , \tag{4.11a}$$

$$R_{\bar{\nu} p} = 0.42 \pm 0.13 \quad \text{(Derrick et al., 1978)} \tag{4.11b}$$

$$R_{\nu}^{n/p} = 1.22 \pm 0.35 \;, \quad \text{(Marriner, 1977)} \tag{4.11c}$$

$$R_{\bar{\nu}}^{n/p} = 0.64 \pm 0.18 \;. \quad \text{(Roe, 1979)} \tag{4.11d}$$

Equations (4.11c) and (4.11d) show that the hadronic neutral current is _not_ isospin pure since $R_{\nu}^{n/p}$ and $R_{\bar{\nu}}^{n/p}$ should both be unity in that case.

To make good use of Eqs. (4.11a) and (4.11b), we use the following equations which are easy to derive,

$$\left| \epsilon_L(u) \right|^2 + \eta \left| \epsilon_L(d) \right|^2 = \frac{9}{8} \left(\eta R_{\nu p} - \frac{1}{9} R_{\bar{\nu} p} \right) , \qquad (4.12a)$$

$$\left| \epsilon_R(u) \right|^2 + \eta \left| \epsilon_R(d) \right|^2 = \frac{3}{8} \left(\eta R_{\bar{\nu} p} - R_{\nu p} \right) , \qquad (4.12b)$$

where

$$\eta = \frac{\int_0^1 x d(x) \, dx}{\int_0^1 x u(x) \, dx} \approx 0.51 , \qquad (4.13)$$

and where the sea of quark-antiquark is ignored. Equations (4.12a) and (4.12b) can be used in conjuction with Eqs. (4.3a) and (4.3b) which determine $\left| \epsilon_L(u) \right|^2 + \left| \epsilon_L(d) \right|^2$ and $\left| \epsilon_R(u) \right|^2 + \left| \epsilon_R(d) \right|^2$ respectively, to obtain $\left| \epsilon_L(u) \right|^2$, $\left| \epsilon_L(d) \right|^2$, $\left| \epsilon_R(u) \right|^2$ and $\left| \epsilon_R(d) \right|^2$ separately.

Unfortunately, $\left| \epsilon_R(u) \right|^2 + \left| \epsilon_R(d) \right|^2 = 0.024 \pm 0.08$ is too small to be a useful equation. We can therefore only use $\left| \epsilon_L(u) \right|^2 + \left| \epsilon_L(d) \right|^2$ and

$|\varepsilon_L(u)|^2 + \eta |\varepsilon_L(d)|^2$ and plot the experimental data in a $|\varepsilon_L(u)|^2$ –

$|\varepsilon_L(d)|^2$ plane as shown in Fig. 4.2 . The intersection of the

$|\varepsilon_L(u)|^2 + |\varepsilon_L(d)|^2$ and $|\varepsilon_L(u)|^2 + \eta |\varepsilon_L(d)|^2$ bands yields (Dydak,

1979)

$$|\varepsilon_L(u)|^2 = 0.15 \pm 0.05 , \qquad (4.14a)$$

$$|\varepsilon_L(d)|^2 = 0.16 \pm 0.07 . \qquad (4.14b)$$

Also shown in the $|\varepsilon_L(u)|^2 - |\varepsilon_L(d)|^2$ plane are the allowed region

for $|\varepsilon_L(u)|^2$ and $|\varepsilon_L(d)|^2$ obtained from a second method, the semi-

inclusive reaction $\overset{(-)}{\nu} + N \to \overset{(-)}{\nu} + \pi^{\pm} +$ any, which historically was

the first to obtain <u>all four</u> chiral couplings. The agreement between

the two independent methods is quite remarkable.

4.4   Semi-inclusive Pion Production on Isoscalar Targets

The reactions involved in the second method are the semi-

inclusive neutral-current pion production on isoscalar targets:

$$\overset{(-)}{\nu} + N \to \overset{(-)}{\nu} + \pi^{\pm} + \text{any} ,$$

where we make use of the quark fragmentation model. How can we extract

isospin informations from this set of reactions? Experimentally, it

is known from charged-current reactions that in $\nu N(\bar\nu N)$ collisions the

probability for a u(d) quark to fragment into a $\pi^{+(-)}$ in the forward

direction is much larger than the probability for a u(d) quark to

fragment into a $\pi^{-(+)}$, i.e,

$$\left(\frac{\pi^+}{\pi^-}\right)_{\nu \to \mu^-} = \frac{D_u^{\pi^+}}{D_u^{\pi^-}} \simeq 3, \qquad (4.15a)$$

$$\left(\frac{\pi^+}{\pi^-}\right)_{\bar{\nu} \to \mu^+} = \frac{D_d^{\pi^+}}{D_d^{\pi^-}} = \frac{D_u^{\pi^-}}{D_u^{\pi^+}} \simeq \frac{1}{3}, \qquad (4.15b)$$

for the kinematical region $z \equiv \dfrac{E_\pi}{E_{had}} \simeq 0.6$ (fraction of total hadronic energy carried by the detected pion). The kinematical region concerned is the "current fragmentation region". To get a feeling for what goes on, recall that the quark beams produced in charged and neutral current reactions come from the following elementary processes (ν-scattering on valence quarks only): $\nu_\mu + d \to \mu^- + u$, $\bar{\nu}_\mu + u \to \mu^+ + d$, $\overset{(-)}{\nu}_\mu + \begin{Bmatrix} u \\ d \end{Bmatrix} \to \overset{(-)}{\nu}_\mu + \begin{Bmatrix} u \\ d \end{Bmatrix}$. For $\overset{(-)}{\nu} \to \mu^{-(+)}$, we get a u-quark (d-quark) beam which would then fragment into a pion characterized by a probability amplitude $D_{u(d)}^\pi(z)$. For neutral-current reactions, both quarks helicities are involved and the right-handed quarks are only one-third as effective as the left-handed ones (recall the factor $\int_0^1 (1-y)^2 dy = \frac{1}{3}$).

In neutral-current reactions on isoscalar targets where there are as many u quarks as d quarks, an asymmetry between π^- and π^+ would give informations on the difference in the magnitude of neutral-current interactions between u and d quarks. Quantitatively, we have (Hung, 1977; Sehgal, 1977; Okada and Pakvasa, 1976)

$$\left(\frac{\pi^+}{\pi^-}\right)_{\nu\to\nu} = \frac{\left[\left|\epsilon_L(u)\right|^2 + \frac{1}{3}\left|\epsilon_R(u)\right|^2\right] D_u^{\pi^+} + \left[\left|\epsilon_L(d)\right|^2 + \frac{1}{3}\left|\epsilon_R(d)\right|^2\right] D_u^{\pi^-}}{\left[\left|\epsilon_L(u)\right|^2 + \frac{1}{3}\left|\epsilon_R(u)\right|^2\right] D_u^{\pi^-} + \left[\left|\epsilon_L(d)\right|^2 + \frac{1}{3}\left|\epsilon_R(d)\right|^2\right] D_u^{\pi^+}},$$

$$(4.16a)$$

$$\left(\frac{\pi^+}{\pi^-}\right)_{\bar\nu\to\bar\nu} = \left\{ \begin{array}{c} \epsilon_L(u) \leftrightarrow \epsilon_R(u) \\ \epsilon_L(d) \leftrightarrow \epsilon_R(d) \end{array} \right\} , \qquad (4.16b)$$

where $D_d^{\pi^+} = D_u^{\pi^-}$ and $D_d^{\pi^-} = D_u^{\pi^+}$ by strong isospin invariance. Equations (4.16a) and (4.16b) are the two extra equations needed in conjuction with Eqs. (4.3a) and (4.3b), to determine completely $\left|\epsilon_L(u)\right|^2$, $\left|\epsilon_L(d)\right|^2$, $\left|\epsilon_R(u)\right|^2$ and $\left|\epsilon_R(d)\right|^2$.

Experimentally, we have (Kluttig et al., 1977)

$$\left(\frac{\pi^+}{\pi^-}\right)_{\nu\to\nu} = 0.77 \pm 0.14 , \qquad (4.17a)$$

$$\left(\frac{\pi^+}{\pi^-}\right)_{\bar\nu\to\bar\nu} = 1.64 \pm 0.36 . \qquad (4.17b)$$

Using the experimental values (4.19a-b) and Eqs. (4.16a), (4.16b) (4.3a) and (4.3b), Sehgal (1977) obtained

$$\left|\epsilon_L(u)\right|^2 = 0.11 \pm 0.03 \qquad (4.18a)$$

$$\left|\epsilon_L(d)\right|^2 = 0.19 \pm 0.03 , \qquad (4.18b)$$

$$\left|\epsilon_R(u)\right|^2 = 0.03 \pm 0.015, \qquad (4.18c)$$

$$\left|\varepsilon_R(d)\right|^2 = 0.00 \pm 0.015 \ . \tag{4.18d}$$

Sehgal's analysis was subsequently confirmed by semi-inclusive data at higher energies coming from the BEBC and FMMS collaborations (Roe, 1979; Deden, 1979) and by deep-inelastic scattering data on protons (Dydak, 1979) discussed in the last subsection as can be seen in Fig. 4.2 .

There is an unavoidable feature which is that, because of the incoherence assumption of the quark-parton model, only $\left|\varepsilon_{L,R}\right|^2$ is determined and not $\varepsilon_{L,R}$. Such a feature gives rise to a sign ambiguity in $\varepsilon_{L,R}$. We have therefore a $2^4 = 16$ fold ambiguity. Fortunately, $\left|\varepsilon_R(d)\right|^2$ happens to be essentially zero, so there is only an eight-fold ambiguity (Hung and Sakurai, 1977 b). The eight solutions shown in Fig. 4.3, are denoted by

$A(A')$: isovector axial-vector dominant: $|\beta|$ large

$B(B')$: isovector vector dominant: $|\alpha|$ large

$C(C')$: isoscalar axial-vector dominant: $|\delta|$ large

$D(D')$: isoscalar vector dominant: $|\gamma|$ large

where the primed solutions are obtained by simultaneous sign reversal of all four constants. The ambiguity between the unprimed and primed solutions is shown in Chapter VII to be resolved by assuming a property called "factorization". Therefore we now concentrate on the resolution of ambiguities among the unprimed solutions.

4.5   Exclusive Processes

To be able to resolve the above ambiguities, we have to resort to exclusive reactions.   There are three classes of experiments for that purpose (i.e. those which have been performed).

i) Experiments which look for $\Delta(1236)$ resonance i.e. $\nu + N \rightarrow \nu + \Delta$, since a clear and strong signal would indicate a dominance of the isovector solution (solutions A and B).   This is precisely what happened since a strong signal has been observed in $\nu + N \rightarrow \nu + \Delta$ by the Gargamelle collaboration (see Fig.   4.4).   Qualitatively speaking, the clear observation of a Δ-resonance in neutral-current reactions indicates that there is a large isovector component in the hadronic neutral current.   This clearly rules out solutions C and D which predict the Δ-production cross section to be an order-of-magnitude smaller (Ecker, 1978; Abbott and Barnett, 1978 a and b; Monsay, 1978).

ii)   Elastic scattering experiments $\nu + p \rightarrow \nu + p$ and $\bar{\nu} + p \rightarrow \bar{\nu} + p$ were of great help in the resolution of solution ambiguities.   The matrix element of the hadronic neutral current between proton states is characterized by

$$<P(p') | J_\lambda^{N.C.} | P(p) > = i\bar{u}_p(p') \left[\gamma_\lambda F_1^N(q^2) + i\sigma_{\lambda\rho} q_\rho \frac{F_2^N(q^2)}{2M_N} + \gamma_\lambda \gamma_5 G_A(q^2) \right] u_p(p),$$

$$(4.19)$$

where $q_\lambda = (p' - p)_\lambda$.   It is convenient to use instead the following form factors: $G_E(q^2)$, $G_M(q^2)$ and $G_A(q^2)$, which at $q^2 = 0$ have the following clean prediction

$$G_E(0) = F_1^N(0) = \frac{1}{2} (\alpha + 3\gamma) , \qquad (4.20a)$$

$$G_M(0) = F_1^N(0) + F_2^N(0) = \frac{4.7}{2} (\alpha + 0.56\gamma) , \qquad (4.20b)$$

$$G_A(0) = \frac{1.25}{2} (\beta + \lambda\delta) , \qquad (4.20c)$$

where λ is a model-dependent parameter. We see that only vector

currents contribute to G_E and G_M while only axial-vector currents

contribute to G_A.

In Eq. (4.20a-c), the factor of 3 comes from the definition of

the isovector and isoscalar currents which are $V_\lambda^{(3)} = \frac{1}{2}(\bar{u}\gamma_\lambda u - \bar{d}\gamma_\lambda d)$

and $V_\lambda^{(0)} = \frac{1}{3}(\bar{u}\gamma_\lambda u + \bar{d}\gamma_\lambda d)$ respectively. The factor 4.7 comes from

the isovector part of nucleon magnetic moments i.e.

$4.7 = 1 + \mu_p - \mu_n$ where $\mu_p = 1.79$, $\mu_n = -1.91$ are the proton and neutron

anomalous magnetic moments respectively. The factor 0.56 comes from

the isoscalar-isovector ratio $(1 + \mu_p + \mu_n)/(1 + \mu_p - \mu_n)$. Also

$1.25 = F + D = g_A/g_V$ comes from the isovector part of the axial-vector

current. The factor $\lambda = g_A^{(S)}/g_A$ is the ratio of isoscalar to isovector

axial-vector constants. In the SU(6) non-relativistic quark model,

we have $\lambda = 3/5 = 0.6$ ($\bar{s}s$ contribution is ignored). Another approach

uses SU(3) symmetry (again ignoring $\bar{s}s$ contribution) giving

$\lambda = (3F - D)/(F + D) = 0.40 \pm 0.04$, where $D/(D + F) = 0.65 \pm 0.01$ is

obtained from semi-leptonic hyperon decay data (consult Hung (1978)

for more details). Because of the uncertainties in λ(Wolfenstein,

1979) sometimes it is preferrable to leave it as a free parameter.

Unfortunately the data is available only for $q^2 \geqslant 0.4$ GeV^2
(for $q^2 \rightarrow 0$ the kinetic energy of the proton becomes zero and it cannot
be seen, so we would have a final state with only invisible particles,
$\overset{(-)}{\nu}$ and p). In one approach (Hung and Sakurai, 1977b), by guessing
the q^2-dependence of the form factors G_E, G_M and G_A^S using CVC (by
comparison with ep and en elastic data) and the charged-current
reaction $\nu_\mu \rightarrow n \rightarrow \mu^- + p$, the integrated cross-sections are used. It
was found, using the data available, that solutions C(C') and D(D')
are ruled out in agreement with the semi-quantitative assessment of
the Δ-production. We are left with solutions A(A') and B(B').

The second approach exploits the q^2-dependence of $d\sigma/dq^2$. By
extrapolating back to $q^2 = 0$, the following values are obtained

$$G_E = 0.5 \begin{smallmatrix} + & 0.25 \\ - & 0.5 \end{smallmatrix},$$
(4.21a)

$$G_M = 1.0 \begin{smallmatrix} + & 0.7 \\ - & 0.8 \end{smallmatrix},$$
(4.21b)

$$G_A^S = 0.5 \begin{smallmatrix} + & 0.4 \\ - & 0.3 \end{smallmatrix},$$
(4.21c)

to be compared with 0, 1.1, 0.6 for solution A and 0.55, 2.2, 0.26
for solution B (Paschos, 1979; Claudsen et al., 1979). Clearly
solution A is favored while solution B is ruled out. Solution D(D')
is marginal.

(iii) The reaction $\bar{\nu}_e + D \rightarrow \bar{\nu}_e + n + p$ at reactor energies
studied by a UC-Irvine group (Pasierb et al., 1979) is sensitive only
to the isovector axial- vector coupling β (at threshold energy

$E_\nu = 2.25$ MeV) – a transition from $^3S_1$(I =0) D state to $^1S_0$(I = 1) n–p

state. This provides a direct measurement of β. Experimentally,

$$|\beta| = 0.9 \pm 0.1 . \qquad (4.22)$$

The values of $|\beta|$ for solutions A(A') and B(B') are 0.92 ± 0.14 and

0.58 ± 0.14. The experimental data is in excellent agreement with

solution A(A') while solution B(B') is clearly ruled out (solutions

C(C') and D(D') are ruled out even more strongly). The results from

exclusive reactions are summarized in table 4.1. Clearly solution

A(A') is the sole survivor. Note that there exists high degrees of

redundancy, a very desirable feature in weak interaction physics.

Also, if you distrust semi-inclusive results, then Δ production and

νp elastic scattering alone select region A but with somewhat larger

errors (Claudsen et al., 1979).

The overall sign of the neutrino-quark coupling constants

(solution A vs solution A') could be determined by interference with

gravity, which is obviously quite a difficult task. By using

factorization and the SLAC asymmetry data, we shall see later in

Chapter VII that it is possible to eliminate solution A'.

The most ambitious fit to all neutrino-hadron data has recently

been performed by Langacker et al. (1979) and reviewed in details by

Kim et al. (1981) (see also Abbott and Barnett, 1978 b). They obtained

a unique solution that is essentially the same as solution A of Hung

and Sakurai (see Fig. 4.5). The results are (with χ^2/DOF = 13.9/24):

| FIT TO DATA | STANDARD MODEL |
| | ($\sin^2 \theta_W = 0.23$) |

$$\epsilon_L(u) = 0.340 \pm 0.033 \qquad\qquad 0.347$$

$$\epsilon_L(d) = -0.424 \pm 0.026 \qquad\qquad -0.423$$

$$(4.23a)$$

$$\epsilon_R(u) = -0.179 \pm 0.019 \qquad\qquad -0.153$$

$$\epsilon_R(d) = -0.017 \pm 0.058 \qquad\qquad 0.077$$

or

$$\alpha = 0.589 \pm 0.067 \qquad\qquad 0.540$$

$$\beta = 0.937 \pm 0.062 \qquad\qquad 1.000$$

$$(4.23b)$$

$$\gamma = -0.273 \pm 0.081 \qquad\qquad -0.153$$

$$\delta = 0.101 \pm 0.093 \qquad\qquad 0.000$$

So we see that the phenomenologically determined parameters are in excellent agreement with the standard model predictions with $\sin^2 \theta_W$ set equal to 0.23. Further tests of $SU(2) \otimes U(1)$ are discussed in Chapter VII.

V. ELECTRON-QUARK INTERACTIONS

5.1 General Remarks on Weak-Electromagnetic Interference

So far we have been concerned with the neutral-current inter-
actions initiated by neutrino beams. Our understanding of the neutral-
current interactions is necessarily incomplete if we restrict ourselves
only to neutrino-induced reactions. For one thing those reactions do
not even answer the vital question of whether or not parity is
conserved in the weak neutral-current interactions. This is because
no pseudoscalar observables are directly detected. Furthermore a
study of parity violation is complicated by the fact that we are
forced to start with neutrinos with only one kind of handedness-
left-handed for $\nu_{\mu,e}$, right-handed for $\bar{\nu}_{\mu,e}$. Indeed it is not
difficult to show that all neutrino-induced neutral-current data
that fit the (parity-violating) standard model can be fitted equally
well using a two-Z boson model of the VV + AA (parity-conserving)
type (Fritzsch and Minkowski, 1976).

In this chapter we turn our attention to the interactions of
electrons with hadrons (ultimately with quarks) as revealed in
inelastic electron-deuteron scattering and in atomic radiative
transitions. The neutral-current interactions involving charged
leptons (e,μ,τ) only will be the subject of the next chapter.

There are two classes of lepton-quark neutral-current interactions-
(a) parity- conserving $V_{lept}V_{quark}$ and $A_{lept}A_{quark}$, and (b) parity-
violating $A_{lept}V_{quark}$ and $V_{lept}A_{quark}$. Until we explore very high

energies with q^2 of order 10^4 GeV^2, the parity-conserving neutral-current interactions are completely masked by the much stronger electromagnetic interactions. On the other hand, the parity-violating neutral-current interactions could be detected by their characteristic signature - the dependence of measurable quantities on pseudoscalars, e.g. electron or photon helicity. Such parity-violating effects arise from weak-electromagnetic interference and are of first order in Fermi's G.

The first proposal to detect weak neutral-current interactions between the electron and the nucleon via weak-electromagnetic interference was made more than twenty years ago by Zeldovich (1959). In a paper exhibiting remarkable foresight, he correctly pointed out that a parity-violating asymmetry as large as 10^{-3} to 10^{-4} could be observable in the elastic scattering of polarized electrons by protons at $q^2 \simeq m_p^2$ provided the weak Lagrangian contains a parity-violating neutral-current piece with strength comparable to the usual charged-current interactions. His argument is very simple; a parity-violating asymmetry arising from weak-electromagnetic interference must go roughly as the ratio of Fermi's G to the one-photon exchange amplitude e^2/q^2,

$$A_{e.m.}A_{weak}/\left[|A_{e.m.}|^2 + |A_{weak}|^2\right] \simeq A_{weak}/A_{e.m.} \sim G/(e^2/q^2)$$

$$(5.1)$$

See Fig.5.1. If we put the numbers in, we get

$$G/(e^2/q^2) \simeq (10^{-5} \times 137/4\pi)q^2/m_p^2$$

$$\simeq 10^{-4}q^2/m_p^2. \qquad (5.2)$$

Zeldovich's expectation was based on an SU(2) symmetric V-A model of the kind considered somewhat earlier by Bludman (1958), which gives essentially the same predictions at low energies as the standard model with $\sin^2\theta_W$ set to zero. It is amusing that the parity-violating asymmetry experimentally measured at SLAC twenty years later is indeed of order 10^{-4} at $q^2 \sim m_p^2$, as we now discuss.

5.2   Inelastic Electron-Deuteron Scattering

Suppose we start with a beam of longitudinally polarized electrons. They may be scattered inelastically by a nucleus. Parity nonconservation due to the weak neutral-current interactions can be detected by studying how the observed cross section depends on the incident electron helicity. If it is found that

$$\sigma(\lambda = 1/2) \neq \sigma(\lambda = -1/2) \qquad (5.3)$$

where λ stands for the incident electron helicity in

$$e^-_{polarized} + nucleus \to e^- + any, \qquad (5.4)$$

we have unambiguous proof for parity violation, independently of any model.

A beautiful experiment performed at SLAC by a SLAC-Yale collaboration in 1978-1979 studied a reaction of this kind (Prescott et al. 1978, 1979). A 16-22 GeV/c beam of $\sim$ 40% polarized electrons was incident on a liquid deuterium target. Scattered electrons in the 11-16 GeV/c range were detected in a "single arm" arrangement with a spectrometer set at 4^0 in the laboratory system.

Before presenting the experimental results, we discuss in some detail what information can be obtained on the neutral-current coupling parameters by performing such an experiment. It has become customary to define the parity-violating asymmetry A as follows.

$$A = [\sigma(\lambda = \tfrac{1}{2}) - \sigma(\lambda = -\tfrac{1}{2})]/[\sigma(\lambda = \tfrac{1}{2}) + \sigma(\lambda = -\tfrac{1}{2})] \qquad (5.5)$$

where $\sigma(\lambda = \pm \tfrac{1}{2})$ is the double differential cross section $d^2\sigma/d\Omega dE'$ for right-(left-) handed electron on deuteron. In terms of the scaling variables x and y defined in Section 4.1, we can derive

$$A(x,y,q^2)/q^2 = a_1(x) + a_2(x)\left\{[1 - (1-y)^2]/[1 + (1-y)^2]\right\}$$

$$(5.6)$$

where the Bjorken scaling hypothesis and $\sigma_S/\sigma_T \simeq 0$, i.e. no longitudinal (spin-zero parton like) contribution to electroproduction, have been assumed. The function a_1 (x) depends on $A_{lept}V_{quark}(\tilde{\alpha}$ and $\tilde{\gamma})$ while $a_2(x)$ depends on $V_{lept}A_{quark}(\tilde{\beta}$ and $\tilde{\delta})$. By changing y, the ratio of the energy transfer to the incident energy, it is possible to determine separately $a_1(x)$ and $a_2(x)$. For isoscalar targets such as deuterium, a_1 and a_2 are actually independent of x:

$$a_1 = (G/\sqrt{2} \ e^2)(9\tilde{\alpha} + 3\tilde{\gamma}/5), \qquad (5.7a)$$

$$a_2 = (G/\sqrt{2} \ e^2)(9\tilde{\beta} + 3\tilde{\delta}/5) \qquad (5.7b)$$

where we may recall from Section 2.1 the definitions of the four coupling constants:

$$
\begin{aligned}
\tilde{\alpha} &: \quad A_{\text{lept}} V_{\text{quark}}^{I=1}, \\
\tilde{\beta} &: \quad V_{\text{lept}} A_{\text{quark}}^{I=1}, \\
\tilde{\gamma} &: \quad A_{\text{lept}} V_{\text{quark}}^{I=0} \\
\tilde{\delta} &: \quad V_{\text{lept}} A_{\text{quark}}^{I=0} .
\end{aligned}
\qquad (5.8)
$$

Equations (5.7) is a model-independent way of writing down expressions obtained by Cahn and Gilman (1978), Yoshimura (1978) and others. Notice that the asymmetry goes like Gq^2, in agreement with (5.2).

We now briefly outline the derivation of (5.6) and (5.7), first within the framework of the quark-parton model. Suppose a right-handed electron is incident; the incoherent contribution to inelastic electron scattering due to quark i of charge z_i can then be written as

$$\sigma^{(i)}(\lambda = \tfrac{1}{2}) = \left| - z_i e^2/q^2 + (4G/\sqrt{2}) \ \varepsilon \ (e_R, q_L^i) \right|^2 (1 - y)^2$$

$$+ \left| - z_i e^2/q^2 + (4G/\sqrt{2}) \varepsilon (e_R, q_R^i) \right|^2 . \qquad (5.9)$$

Here $\varepsilon(e_R, q_{L,R}^i)$ are appropriate four-Fermion neutral-current

coupling constants in chiral (rather than parity-eigenstate) notation;

they can be written as linear combinations of $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}$ and $\tilde{\delta}$ and their

parity-conserving analogs. A similar expression can be obtained for

the $\lambda = -\frac{1}{2}$ cross section. After summing over the u and d quark

contributious and taking the difference between $\sigma(\lambda = \frac{1}{2})$ and

$\alpha(\lambda = -\frac{1}{2})$, the terms that depend on the parity-conserving parameters

drop out, as they must, and we are left with (5.6) and (5.7). The

independence of a_1 and a_2 on x can be traced to the fact that the d

and u quark distributions are identical for isoscalar targets.

We may ask to what extent these formulas depend on the details

of quark parton concepts. Fortunately they also follow from rather

general considerations (Bjorken 1978b, Wolfenstein 1978). First of

all, scaling violation effects of the kind predicted by QCD affect

the weak and electromagnetic interactions in the same way and there-

fore cancel as we are here concerned only with the ratio of the weak

to the electromagnetic amplitude. Second, for the a_1 term where the

vector part of the hadronic current is relevant, we can easily see

that (5.7a) would be an exact relation, depending only on the CVC

hypothesis if there were no isoscalar current contributions. When

the isoscalar pieces are not ignored, the only quark-parton like

idea needed is that the structure function contribution of $\bar{u}u - \bar{d}d$

is the same as that of $\bar{u}u + \bar{d}d$, a rather mild form of the incoherence

assumption inherent in the quark parton model. As for the a_2 term

which arises from $V_{quark}^{(e.m.)} A_{quark}$ interferance, it is possible, again

with the incoherence assumption, to relate it to the difference between the charged-current inelastic neutrino-nucleon scattering cross sections $\sigma_{CC}(\nu N) - \sigma_{CC}(\bar{\nu}N)$; to the extent that the quark model can account for this difference, it is legitimate to apply the model for computing a_2, which, as we will see, is in any case small.

The first SLAC data, which became available in the summer of 1978, was obtained at one value of y(viz. $y \simeq 0.21$) and therefore determined just one linear combination of a_1 and a_2, viz, $a_1 + 0.23a_2$. Subsequently the y dependence has been studied, as shown in Fig. 5.2. A two-parameter fit may be made to separate a_1 and a_2. In this way they have obtained (Prescott et al. 1979)

$$a_1 = (-9.7 \pm 2.6) \times 10^{-5} \text{GeV}^{-2}, \tag{5.10a}$$

$$a_2 = (4.9 \pm 8.1) \times 10^{-5} \text{GeV}^{-2}, \tag{5.10b}$$

or, in terms of the coupling parameters

$$\tilde{\alpha} + \tilde{\gamma}/3 = -0.60 \pm 0.16, \tag{5.11a}$$

$$\tilde{\beta} + \tilde{\delta}/3 = 0.31 \pm 0.51 . \tag{5.11b}$$

The errors in (5.11a) and (5.11b) are highly correlated, as can be seen from Fig. 5.3. The best determined quantity is a particular linear combination:

$$(\widetilde{\alpha} + \widetilde{\gamma}/3) + 0.25 \ (\widetilde{\beta} + \widetilde{\delta}/3) = -0.53 \pm 0.05 \qquad (5.12)$$

The first model-independent statement we can make is that parity violation is established beyond any shadow of doubt. Moreover, because the asymmetry is, in fact, of order Gq^2/e^2, the observed effect cannot be due to higher order charged-current interactions. Models in which parity is conserved in the neutral-current interactions are now only of historical interest. Second, the a_1 term is "large" while the a_2 term is "small" (compatible with zero). This shows that the $A_{lept}V_{quark}$ interactions are much stronger than the $V_{lept}A_{quark}$ interactions, in contradiction with models, e.g. the "hybrid model", in which the A_{lept} contribution varnishes indentically.

Let us now compare these results with standard model predictions. This model relates the four coupling parameters to just one parameter $\sin^2\theta_W$:

$$\widetilde{\alpha} = - \ (1 - 2 \ \sin^2\theta_W),$$

$$\widetilde{\beta} = - \ (1 - 4 \ \sin^2\theta_W),$$

$$\widetilde{\gamma} = \frac{2}{3} \ \sin^2\theta_W,$$

$$\widetilde{\delta} = 0 \qquad (5.13)$$

So a_1 and a_2 are functions of $\sin^2\theta_W$ only. Notice in particular that for $\sin^2\theta_W \simeq \frac{1}{4}$, the a_2 term is predicted to be vanishingly small, in agreement with (5.10b). More generally, a one-parameter fit to all the data points in Fig. 5.2 can be made. The result is

$$\sin^2\theta_W = 0.224 \pm 0.12 \pm 0.008 \tag{5.14}$$

where the first error is statistical and the second, systematic. The fact that this entirely different way of determining $\sin^2\theta_W$ agrees with the determinations of $\sin^2\theta_W$ from the various neutrino-induced reactions discussed in the previous chapters is taken to be an extraordinary triumph of the standard model.

Historically, at the time the electron-deuteron asymmetry measurement was being performed, there were conflicting atomic parity experiments; as a result, a multitude of electroweak models were proposed to accommodate this or that experimental result. The SLAC experiment was most decisive in showing that the standard model was practically the only surviver.

From a purely phenomenological point of view, however, our task is not yet completely over. There are altogether as many as four coupling constants that characterize the parity-violating electron-quark interactions while the SLAC asymmetry experiment succeeded in determining one linear combination of the $A_{lept}V_{quark}$ parameters with a relatively small error and one linear combination of the

$V_{lept}A_{quark}$ parameters with a much larger error. For a complete and separate determination of all four parameters, we need more experiments. Comparison of the ep and eD asymmetries in inelastic electron scattering and studies of parity violation in _elastic_ ep, eD and e nucleus scattering at moderate energies can, in principle, provide us with additional information. At this moment, however, the only experimental information beyond the SLAC eD asymmetry comes from the realm of atomic physics, to which we now turn our attention.

5.3 Parity Violation in Atoms

It was remarked in Section 5.1 that a parity-violating asymmetry due to weak-electromagnetic interference goes like $10^{-4}q^2/m_p^2$. In atomic physics q^2 is of order $1/R_{atom}^2 \sim (m_e/137)^2$ where R_{atom} is a typical atomic dimension. Plugging in the numbers, we expect the asymmetry to be of order 10^{-15}, much too small to be measurable. Fortunately, however, there are methods available to enhance parity-violating asymmetry, as we will see shortly.

Let us first examine how parity violation arises in atomic transitions. If there are parity violating interactions between the electron and the nucleon, atomic levels are expected to contain small opposite-parity components. To estimate the admixture we begin by constructing an effective parity-violating potential between an atomic electron and the nucleus. Parity violating effects due to the interactions among the atomic electrons themselves can be shown to be negligible in comparison.

Nuclei are made up of nucleons, and nucleons are believed to be made up of quarks. So parity-violation experiments in atomic transitions ultimately determine the four coupling parameters $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}$ and $\tilde{\delta}$ of the electron-quark interactions. To atomic physicists, however, the nucleon does look elementary; as a result, the electron-nucleon constants C_{1p}, C_{1n}, C_{2p} and C_{2n} also appear in the literature.

$$\tilde{\alpha} = - (C_{1p} - C_{1n}), \qquad 1.25\tilde{\beta} = - (C_{2p} - C_{2n}),$$

$$\tilde{\gamma} = - \frac{1}{3} (C_{1p} + C_{1n}), \qquad 1.25\lambda)\tilde{\delta} = -(C_{2p} + C_{2n}). \qquad (5.15)$$

The reason for the appearance of the axial-vector factors, 1.25 and 1.25λ was already discussed in connection with elastic νp scattering. (See Section 4.4).

The fact that the nucleons are slow implies that for the A_{lept} V_{quark} interactions ($\tilde{\alpha}$ and $\tilde{\gamma}$) only the time component of V_{quark} is of importance; likewise, for the $V_{lept}A_{quark}$ interactions ($\tilde{\beta}$ and $\tilde{\delta}$), only the space components are important. All this is familiar from nuclear beta decay. However, unlike the beta decay case, there is coherence for the vector part because the "charges" are additive. The matrix element of the time component of V_{quark} goes like the number of nucleons; more precisely it is given by the "weak charge" Q_W:

$$Q_W(Z,N) = - [\tilde{\alpha}(Z-N) + 3\tilde{\gamma}(Z+N)] \qquad (5.16)$$

where Z and N stand for the numbers of protons and neutrons, respecitvely. In the standard model this can be written as

$$Q_W(Z,N)\big|_{standard} = - [Z(4 \sin^2\theta_W - 1) + N], \qquad (5.17)$$

so, with $\sin^2\theta_W \simeq \frac{1}{4}$, Q_W goes roughly like the number of neutrons. In contrast, because the nucleon spins tend to cancel with each other only the nucleons outside the closed shells contribute to the matrix elements of $A_{quark}^{I = 1,0}$; for heavy atoms the $V_{lept}A_{quark}$ interactions are therefore relatively unimportant. As for the lepton side, we note, for example, that the time component of the axial-vector current gives

$$\bar{e}\gamma_4\gamma_5 e \rightarrow - (\vec{p} + \vec{p}') \cdot \vec{\sigma}_e/2m_e \qquad (5.18)$$

in the plane-wave representation. Putting everything together we obtain a short-ranged parity-violating potential between the electron and the nucleus, as derived by Bouchiat and Bouchiat (1974):

$$H_{PV} = (G/4\sqrt{2})Q_W \vec{\sigma}_e \cdot \{\delta^3(\vec{x}), \vec{p}_e/m_e\}$$

$$+ (G/2\sqrt{2})\left[\{\vec{p}_e/m_e, \delta^3(\vec{x})\} - i[\delta^3(\vec{x}), (\vec{\sigma}_e \times \vec{p}_e)/m_e]\right]$$

$$\times 1.25 [(\widetilde{\beta} + \lambda\widetilde{\partial})\vec{s}_p + (-\widetilde{\beta} + \lambda\widetilde{\partial})\vec{s}_n], \qquad (5.19)$$

with Q_W given as in (5.16) and λ defined by (4.20c). Notice that $\vec{S}_p$ and $\vec{S}_n$ stand for the total spin sums of the protons and neutrons respectively. It may be mentioned that the nonrelativistic reduction made for the electron is for illustration only. In actual calculations we must use potentials appropriate for the Dirac electron.

Because of H_{pv}, an atomic level is no longer expected to be a pure eigenstate of parity. Let $|i>$ be a simultaneous eigenstate of energy and parity in the absence of H_{pv}. With H_{pv} added, however, the exact energy eigenstate for this system contains a tiny admixture of opposite-parity components. The standard perturbation theory technique gives

$$|i> \to |i'> = |i> + \sum_{k \neq i} |k><k|H_{pv}|i>/(E_i - E_k) \quad (5.20)$$

where the states $|k>$ are opposite parity states now connected to $|i>$ by virtue of H_{pv}.

Consider now a radiative transition between $|i>$ and $|f>$, or, more precisely, because of parity mixing, between $|i'>$ and $|f'>$. It may, for instance, be an M1 transition when $|i>$ and $|f>$ have the same parity. The transition matrix element is

$$<f'|H_{e.m.}|i'> \simeq <f|H_{e.m.}|i>$$

$$+ \sum_{k \neq i} <f|H_{e.m.}|k><k|H_{pv}|i>/(E_i - E_k)$$

$$+ \sum_{k \neq f} <f|H_{pv}|k><k|H_{e.m.}|i>/(E_f - E_k), \quad (5.21)$$

where we have ignored terms of order G^2. When $|i>$ and $|f>$ are connected by an M1 transition, the first term $<f|H_{e.m.}|i>$ is the usual M1 matrix element. The second and third term, however, contain $<f|H_{e.m.}|k>$ and $<k|H_{e.m.}|i>$, which are E1 matrix elements; this is because the parity of $|k>$ is opposite to that of $|i>$ and $|f>$. As a result, a nominal M1 transition actually acquires a small amount of E1 components with coefficients going like $<k|H_{pv}|i>/(E_i - E_k)$ etc. This can lead to an observable parity-violating effect. For example, the photon absorption cross section turned to a particular radiative transition may depend on the incident photon helicity (circular polarization).

In practice we must work with transitions where parity-violation asymmetries are large enough to be measurable. Even though atomic parity-violation experiments with hydrogen were discussed a long time ago (Zeldovich 1959, Michel 1965), practical proposals on atomic parity experiments were not made until the pioneering work of Bouchiat and Bouchiat (1974). In this paper the authors made the crucial observation that parity-violating effects can be greatly enhanced by working with heavy (high Z) atoms. We already remarked that as far as the $A_{lept}V_{quark}$ interactions are concerned, the contributions from the individual nucleons add coherently. Because N is roughly proportional to Z, the formula for Q_W (5.16) gives one factor of Z, i.e. the "weak charge" grows roughly as the electric charge. Second, the parity-violating potential (5.19) is proportional to the electron momentum operator, and the "velocity" of the valence

electron is known to scale like Z. Third, because of the short-ranged

nature of H_{pv}, the parity violating matrix element depends on $|\psi|^2$ of

the valence electron evaluated at the origin; with screening taken

into account, this can be shown to vary like Z, which, incidentally,

has been checked through a study of the Z dependence of hyperfine

structure (the Fermi-Segré rule). So altogether we gain Z^3. For

bismuth Z^3 is as large as $\sim 5.7 \times 10^5$.

There are two other ways to enhance parity-violation effects.

For a nonminal Ml transition parity-violation asymmetry goes as

$\mathcal{E}_{pv}/m$ where $\mathcal{E}_{pv}$ and m are the matrix elements for the parity-

violating El and the parity-conseving Ml amplitude, respectively.

So it helps to make m small by choosing a transition where the parity-

conserving transition is suppressed, as in so-called "forbidden" Ml

transitions. Another consideration that follows from (5.21) is

that the parity-mixing coefficient can be made large when the energy

denominator is small. This can be accomplished by choosing a level

with an almost degenerate opposite-parity level nearby and further

controlling the various hyperfine sublevels by subjecting the atom

to an external magnetic field until some of them actually cross.

The key to hydrogen and deuterium experiments now in progress

exploit this technique (Lewis and Williams 1975, Hinds and Hughes

1977).

Conceptually the simplest atomic parity experiment is of the

type originally conceived by Bouchiat and Bouchiat (1974) where

atoms - in practice metallic vapor - are irradiated by a circularly polarized laser beam tuned to some particular radiative transition where parity-violating asymmetry is predicted to be large enough to be measurable. The photon absorption cross section σ is then measured by detecting resonance fluorescence due to the decay of the level excited by the laser beam. A nonvanishing value of "circular dichroism"

$$A = [\sigma \ (\lambda = 1) - \sigma(\lambda = -1)]/[\sigma(\lambda = 1) + \sigma(\lambda = -1)] \quad (5.22)$$

where λ stands for the incident photon helicity would conclusively show parity violation in atoms. Notice that this expression is quite analogous to (5.5); the only difference is that here we are talking about a polarized photon beam of a few eV instead of a polarized electron beam of 19 GeV/c. For a nominal M1 transition, A is given by

$$A = 2 \ \text{Im}(\mathscr{E}_{pv}^{c} m)/\left[|m|^2 + |\mathscr{E}_{pv}^{c}|^2\right]$$

$$\simeq 2 \ \text{Im} \ \mathscr{E}_{pv}^{c}/m. \quad (5.23)$$

For typical transitions considered to be feasible experimentally, A is of order $10^{-4} - 10^{-7}$.

No experiments that exploit the original Bouchiat-Bouchiat proposal have been successfully completed to date. The first

positive evidence for parity violation in atoms was published by a Novosibirsk team (Barkov and Zoltorev, 1978 and 1979) who studied optical rotation in bismuth atoms. Similar experiments have also been performed at Seattle (Lewis et at. 1977; Hollister et al. 1980) and Oxford (Baird et al. 1977; Baird 1980). The basic physics behind these optical experiments goes as follows (Khriplovich 1974). First, recall the Lorentz formula

$$n_\pm = 1 + (2\pi/\omega^2)N\, f_\pm(\omega) \qquad (5.24)$$

which relates the index of refraction $n_\pm$ of a medium with N atoms per unit volume to the forward scattering amplitude of light $f_\pm(\omega)$ where ± denotes the photon helicity. Near the resonance line, the amplitude $f_\pm(\omega)$ is expected to take a typical resonance form:

$$f_\pm(\omega) = -\omega^2[\,|m|^2 \pm 2\mathrm{Im}(\mathcal{E}_{pv}\,m)\,]/[\omega - \omega_2 + i\frac{\Gamma}{2}] \qquad (5.25)$$

The fact that the right and left circularly polarized beams of light propagate with slightly different indices of refraction implies that when we start with a linearly polarized beam of light, which is a superposition of two light beams with opposite circular polarization, the plane of polarization rotates by a tiny amount. This rotation angle can be worked out to be

$$\phi_{pv} = (\omega\ell/2)\mathrm{Re}(n_+ - n_-) \qquad (5.26)$$

where ℓ is the path length. It is amusing that the original Bouchiat-Bouchiat proposal is to measure the difference in the imaginary part of the index of refraction between the two photon helicity states while the optical rotation experiments measure the difference in the real part. When attenuation of the beam due to absorption of

light is taken into account, an optimal condition is obtained when ℓ is comparable to absorption length ℓ_o. From (5.24) and (5.26) we see that the rotation angle follows a dispersion-like dependence on ω and exhibits a strikingly asymmetric curve vanishing exactly at $\omega = \omega_r$. The maximum difference in rotation angle as we sweep through the resonance region by varying ω can be derived to be

$$\Delta \phi_{pv} \Big|_{\ell = \ell_o} = Im(\mathcal{E}_{pv})/m \equiv R \qquad (5.27)$$

at one unit of absorption length. It is this quantity R, referred to as "rotation parameter", that is usually quoted in the various optical rotation experiments.

All three groups who studied optical rotation have used bismuth (Bi) atoms. The ground state of Bi belongs to the configuration $6p^3$, i.e. 3 valence electrons in the 6p shell outside the closed shells composed of 80 electrons. The Novosibirsk and Oxford groups utilize the 647 nm M1 transition $(4S_{3/2} \to {}^2D_{5/2})$ while the Seattle group, the 876 nm M1 transition $(4S_{3/2} \to {}^2D_{3/2})$. Both lines can be conveniently studied using tunable lasers.

The most recent results of the three groups are summarized in Table 5.1. Also shown are the standard-model predictions with

$\sin{}^2\theta_W = 0.23$ based on the atomic physics calculations of Novikov,

Sushkov and Khriplovich (1976), Sandars (1980) and Martensson,

Henley and Wilets (1980). It is beyond the scope to this review to

discuss various problems connected with the atomic physics calculations.

The reader is referred to a review paper of Fortson and Wilets (1980)

which contains additional references on the atomic physics calculations.

An atomic parity experiment of a different kind has been

successfully performed by Commins and collaborators at Berkeley

(Conte et al. 1979, Bucksbaum, Commins and Hunter 1981). As in the

original Bouchiat-Bouchiat proposal, the Berkeley experiment starts

with a circularly polarized beam incident on metallic vapor (Thallium

in this case). But instead of studying $\mathcal{E}_{pv}\mathcal{m}$ interference directly

by measuring circular dithroism, the atoms are placed in a static

electric field, and interference is looked for between $\mathcal{E}_{pv}$ and a

Stark-induced El amplitude; the actual quantity detected is the

polarization $\langle\vec{J}\rangle$ of the level excited by circularly polarized light.

Because this experiment is described in detail in a review paper by

Commins and Bucksbaum (1980) that appeared in this annual series,

we do not discuss the experimental method any further. The most

recent result of this group, when expressed in terms of circular

dichroism A as defined in (5.22) is

$$A(6{}^2P_{1/2} \to 7{}^2P_{1/2}) = 2.9{}^{+1.0}_{-0.9} \times 10^{-3} . \tag{5.28}$$

In principle the atomic physics calculations on thallium atoms
should be less difficult than in the bismuth case because there is only
one valence electron in the 3p configuration. The best theoretical
estimate for this transition is quoted to be (Neuffer and Commins 1977;
Sushkov, Flambaum and Khriplovich 1976)

$$A_{theory} \ (6^2P_{1/2} \rightarrow 7^2P_{1/2}) = (2.1 \pm 0.7) \times 10^{-3} \ . \qquad (5.29)$$

What can we conclude from the three Bi experiments of Table 5.1
and the Berkeley Tℓ experiment? Even though there is still some
discrepancy, viz. between Novosibirsk and Oxford for the 647 nm
transition of Bi, it does appear that the various experimental results
are slowly converging towards the theoretical predictions based on
the standard model. Considering that the atomic physics calculations
themselves may have uncertainties of $\sim 30\%$, the agreement with the
standard model must be regarded as satisfactory. To do better we may
have to await precise results of very difficult hydrogen and deuterium
experiments where there are no uncertainties in atomic physics
calculations.

5.4. Determination of the Parity-violating Coupling Parameters

We have emphasized that both the electron-deuteron experiment
and the atomic physics experiments are in good agreement with the
standard model with the value of $\sin^2\theta_W$ independently determined
from the neutrino-induced reactions. It is nevertheless worth pressing
our purely model-independent analysis based on $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}$, and $\tilde{\delta}$ to examine

to what extent these experiments determine the four constants

individually.

As pointed out in Section 5.2, the SLAC measurement of the a_1

coefficient determines the linear combination $3\tilde{\alpha} + \tilde{\gamma}$. In contrast,

the heavy atom (Bi, Tℓ) experiments measure Q_W [see (5.16)] given by

$$Q_W(\text{Bi}) = 43\tilde{\alpha} - 627\tilde{\gamma} = 43(\tilde{\alpha} - 14.6\tilde{\gamma})$$

$$Q_W(\text{Tℓ}) = 42\tilde{\alpha} - 612\tilde{\gamma} = 42(\tilde{\alpha} - 14.6\tilde{\gamma}) \ . \qquad (5.30)$$

We see that the linear combinations measured by the SLAC experiment

and the heavy atom experiments are almost "orthogonal". The SLAC

experiment is more sensitive to the isovector constant $\tilde{\alpha}$ while the

heavy atom experiments are more sensitive to the isoscalar constant

$\tilde{\gamma}$. They therefore provide complementary information.

The best way to see the significance of the two types of

experiments is to display the experimental constraints in an $\tilde{\alpha} - \tilde{\gamma}$

plane (Bjorken 1978b; Sakurai 1978; Hung and Sakurai 1979). See

Fig. 5.4. For the heavy atom experiments we use the Seattle Bi(876 nm)

(together with the average of the three atomic physics calculations

of Table 5.1) and the Berkeley Tℓ result. The Bi (647 nm) results

of Oxford and Novosibirsk are not used here in view of the discrepancy

in the data by nearly a factor of two. In terms of Q_W, the recent Seattle

result for R(876nm) gives

$$Q_W(\text{Bi}) = - 115 \pm 19 \ , \qquad (5.31a)$$

while the Berkely result implies

$$Q_W(\text{Tℓ}) = - (170 \ {}^{+59}_{-53}) \qquad (5.31b)$$

Using these values, we can draw "atomic-physics bands" as shown in Fig. 5.4.

We see that the intersection of the SLAC a_1 band and the heavy atom band determine <u>uniquely</u> $\tilde{\alpha}$ and $\tilde{\gamma}$ <u>including signs</u>:

$$\tilde{\alpha} = -0.67 \pm 0.19$$
$$\tilde{\gamma} = 0.22 \pm 0.12 \tag{5.32}$$

Notice that no gauge theory concepts have been used to extract the two parameters of the $A_{lept}V_{quark}$ interactions.

It would be nice if we could carry out analogous analyses for the $V_{lept}A_{quark}$ interactions. Unfortunately without hydrogen and deuterium experiments we cannot pin down $\tilde{\beta}$ and $\tilde{\delta}$ separately. All we know at this moment is the linear combination (5.11b) with a large error based on the SLAC a_2 measurement.

VI. CHARGED LEPTON INTERACTIONS

6.1 Weak Interaction Predictions for Muon Pair Production and Bhabha Scattering

In this chapter we concern ourselves with the neutral-current reactions involving charged leptons only (e,μ,τ). At this moment this subject is an exclusive domain of electron-positron colliding beam physics. The basic processes of interest are electron-positron annihilation into muon pairs

$$e^+ + e^- \rightarrow \mu^+ + \mu^- \tag{6.1}$$

and Bhabha scattering

$$e^+ + e^- \rightarrow e^+ + e^- . \tag{6.2}$$

In (6.1) the neutral current appears in the s channel, hence the q^2 for the current is time-like; in contrast in (6.2) the neutral current appears in both the s channel and t channel, so the q^2 for the current can be space-like as well as time-like. As already mentioned in Section 2.1, with μe universality assumed, there are three low-energy parameters, h_{VV}, h_{AA} and h_{VA}, to be determined. We can also extend our consideration to the τ lepton whose neutral-current coupling can be studied in

$$e^+ + e^- \rightarrow \tau^+ + \tau^- \tag{6.3}$$

in a manner completely analogous to a study of (6.1).

The history of this subject is again rather old. As early as 1961, when electron-positron colliding-beam facilities were first contemplated at Orsay and Frascati, Cabibbo and Gatto (1961) worked out the phenomenology of weak effects in muon pair production (6.1) within the framework of a weak-interaction model containing a pure V-A neutral current. Now seven machine generations later - ACO, VEPP, ADONE, CEA-Bypass, SPEAR, DORIS, DCI - a study of neutral current effects via weak-electromagnetic interference in these reactions has finally become a reality at PETRA (and also PEP). We may recall the $10^{-4} q^2$ rule of Section 5.1; with $s \simeq 1000$ GeV^2 typically at PETRA and PEP, we should find $\sim 10\%$ effect.

Let us first look at the muon pair reaction (6.1). If neither the electron, nor the positron beam is polarized, there are two kinds of effects we can look at. First, we may examine whether the observed total cross section for (6.1) deviates from the QED prediction. For $s \ll m_Z^2$, this measures h_{VV} as follows (Kinoshita et al. 1970; Wolfenstein 1974)

$$(\sigma - \sigma_{QED})/\sigma_{QED} = -4(G/\sqrt{2} \; e^2) \; h_{VV} s \qquad (6.4)$$

Second, we may examine neutral-current effects in the angular distribution. Again, with $s \ll m_Z^2$, we can write the angular distribution of the outgoing μ^- relative to the incident e^- direction as (Kinoshita et al. 1970; Wolfenstein 1974)

$$d\sigma/d(\cos\theta) = (\pi\alpha^2/2s)[A(1 + \cos^2\theta) + B \cos\theta] \qquad (6.5)$$

where

$$A = 1 - 4(G/\sqrt{2}\ e^2)h_{VV}s, \qquad (6.6a)$$

$$B = - 8(G/\sqrt{2}\ e^2)h_{AA}s\ . \qquad (6.6b)$$

Notice that the $\cos\theta$ term is due solely to the weak-electromagnetic interference; in practice, however, a $\cos\theta$ term arising from higher order (two-photon exchange) electromagnetic interactions must be subtracted out before applying our formulas. This leads to an integrated forward-backward asymmetry

$$<A_{FB}> \equiv (F - B)/(F + B)$$

$$\simeq -3(G/\sqrt{2}\ e^2)h_{AA}s. \qquad (6.7)$$

Finally, to determine the third parameter h_{VA}, one must measure the longitudinal polarization of the outgoing muon or study the cross section dependence on the helicity of the incident electron (or positron).

Let us now turn to the standard-model predictions for the coupling parameters. Using the $T_{3L} - Q \sin^2\theta_W$ rule of Section 2.2,

we readily obtain

$$h_{VV} = \frac{1}{4} (1 - 4 \sin^2\theta_W)^2 , \qquad (6.8a)$$

$$h_{AA} = \frac{1}{4} , \qquad (6.8b)$$

$$h_{VA} = \frac{1}{4} (1 - 4 \sin^2\theta_W) . \qquad (6.8c)$$

Because $\sin^2\theta_W$ is known to be close to 1/4 from other experiments, we expect h_{VV} to be essentially zero. So the fractional deviation from the QED cross section, given by (6.4), is predicted to be too small to be measureable. On the other hand, a sizable effect is predicted for $< A_{FB} >$. At $s = 1347$ GeV^2, corresponding to the highest (at this writing) PETRA energy of 36.7 GeV, the formulas (6.7) and (6.8b) predict $< A_{FB} >$ to be as large as -8.7%.

The formulas (6.4)-(6.7) ignore finite Z mass effects due to the Z boson propagator. If s is still small but not completely negligible compared to m_Z^2, we need to multiply the predictions by $m_Z^2/(m_Z^2 - s)$. So this has a tendency of <u>enhancing</u> the weak-interaction effects. For example, with m_Z taken to be 90 GeV, our asymmetry prediction at $s = 1347$ GeV^2 changes from -8.7% to -10.4%.

The formulas analogous to (6.5) and (6.6a,b) can be derived for Bhabha scattering:

$$[(d\sigma/d\Omega) - (d\sigma/d\Omega)_{QED}]/(d\sigma/d\Omega)_{QED}$$

$$= (2/\sqrt{2}\pi\alpha) \left\{ h_{VV}[(3+\cos^2\theta)/(1-\cos\theta)] \left[\frac{1}{2}(3+\cos\theta) G(s,m_Z^2) \right. \right.$$

$$\left. + \cos\theta\, G(s,m_Z^2) \right] - h_{AA}\left[\frac{1}{2}(7+4\cos\theta+\cos^2\theta)G(s,m_Z^2) \right]$$

$$\left. - h_{AA}[(1+3\cos^2\theta)/(1-\cos\theta)]\, G(t,m_Z^2) \right\}/[(3+\cos^2\theta)/(1-\cos\theta)]^2$$

$$+ \text{terms of order } G^2 \qquad\qquad (6.9)$$

which is a model-independent way of writing down an expression first obtained (correctly) by Bundy and McDonald (1974). In (6.9) the Fermi constant G is written in such a way to exhibit its effective s or t dependence

$$G(s,m_Z^2) \equiv G/(1 - s/m_Z^2) \xrightarrow{s \ll m_Z^2} G ,$$

$$G(t,m_Z^2) \equiv G/(1 - t/m_Z^2) \xrightarrow{|t| \ll m_Z^2} G , \qquad (6.10)$$

where t, in terms of the center-of-mass angle θ, is given by

$$t = -s \sin^2(\theta/2) \qquad\qquad (6.11)$$

6.2   Recent PETRA Results

Since the early days of colliding-beam physics, it has been
traditional to parametrize possible departures from the QED predictions
in muon pair production (6.1) and Bhabha scattering (5.2) by modifying
the photon propagator, e.g. in the s channel for the muon pair reaction
(6.1), as follows:

$$1/s \rightarrow 1/s \mp 1/(s - \Lambda_\pm^2) \ . \tag{6.12}$$

When the reactions (6.1) and (6.2) were studied at PETRA with s typically
around 1000 GeV^2, new lower limits on $\Lambda_\pm$ were reported by various
collaboration groups, typically in the 100 - 200 GeV range (Barber
et al. 1979; Böhm 1980).

Such an analysis, however, presupposes that the electron and
muon do not enjoy interactions other than pure QED. In reality, these
charged leptons are known to participate in the weak neutral-current
interactions, presumably mediated by an object of mass $\sim$ 100 GeV
with a dimensionless coupling constant of order e. With the lower
limit on $\Lambda_\pm$ now reaching and even surpassing the conjectured Z mass,
we should regard the so-called QED tests as tests of electroweak
models at unprecedentedly high values of q^2 or s. (Wright and
Sakurai, 1980).

Specifically, for the muon pair production reaction (6.1),
the QED parameterization (6.12) leads to

$$(\sigma - \sigma_{QED})/\sigma_{QED} \simeq \pm 2s/\Lambda_{\pm}^2 \qquad (6.13)$$

for $\Lambda_{\pm}^2 \gg s$. Comparing this with the neutral-current formula (6.4), we see that a lower limit on Λ_- implies an upper limit on h_{VV}, or within the context of the standard model, an allowed range for $\sin^2\theta_W$. For example,

$$\Lambda_- > 160 \text{ GeV} \Rightarrow h_{VV} < 0.22, \quad 0.02 < \sin^2\theta_W < 0.48 \ . \quad (6.14)$$

The fact that no QED violations have been observed in the integrated cross section for muon pair production may be interpreted to mean that h_{VV} is small, as expected from the standard model ($h_{VV} = 0.0016$ for $\sin^2\theta_W = 0.23$).

Attempts have also been made at PETRA to measure a possible forward-backward asymmetry in the angular distribution of the muon pair reaction (6.1). As emphasized in the previous section, this is sensitive to h_{AA}. At the time of the Madison Conference (July 1980) the "combined" PETRA results - based on work of the JADE, Mark J, PLUTO and TASSO Collaborations - were quoted as

$$< A_{FB} > = (- 0.9 \pm 4.9)\% \qquad (6.15)$$

where the standard model predicts -6% for the energy and angular range covered by the measurements (Böhm 1980).

More recently, the JADE Collaboration has extracted h_{VV} and
h_{AA} by analyzing simultaneously the muon pair data and the Bhabba
scattering data (Bartel et al. 1981, see also earlier Mark J data of
Barber et al. 1980). Assuming $s \ll m_Z^2$, they obtain

$$h_{VV} = 0.01 \pm 0.08, \quad (0.0016 \text{ for } \sin^2\theta_W = 0.23)$$

$$h_{AA} = 0.18 \pm 0.16 \quad (0.25) \tag{6.16}$$

where the values in parentheses are the standard-model predictions.
The allowed region in an h_{VV}-h_{AA} plane is indicated in Fig. 6.1. The
errors are still large, but it is amusing that the "origin" $h_{VV} = h_{AA} = 0$,
i.e. no neutral-current interactions, is not favored by the data. We
may also note that some once-popular models - e.g. the SU(2) symmetric
V-A model that predicts $h_{VV} = h_{AA} = 1/4$ and the 'hybrid model" that
predicts $h_{VV} = 1/4$, $h_{AA} = 0$ - are now excluded by the PETRA data alone.

As will be emphasized in the next chapter, given the successes of
the standard model in the νq, νe and eq sectors, the standard model
predictions for h_{VV} and h_{AA} are the predictions of any single Z boson
model. We may therefore ask what new information can be obtained by
studying weak interaction effects in these reactions. We should keep
in mind, however, that here the weak-neutral-current interactions are
being explored at values of q^2 much higher than are accessible in
other experiments. These reactions are therefore particularly
sensitive to possible nonlocal effects due to the weak boson

propagator, etc. For example, we could imagine a phenomenological (nongauge) model in which the low-energy predictions for the neutral-current interactions coincide exactly with those of the standard model with, say, $\sin^2\theta_W = 0.23$ and yet Weinberg's Z mass prediction may fail (Bjorken 1978a, 1979; Hung and Sakurai 1978). The asymmetry prediction in muon pair production in such a phenomenological model is given in Fig. 6.2 for various values of m_Z, now taken as a completely free parameter. With $m_Z \simeq 50$ GeV, the asymmetry is predicted to be about 3 times larger at typical PETRA energies. Analyses of the JADE data within such a theoretical framework have led to (Bartel et al. 1981)

$$m_Z > 51 \text{ GeV (95\%CL)} \; . \tag{6.17}$$

The PETRA data were also used to examine a possible modification of the standard model as follows

$$\mathcal{L}_{eff} = (4G/\sqrt{2})\left[(J_\lambda^3 - \sin^2\theta_W J_\lambda^{e.m.})^2 + c(J_\lambda^{e.m.})^2\right] \; . \tag{6.18}$$

This structure is actually predicted in some larger group models, $SU(2) \otimes U(1) \otimes U(1)$ or $SU(2) \otimes SU(2) \otimes U(1)$ (de Groot, Gounaris and Schildknecht 1979, Barger, Keung and Ma 1980). Except for the c term, this effective interaction is precisely that of the standard model. Now the extra c term makes no contribution to neutrino-induced reactions because the neutrino has no electric charge; it

does not show up in the SLAC asymmetry either because of parity. In fact, the ideal place to test its presence is precisely muon pair production and Bhabha scattering. The current limit on c from a combined analysis of the various Collaborations at PETRA is quoted to be (Böhm 1980)

$$c < 0.03. \quad (95\% \text{ CL}) . \qquad (6.19)$$

This is a nontrivial result. We cannot tolerate an extra $(J^{e.m.})^2$ term entering with strength of a few % of the normal term.

VII. COMPARISONS WITH MODELS

7.1 Factorization Tests

This chapter is concerned with comparisons of the experimentally determined coupling parameters with the predictions of various weak-interaction models. Before turning to electroweak gauge models, we first discuss the <u>factorization</u> constraints that must be satisfied by any model in which the neutral-current interactions are mediated by a single Z boson.

With the single Z boson hypothesis the interactions appearing in the neutral-current pyramid of Fig. 2.1 can be completely determined by specifying seven parameters – the couplings of Z to $\bar{\nu}_L \nu_L$, and to $\bar{u}u$, $\bar{d}d$ and $\bar{e}e$ (vector and axial-vector for each) – where we have assumed μe universality. On the other hand, there are as many as 13 phenomenological parameters in the neutral-current pyramid. We therefore expect that in a single Z model there must be six independent "factorization" relations among the thirteen parameters (Hung and Sakurai 1977a).

As an example to illustrate the factorization constraints, let us look at Fig. 7.1. The product of the $\nu\nu$ scattering amplitude and the eq scattering amplitude in appropriate helicity states must be equal to the product of the νq scattering amplitude and the νe scattering amplitude. In our normalization convention this means

$$c_\nu^2 = 2g_A \alpha / \tilde{\alpha} = 2g_A \gamma / \tilde{\gamma} = 2g_V \beta / \tilde{\beta} = 2g_V \delta / \tilde{\delta} \qquad (7.1)$$

where we have introduced the $\nu\nu$ interaction constant c_ν^2 via

$$\mathcal{L}_{eff} = - (Gc_\nu^2/\sqrt{2})[\bar{\nu}\gamma_\lambda(1 + \gamma_5)\nu][\bar{\nu}\gamma_\lambda(1 + \gamma_5)\nu]. \qquad (7.2)$$

In the standard model, or, more generally, in any $SU(2) \otimes U(1)$ with the Higgs in doublets, c_ν^2 is unity.

Upon eliminating c_ν^2, we may rewrite (7.1) as

$$\tilde{\gamma}/\tilde{\alpha} = \gamma/\alpha , \qquad (7.3a)$$

$$\tilde{\delta}/\tilde{\beta} = \delta/\beta , \qquad (7.3b)$$

$$g_V/g_A = \alpha\tilde{\beta}/\beta\tilde{\alpha} . \qquad (7.3c)$$

Equations (7.3a) and (7.3b) state that the isoscalar-isovector ratios measurable in the eq interactions must be equal to the corresponding ratios measurable in the νq interactions while (7.3c) states that the V- to -A ratio in νe scattering can be inferred from the νq and eq data.

To what extent do the data available up to now support these factorization relations? First, as a test of (7.3a), let us go back to Fig. 5.4. It is seen that the intersection of the SLAC a_1-asymmetry band and the atomic parity band lies comfortably within the region allowed by the right-hand side of (7.3a) as evaluated

using Solution A of the νq sector. As a second test of factorization, we combine (7.3a) – (7.3c) to write

$$g_V/g_A = [(\alpha + \gamma/3)(\widetilde{\beta} + \widetilde{\delta}/3)]/[(\widetilde{\alpha} + \widetilde{\gamma}/3)(\beta + \delta/3)] \quad (7.4)$$

where the right-hand side can be evaluated from the νq data and the SLAC asymmetry with $a_1 - a_2$ separation. The region allowed by the right-hand side of (7.4) is shown in the $g_V - g_A$ plane of Fig. 3.1. We see that the vector dominant solution in νe scattering is comfortably ruled out by factorization while the axial-vector dominant solution is in beautiful agreement; so, within the framework of single Z boson models, the VA ambiguity in νe scattering mentioned in Section 3 is unambiguously resolved in favor of the axial-vector dominant solution. This choice agrees with the standard model.

We remarked earlier that there should be altogether six factorization relations while we have written down only three so far, vis. (7.3 a,b,c). The remaining three factorization relations all concern the $(\bar{e}e)$ $(\bar{\mu}\mu)$ constants, h_{VV}, h_{AA} and h_{VA}:

$$h_{VV} = g_V^2/c_\nu^2 \quad , \qquad\qquad\qquad (7.5a)$$

$$h_{AA} = g_A^2/c_\nu^2 \quad , \qquad\qquad\qquad (7.5b)$$

$$h_{VA}^2 = h_{VV}h_{AA} \quad , \qquad\qquad\qquad (7.5c)$$

where c_ν^2 may be obtained from (7.1). Even though the errors are still large, the JADE values of h_{VV} and h_{AA} appear to suggest that h_{AA} is larger than h_{VV} [see (6.16)]. With factorization this would, in turn, imply $g_A^2 > g_V^2$. So, if future experiments confirm $h_{AA} > h_{VV}$, there will be an additional argument in favor of the axial-dominant solution in νe scattering, independently of the eq data. Historically Barber et al. (1980) first reached this conclusion by equating h_{AA} and h_{VV} (as obtained by the Mark J Collaboration) directly to g_A^2 and g_V^2. As is apparent from (7.5 a,b) such an approach presupposes $c_\nu^2 = 1$. Even though c_ν^2 is unity in many gauge models and a simultaneous analysis of the νe, νq and eq constants can be shown to lead to $c_\nu^2 \simeq 1$ (see below), it is contrary to the spirit of a purely phenomenological analysis to assume $c_\nu^2 = 1$ from the very beginning.

As for (7.5c), we cannot test it until the spin dependence of electron-positron annihilation into muon pairs is studied. In view of the sucesses of the standard model in other reactions, it is likely that this relation will be trivially satisfied by having both sides vanishingly small.

So far we have been concerned with testing factorization. If we $\underline{assume}$ factorization, many additional results of interest can be obtained. First, the $V_{lept}A_{quark}$ constants, $\tilde{\beta}$ and $\tilde{\delta}$, which have not been separately determined, can now be inferred using

$$\left\{\begin{matrix} \tilde{\beta} \\ \tilde{\delta} \end{matrix}\right\} = \left\{\begin{matrix} \beta \\ \delta \end{matrix}\right\} (g_V/g_A)(\tilde{\alpha} + \tilde{\gamma}/3)/(\alpha + \gamma/3) \qquad (7.6)$$

Second, the strength of the $\nu\nu$ interaction constant c_ν^2 can be

evaluated using

$$c_\nu^2 = 2g_A(\alpha + \gamma/3)/(\tilde{\alpha} + \tilde{\gamma}/3) \qquad (7.7)$$

to

$$c_\nu^2 = \pm (0.86 \pm 0.28) \qquad (7.8)$$

Because h_{AA} is likely to be positive according to the colliding

beam data, the factorization (7.5b) forces us to choose the positive

sign in (7.8); both $h_{AA} > 0$ and $c_\nu^2 > 0$ follow from the requirement that

the Z boson appears in the s channel of the e^+e^- or the $\bar{\nu}\nu$ interactions

as a physical particle, not as a ghost. Knowing c_ν^2 to be positive,

we can use (7.1) to resolve the sign ambiguity in the νq parameters

α, β, γ and δ; recall, in contrast, that there is no sign ambiguity in

the νe scattering constants g_V and g_A, nor in the eq constants $\tilde{\alpha}, \tilde{\beta}, \tilde{\gamma}$

and $\tilde{\delta}$.

When all this is done, we see that the 13 constants of the neutral-

current pyramid are completely determined including their signs. In

Table 7.1 we present the present status of the 13 parameters obtained

without (model independent) and with (factorization dependent) the

factorization constraints taken into account. Also shown are the

predictions of the standard model. We see that the coupling parameters

determined without recourse to gauge-theory constraints are in

excellent agreement with the standard model predictions with $\sin^2\theta_W$ set equal to 0.23.

7.2   Comparisons with SU(2) ⊗ U(1)

Throughout this review paper we have presented comparisons of the experimental data with the one-parameter standard model based on SU(2) ⊗ U(1). As emphasized in Section 2.2, the standard model assumes the following multiplet assignments:

(i)   the isodoublet assignment for the Higgs boson,

(ii)   the isodoublet assignment for left-handed fermions $(T_{3L} = \pm\frac{1}{2})$,

(iii)   the isosinglet assignment for right-handed fermions $(T_{3R} = 0)$.

If we relax any one or more of the above assumptions, different predictions are possible while still staying withing the general framework of SU(2) ⊗ U(1).

Suppose we relax (i). The overall strength of the neutral-current interactions, commonly denoted by ρ, then becomes a free parameter as far as low energy phenomenology is concerned. Because of the success of the V-A rule in the charged-current interactions, it is customary to assume that the left-handed fermions belong to weak-isospin doublets. So we do not relax (ii) but continue assuming

$$T_{3L}(u) = T_{3L}(\nu) = 1/2, \ T_{3L}(d) = T_{3L}(e) = -1/2. \quad (7.9)$$

In contrast, (iii) has often been relaxed; there have been many papers where some right-handed fermions are assumed to be in nonsinglets.

As an illustration of all this, let us look at the νe constants
g_V and g_A. When both ρ and T_{3R} are taken to be free parameters, the predictions for g_V and g_A are obtained as follows:

$$g_V = \rho \, [- \tfrac{1}{2} + T_{3R}(e) + 2 \, \sin^2\theta_W],$$

$$g_A = \rho \, [- \tfrac{1}{2} - T_{3R}(e)]. \qquad\qquad (7.10)$$

The once popular "hybrid model" (Cheng and Li, 1977) was based on the idea that the right-handed electron forms a doublet together with a very massive neutral lepton N^o so that

$$T_{3R}(N^o) = 1/2, \quad T_{3R}(e^-) = -1/2 \qquad\qquad (7.11)$$

It is seen that in such a model the coupling of the electron is purely vectorial ($g_A = 0$), which is now ruled out by the SLAC asymmetry etc.

We may analyze all available experimental data using a five-parameter SU(2) ⊗ U(1) model with ρ, $\sin^2\theta_W$, $T_{3R}(u)$, $T_{3R}(d)$ and $T_{3R}(e)$ as free parameters to be determined. A very extensive analysis along this line has been made by Kim et al. (1981) [see also earlier attempts by Hung and Sakurai (1977b) and Roos and Liede (1979)] who obtain

$$\rho = 1.018 \pm 0.045, \qquad \sin^2\theta_W = 0.249 \pm 0.031,$$

$$T_{3R}(u) = -0.010 \pm 0.040, \qquad T_{3R}(d) = -0.101 \pm 0.058$$

$$T_{3R}(e) = 0.039 \pm 0.047 \ . \qquad\qquad (7.12)$$

The T_{3R}'s are all seem to be close to zero and quite far from $\pm 1/2$. So the data are consistent with the absence of right-handed doublets, or, more quantitatively, the mixing between the usual doublets and possible singlets is less than 10.3% for u, 34.8% for d and 6.4% for e.

Next we may set the T_{eR}'s to be all zero. A two-parameter fit with ρ and $\sin^2\theta_W$ as adjustable parameters then leads to [Kim et al. (1981); see also earlier attempts by Sehgal (1978), Roos and Liede (1979)]

$$\rho = 1.002 \pm 0.015 \ (\pm 0.011),$$

$$\sin^2\theta_W = 0.234 \pm 0.013 \ (\pm 0.009)$$

$$(\chi^2/_{DOF} = 33.1/44),$$

where the errors in parentheses are due to theoretical uncertanities. This fit is in excellent agreement with the hypothesis $\rho = 1$. Finally, when a fit solely to $\sin^2\theta_W$ is made, these authors obtain for the

only parameter of the standard model

$$\sin^2\theta_W = 0.233 \pm 0.009 \ (\pm \ 0.005)$$

$$(\chi^2/\text{DOF} = 33.1/45) \ . \qquad\qquad (7.14)$$

Another way to test the standard model is to list the values of $\sin^2\theta_W$ obtained in various experiments. This approach is summarized in Table 7.2, which is an update of a compilation made by Baltay (1980).

7.3 Right-left Symmetric Models

Even though both charged-current and neutral-current phenomena violate parity, it is attractive from a certain point of view to suppose that the basic Lagrangian of the world is right-left symmetric before spontaneous breakdown. Right-left symmetric models based on $SU(2)_L \otimes SU(2)_R \otimes U(1)$ are motivated by such a belief [Pati and Salam, 1974; Fritzsch and Minkowski, 1976; Mohapatra and Sidhu, 1977]. It is beyond the scope of this review to discuss in detail how to arrange the Higgs mechanism in such a way to make the weak bosons associated with the right-handed currents much heavier. Here we concentrate on the low q^2 neutral-current interactions of such models, which can be shown to be parametrized by four constants, ρ_L, ρ_R, ρ_{LR} and $x(\equiv \sin^2\theta_W)$ [Ecker, 1979; Bajaj and Rajasekaran, 1980; Sidhu, 1980]:

$$\mathcal{L}_{eff} = (4G/\sqrt{2})\,[\rho_L(J^3_{\lambda L} - xJ^{e.m.}_\lambda)^2 + \rho_R(J^3_{\lambda R} - xJ^{e.m.}_\lambda)^2$$

$$- 2\rho_{LR}(J^3_{\lambda L} - xJ^{e.m.}_\lambda)(J^3_{\lambda R} - xJ^{e.m.}_\lambda)]\,. \qquad (7.15)$$

This class of models has two Z bosons whose masses m_1 and m_2 are given by

$$(37.3 \text{ GeV}/m_1)^2 + (37.3 \text{ GeV}/m_2)^2 = x[(1 - x)(\rho_L + \rho_R) + 2x\rho_{LR}],$$

$$(37.3 \text{ GeV}/m_1)^2(37.3 \text{ GeV}/m_2)^2 = x^2(1 - 2x)(\rho_L\rho_R - \rho^2_{LR})\,.$$
$$(7.16)$$

In the limit $\rho_L \to 1$, $\rho_R \to 0$, $\rho_{LR} \to 0$, one of the Z boson masses, say m_2, becomes infinite and we recover the one-parameter prediction of the standard model.

The experimental data may be used to determine the four parameters of the right-left symmetric models. A recent analysis by Sehgal (1980) leads to

$$\sin^2\theta = 0.25 \pm 0.02,$$
$$\rho_L = 1.0 \pm 0.06\,,$$
$$\rho_R = -0.2 \pm 0.2,$$
$$\rho_{LR} = -0.05 \pm 0.06\,. \qquad (7.17)$$

In terms of the weak boson masses, this fit implies

$$m_1 = (86 \pm 3)\,\text{GeV}, \qquad m_2/m_1 > 3.3 \qquad\qquad (7.18)$$

Dramatic differences between this class of models and the standard model will show up only at energies far above the first (lower) Z boson mass.

VIII.   OUTLOOK

We have discussed a variety of neutral-current phenomena spanning an enormously wide range of energy – from radiative atomic transitions in the few eV range and reactor processes initiatated by antineutrinos of a few MeV to neutrino-induced reactions up to a few hundred GeV and electron-positron colliding-beam processes with q^2 exceeding 1000 GeV^2 in the timelike direction. As repeatedly emphasized throught this review, these neutral-current experiments provide a spectacular confirmation of the low-energy form of the standard SU(2) × U(1) model due to Glashow, Weinberg, Salam and Ward.

We may naturally ask if there is any point in accumulating more data on the neutral-current interactions. The answer to this question is affirmative for several reasons. First, there are still undetermined coupling parameters. For instance, further work in atomic physics – hydrogen and deuterium experiments in particular – is desirable to obtain the values of the $V_{lept}A_{quark}$ parameters $\tilde{\beta}$ and $\tilde{\delta}$; these constants are expected to provide additional tests for factorization. Untouched in this review – because of lack of data – are the neutral-current couplings of strange and heavy quarks, i.e. the couplings of Z to $\bar{s}s$, $\bar{c}c$, $\bar{b}b$, $\bar{t}t$, etc; along this line a study of weak-interaction effects in the vicinity of a yet-to-be-discovered toponium (a $t\bar{t}$ bound state analogous to ψ/J) will be worth while (Bernabéu and Pascual, 1979 and 1980; Sehgal and Zerwas, 1980).

Possible small deviations from the simple standard-model predictions are also of great theoretical interest. Some departure from the realtion $\rho = 1$ is expected due to weak-isospin violation whenever there is a very massive charged lepton forming a weak-isospin doublet together with a massless or nearly massless neutrino (Veltman, 1977; Chanowitz et al., 1978). Turning the argument around, the fact that the experimentally determined value of ρ is close to unity may be used to obtain an <u>upper</u> limit for the mass of a very heavy lepton

$$m_L < 500 \text{ GeV} .\tag{8.1}$$

The higher order electroweak corrections to the neutral-current parameters are unambiguously calculable in electroweak gauge theory. For example, the simple standard-model prediction

$$\alpha - 3\gamma = 1 \tag{8.2}$$

is altered due to such higher order corrections to (Marciano and Sirlin 1980)

$$\alpha - 3\gamma = 0.974 . \tag{8.3}$$

Even more important are the corrections due to gluonic (QCD) inter-actions. The isoscalar axial-vector constant δ for the coupling of

$\bar{u}u + \bar{d}d$, predicted to be zero in the absence of such corrections, may change to $0 \cdot 1$ because $\bar{u}u + \bar{d}d$ can communicate with $\bar{c}c - \bar{s}s$ via OZI forbidden gluonic intermediate states (Collins, Wiczek and Zee, 1978).

Even though $\sin^2\theta_W$ is an arbitrary parameter in the standard $SU(2) \otimes U(1)$ model, when this model is embedded in the larger framework of "Grand Unification" that unites QFD with QCD, a precise prediction for $\sin^2\theta_W$ becomes possible. In particular, the SU(5) model Georgi and Glashow (1974), which is the simplest model of grand unification, predicts (Marciano and Sirlin, 1981)

$$\sin^2\theta_W^{(\nu_\mu; \ell)}(0) = 0.2104 \quad , \tag{8.4a}$$

for ν_μ - lepton scattering at $Q^2 = 0$ and

$$\sin^2\theta_W^{(\nu; h)}(Q^2 = 20 GeV^2) = 0.2098, \tag{8.4b}$$

for deep-inelastic ν-hadron scattering, where a particular value of QCD scale is chosen, namely $\Lambda_{\overline{ms}} = 0.4$ GeV. For this reason a precise experimental value of $\sin^2\theta_W$ is of great interest. We should remind the reader that the first serious calculation of $\sin^2\theta_W$ was made by George, Quinn and Weinberg (Georgi et al., 197) who predicted $\sin^2\theta_W \simeq 0.20$. The above predicted values of $\sin^2\theta_W$, (8.4a,b), are the most recent ones and are consistent with the best experimental value presented in Section 7.2.

We have concentrated in this review on the low energy or low q^2 manifestations of the neutral-current interactions. The range of q^2 explored in the experiments carried out so far is so low that we cannot yet discriminate between the zero-range "Fermi model" and a weak boson model with $m_Z \sim 100$ GeV. The next decisive step would be to observe deviations from the zero-range model and show that the neutral-(and charged-) current interactions are mediated by Z^o(and $W^{\pm}$) of <u>finite</u> mass.

The successes of the standard model, spectacular as they are, all refer to the low q^2 predictions of the model. It may be mentioned in this connection that these low q^2 successes may follow equally well in a more phenomenological (nongauge) model based on <u>global</u> SU(2) symmetry and γ-W^o mixing (Bjorken 1978a, 1979; Hung and Sakurai 1978); in such an approach the W and Z mass predictions of the standard model need not be obeyed. It is therefore crucial to prove (or disprove) the existence of the $W^{\pm}$ and Z bosons with the mass values predicted by Weinberg (1967). Only then can we say that we have a unified description of the electromagnetic and weak interactions.

Much has been done since the Gargamelle discovery of the neutral-current interaction in 1974. Much is yet to be done, especially to examine whether the weak interactions are indeed mediated by the W and Z bosons with the properties predicted by the standard model.

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Figure legends

Fig. 2.1   Neutral-current coupling pyramid.

Fig. 3.1   Determination of the neutrino-electron parameters g_V and

g_A. The shaded region indicates the constraint imposed by factori-

zation (see Sect. 7.1).

Fig. 4.1   The "Weinberg nose" in the $R_\nu - R_{\bar\nu}$ plane.

Fig. 4.2   Determination of $|\varepsilon_L(u)|^2$ and $|\varepsilon_L(d)|^2$ from the Hydrogen

ABCMO collaboration data and the Neon ABCLOS collaboration data

(isoscalar target). Also shown is the allowed (shaded) region

obtained from semi-inclusive pion production data (see Sect. 4.4).

Fig. 4.3   Allowed regions in neutrino-hadron coupling constant planes

obtained from semi-inclusive pion production data.

Fig. 4.4   The mass-distributions of the $\pi^0 p(\pi^- p)$ system in νN

collisions.

Fig. 4.5   The surviving Sol. A in the neutrino-hadron coupling plane.

Fig. 5.1   Weak-electromagnetic interference in the electron-proton

interaction.

Fig. 5.2   The y dependence of the parity-violating asymmetry A in

inelastic electron-deuteron scattering.

Fig. 5.3   Determination of a_1 and a_2 in inelastic electron-

deuteron scattering.

Fig. 5.4   Determination of the $A_{lept}V_{quark}$ parameters, $\tilde\alpha$ and $\tilde\gamma$.

Fig. 6.1   Determination of the $(\bar ee)(\bar\mu\mu)$ parameters, h_{VV} and h_{AA}.

Fig. 6.2   Angular-asymmetry prediction for $e^+ + e^- \to \mu^+ + \mu^-$ in

single Z models that are constrained to give the same low-energy

predictions as the standard model with $\sin^2\theta_W = 0.23$.

Fig. 7.1. Kindergarten approach to factorization.

	SU(2) $\otimes$ U(1) (General)	SU(2) $\otimes$ U(1) (minimal)
α	$\rho[1 + (T_{3R}^u - T_{3R}^d) - 2x]$	$1 - 2x$
β	$\rho[1 - (T_{3R}^u - T_{3R}^d)]$	1
γ	$\rho[(T_{3R}^u + T_{3R}^d) - \frac{2}{3}x]$	$-\frac{2}{3}x$
δ	$-\rho[T_{3R}^u + T_{3R}^d]$	0
g_ν	$\rho[-\frac{1}{2} + T_{3R}^e + 2x]$	$-\frac{1}{2} + 2x$
g_A	$-\rho[\frac{1}{2} + T_{3R}^e]$	$-\frac{1}{2}$
h_{VV}	$\rho[-\frac{1}{2} + T_{3R}^e + 2x]^2$	$[-\frac{1}{2} + 2x]^2$
h_{AA}	$\rho[\frac{1}{2} + T_{3R}^e]^2$	$\frac{1}{4}$
h_{VA}	$-\rho[\frac{1}{2} + T_{3R}^e][-\frac{1}{2} + T_{3R}^e + 2x]$	$\frac{1}{4} - x$
$\tilde{\alpha}$	$-2\rho[\frac{1}{2} + T_{3R}^e][1 + (T_{3R}^u - T_{3R}^d) - 2x]$	$-(1 - 2x)$
$\tilde{\beta}$	$2\rho[-\frac{1}{2} + T_{3R}^e + 2x][1 - (T_{3R}^u - T_{3R}^d)]$	$-(1 - 4x)$
$\tilde{\gamma}$	$-2\rho[\frac{1}{2} + T_{3R}^e][T_{3R}^u + T_{3R}^d - \frac{2}{3}x]$	$\frac{2}{3}x$
$\tilde{\delta}$	$-2\rho[-\frac{1}{2} + T_{3R}^e + 2x][T_{3R}^u + T_{3R}^d]$	0

Table 2.1 Neutral-current parameters in the SU(2) $\times$ U(1) model. ($x = \sin^2\theta_w$)

Experiment	Sample of $\nu_\mu + N \rightarrow \mu^- + X$	$\nu_\mu e$ candidates	Background	σ/E [a]
GGM CERN-PS		1	0.3 ± 0.1	< 3 (90% c.l.)
AP Counter exp.		32	20.5 ± 2.0	1.1 ± 0.6
GGM CERN-SPS	64,000	9	0.5 ± 0.2	2.4 $^{+1.2}_{-0.9}$
CB FNAL 15'	83,700	8	0.5 ± 0.5	1.8 ± 0.8
CHARM Counter exp.	56,000	11	4.5 ± 1.4	2.6 ± 1.6
Average of the experiments				1.6 ± 0.4
Prediction of the standard model ($\sin^2\theta_W = 0.23$				1.5

a) in units of 10^{-43} cm^2/GeV.

Table 3.1 Data on $\nu_\mu e$ scatterings

Experiment	Sample of $\bar{\nu}_\mu + N \to \mu^+ + X$	$\bar{\nu}_\mu e$ candidates	Background	σ/E [a]
GGM CERN-PS		3	0.4 ± 0.1	$1.0 \begin{array}{c}+ 2.1 \\ - 0.9\end{array}$
AP Counter exp.		17	7.4 ± 1.0	2.2 ± 1.0
GGM CERN-SPS	7400	0	< 0.03	< 2.7 (90% c.l.)
FMMS FNAL 15'	8400	0	0.2 ± 0.2	< 2.1 (90% c.l.)
BEBC TST CERN-SPS	7500	1	0.5 ± 0.2	<3.4 (90% c.l.)
Average of the experiments [b]				1.3 ± 0.6
Prediction of the standard model ($\sin^2\theta_W$ = 0.23)				1.3

Table 3.2 Data on $\bar{\nu}_\mu e$ scatterings.

a) In units of 10^{-42} cm^2/GeV.

b) This average is obtained by adding the number of events observed in the experiments and dividing by the sum of the effective antineutrino fluxes.

		A(A')	B(B')	C(C')	D(D')
Δ	Production	✓	✓	X	X
$\left(\overset{-}{\nu}\right)_p$	elastic	✓	X	X	✓(?)
D	disintegration	✓	X	X	X

Table 4.1

647nm	Novosibirsk (Barkov and Zolotorev 1979)	-20.6 ± 3.2
	Oxford (Baird 1980)	-10.7 ± 1.5
	Theory (Novikov et al., 1976)	-17
	Theory (Sandars 1980)	-13
	Theory (Martensson et al., 1981)	-11.1
876nm	Seattle (Hollister et al. 1980)	-9.3 ± 2.9
	Theory (Novikov et al. 1976)	-13
	Theory (Sandars 1980)	-11
	Theory (Martensson et al., 1981)	-8.3

Table 5.1   The bismuth rotation parameter R in units of 10^{-8}.
The theoretical calculations are for $\sin^2\theta_W = 0.23$.

Table 7.1 Neutral-Current Parameters

		Model independent	Factorization dependent	Standard Model	$\sin^2\theta_W = 0.23$
νq	α	$\pm(0.589 \pm 0.067)$	0.589 ± 0.067	$1-2\,\sin^2\theta_W$	0.54
	β	$\pm(0.937 \pm 0.062)$	0.937 ± 0.062	1	1
	γ	$\mp(0.273 \pm 0.081)$	$-(0.273 \pm 0.081)$	$-\frac{2}{3}\,\sin^2\theta_W$	-0.153
	δ	$\pm(0.101 \pm 0.093$	0.101 ± 0.093	0	0
νe	g_V	0.06 ± 0.08 or -0.52 ± 0.06	0.043 ± 0.063	$-\frac{1}{2}(1-4\,\sin^2\theta_W)$	-0.04
	g_A	-0.52 ± 0.06 or 0.06 ± 0.08	-0.545 ± 0.056	$-\frac{1}{2}$	-0.5
eq	$\tilde{\alpha}$	$-0.67 \quad 0.19$	-0.68 ± 0.19	$-(1-2\,\sin^2\theta_W)$	-0.54
	$\tilde{\beta}$	——	0.06 ± 0.21	$-(1-4\,\sin^2\theta_W)$	-0.08
	$\tilde{\gamma}$	0.22 ± 0.12	0.24 ± 0.10	$\frac{2}{3}\,\sin^2\theta_W$	0.153
	$\tilde{\delta}$	——	0.00 ± 0.02	0	0
eq	$\tilde{\alpha} + \frac{1}{3}\tilde{\gamma}$	-0.06 ± 0.16	-0.60 ± 0.16	$-\left(1 - \frac{20}{9}\,\sin^2\theta_W\right)$	-0.489
	$\tilde{\beta} + \frac{1}{3}\tilde{\delta}$	0.31 ± 0.51	0.06 ± 0.21	$-(1-4\,\sin^2\theta_W)$	-0.08
ee $e\mu$	h_{VV}	0.01 ± 0.08	0.02 ± 0.04	$\frac{1}{4}(1-4\,\sin^2\theta_W)^2$	0.0016
	h_{AA}	0.18 ± 0.16	0.35 ± 0.11	$\frac{1}{4}$	0.25
	h_{VA}	——	$-0.27 \quad 0.04$	$\frac{1}{4}(1-4\,\sin^2\theta_W)$	0.02

$\overset{(-)}{\nu}_\mu + N \to \overset{(-)}{\nu}_\mu + \text{any}$	$\begin{cases} 0.229 \pm 0.09(\pm 0.005) \\ 0.230 \pm 0.023^* \end{cases}$
$\nu_\mu + p \to \nu_\mu + p$	0.26 ± 0.06
$\nu_\mu + N \to \nu_\mu + N + \pi^o$	0.22 ± 0.09
$\bar{\nu}_\mu + N \to \bar{\nu}_\mu + N + \pi^o$	$0.15 - 0.52$
$e^- + d \to e^- + \text{any}$	0.224 ± 0.020
$\bar{\nu}_\mu + e^- \to \bar{\nu}_\mu + e^-$	$0.23 \, {}^{+0.09}_{-0.23}$
$\nu_\mu + e^- \to \nu_\mu + e^-$	$0.22 \, {}^{+0.08}_{-0.05}$
$\bar{\nu}_e + e^- \to \bar{\nu}_e + e^-$	0.29 ± 0.05
$e^+ + e^- \to e^+ + e^-, \; \mu^+ + \mu^-$	0.25 ± 0.15

Table 7.2   Various ways to determine $\sin^2\theta_W$.

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*

from the Paschos-Wolfenstein relation
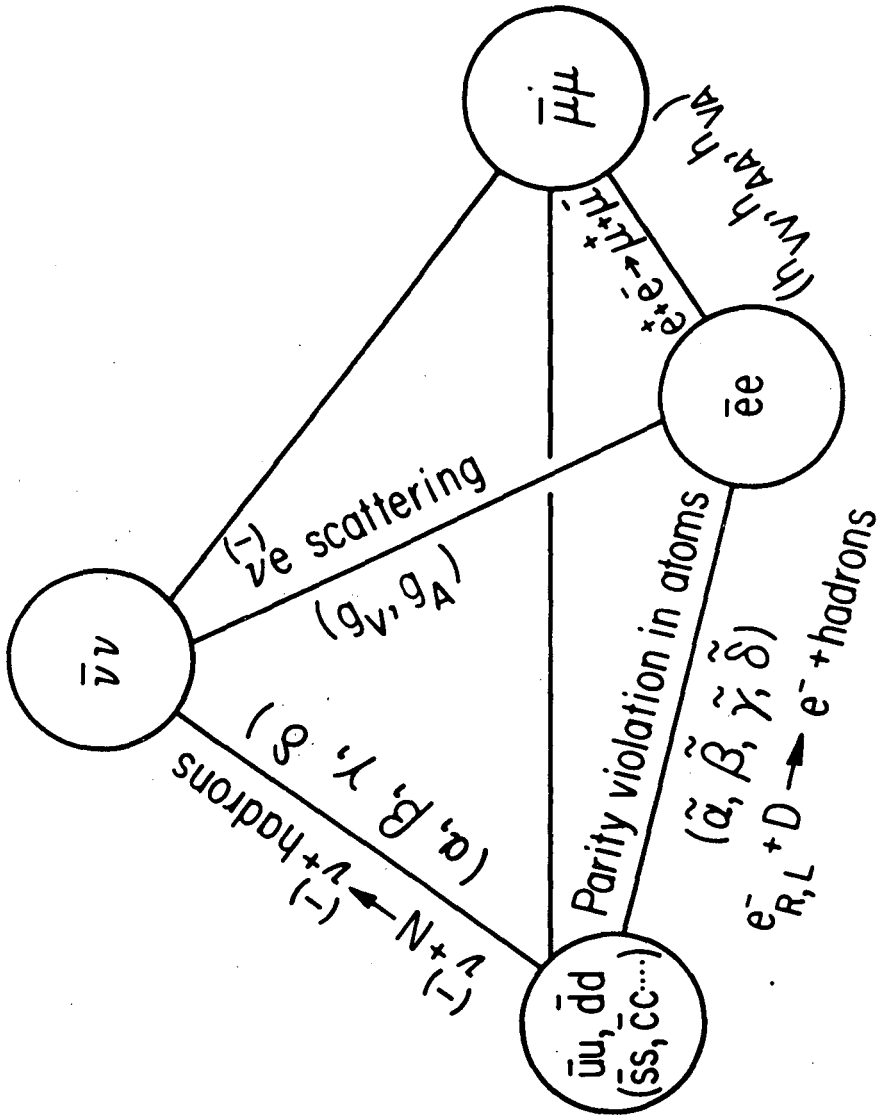
Figure 2.1
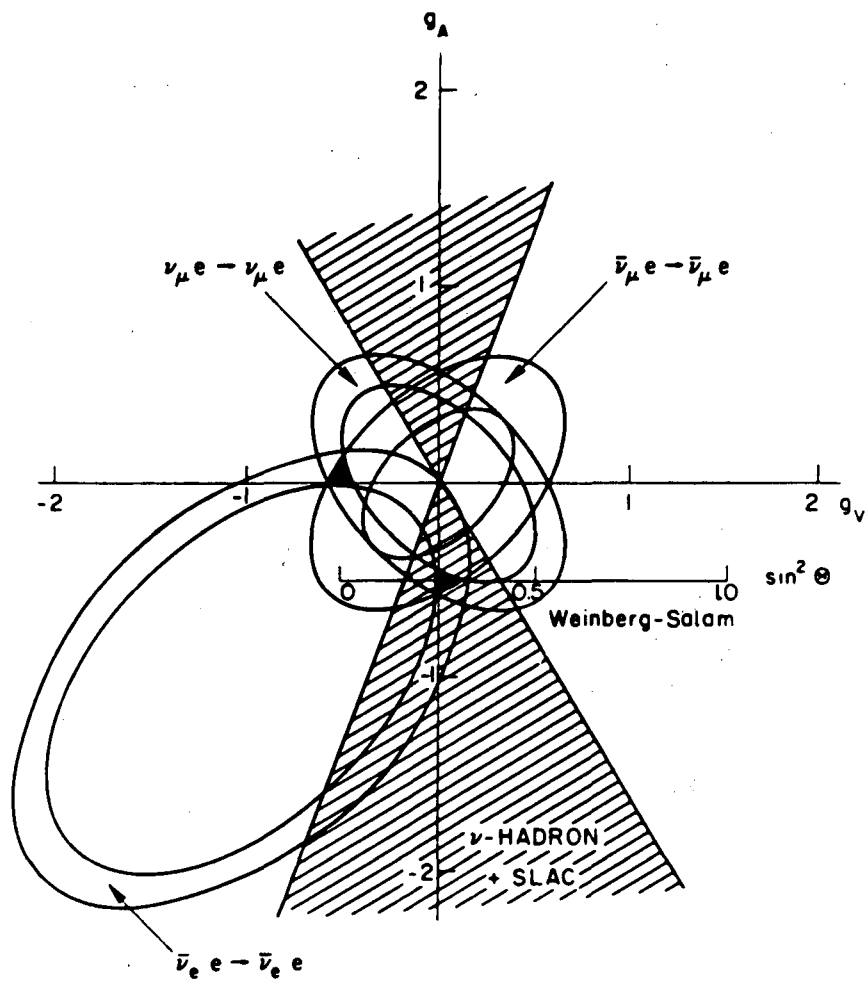
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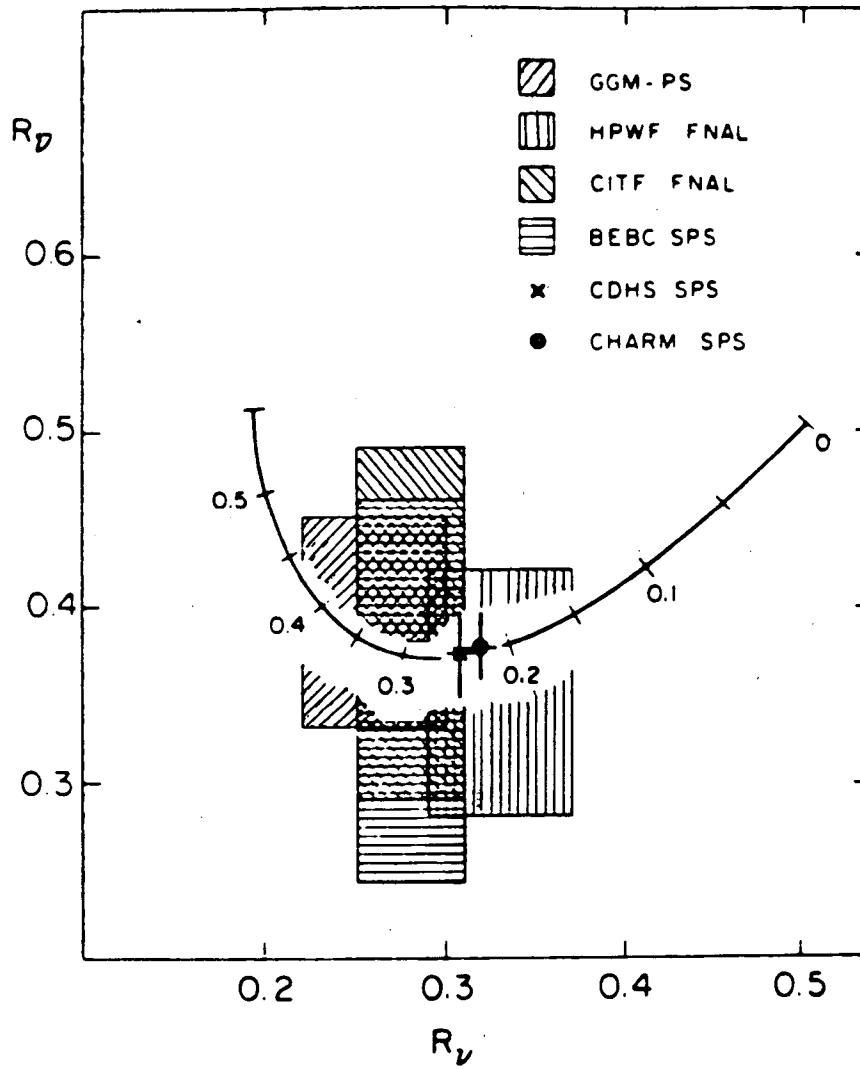
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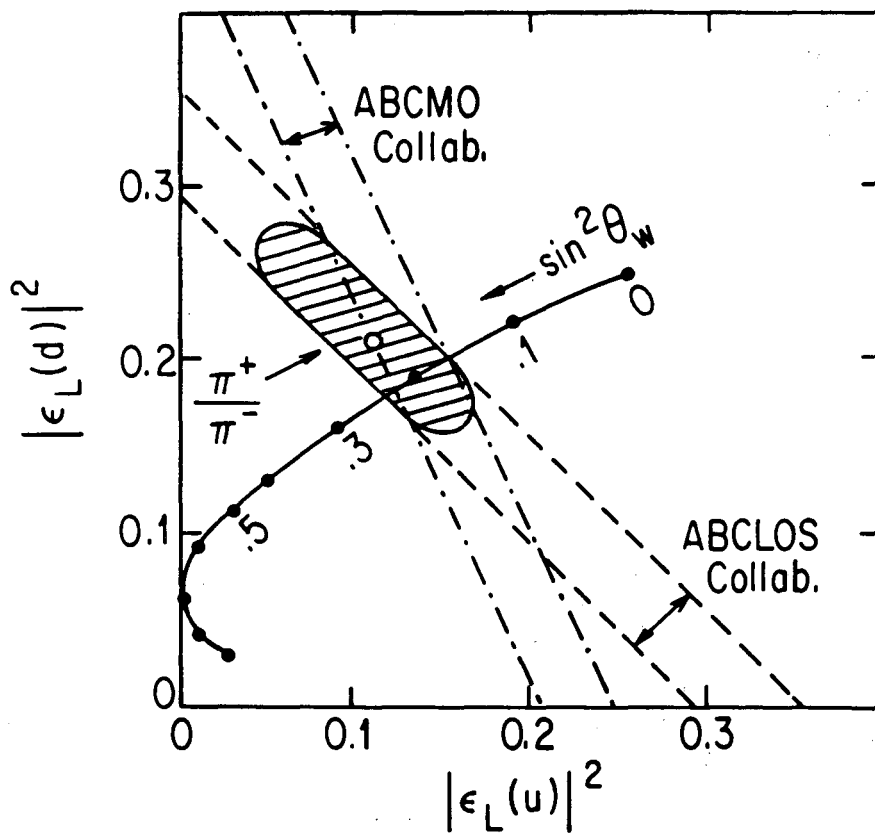
Figure 4.2
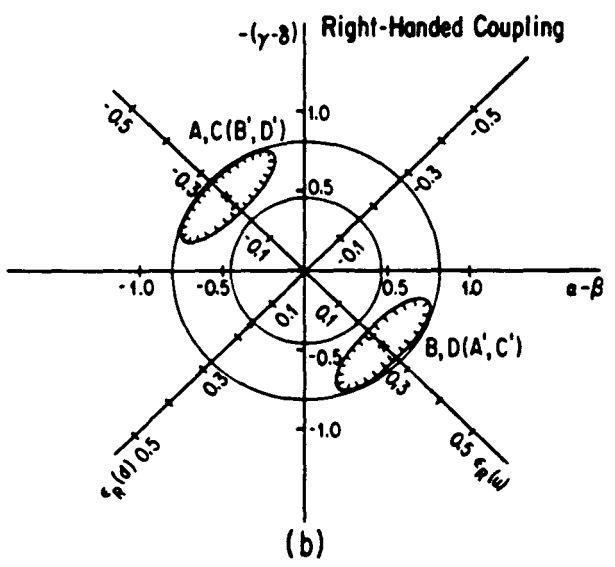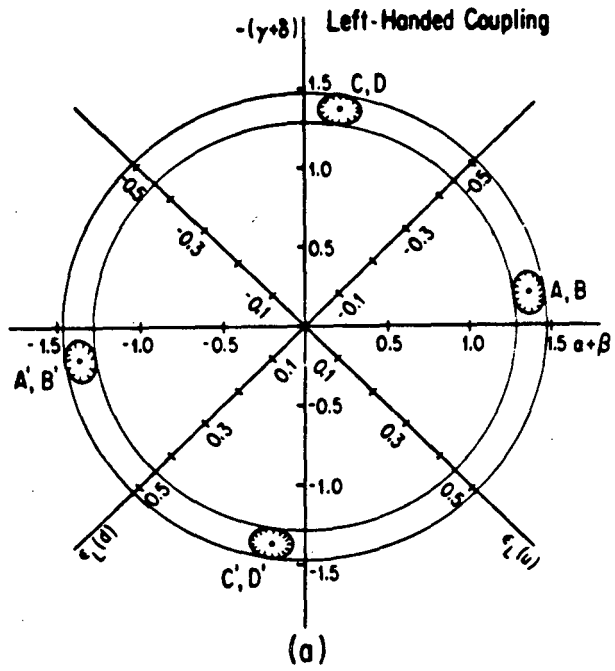
Figure 4.3

Figure 4.4

(a)

(b)
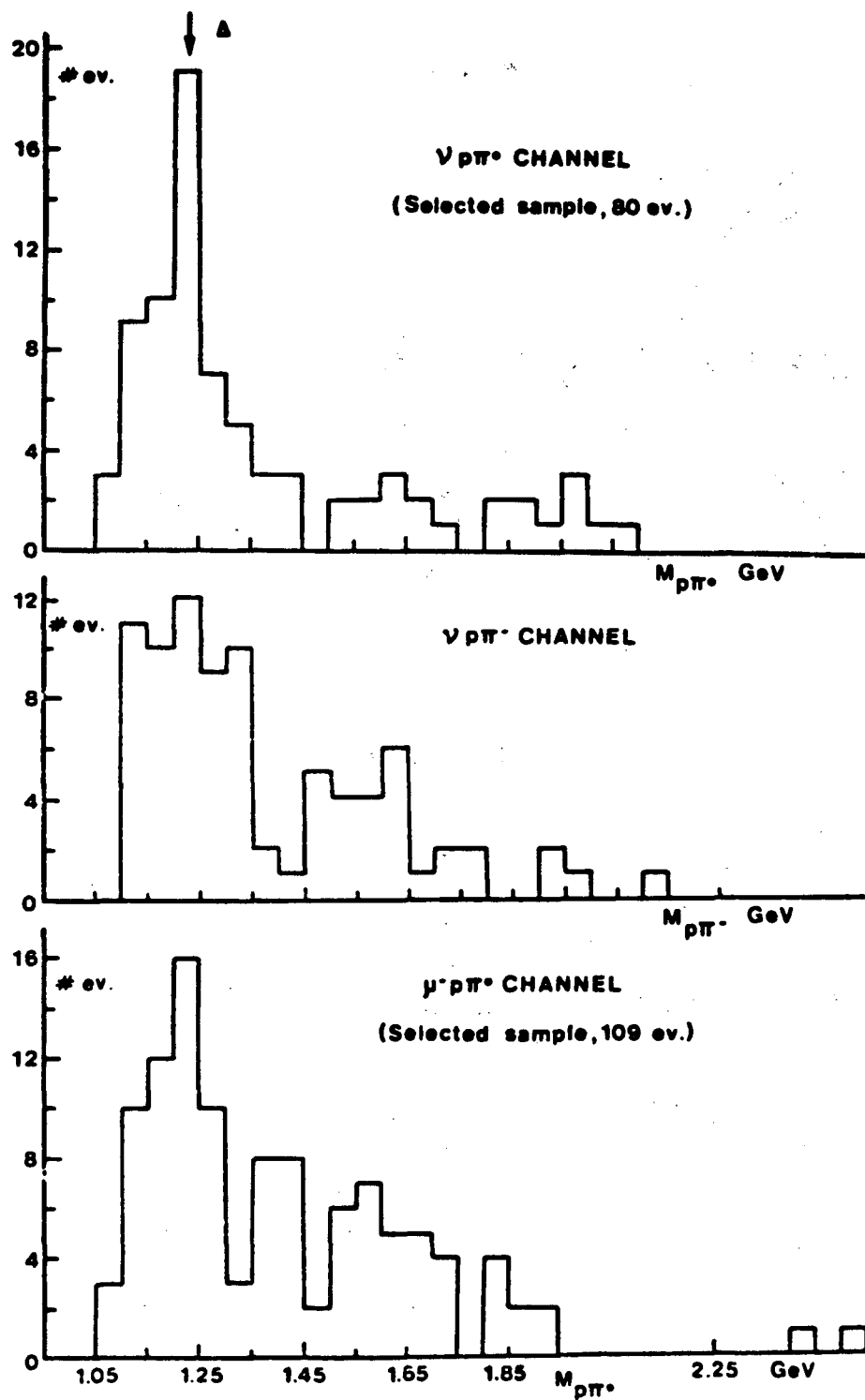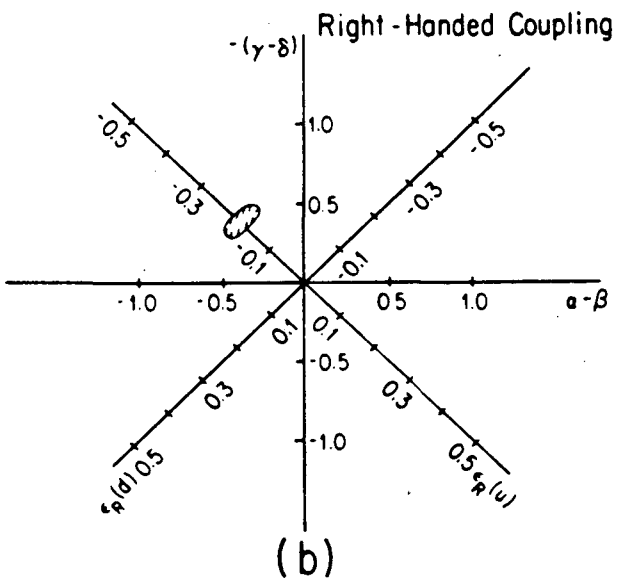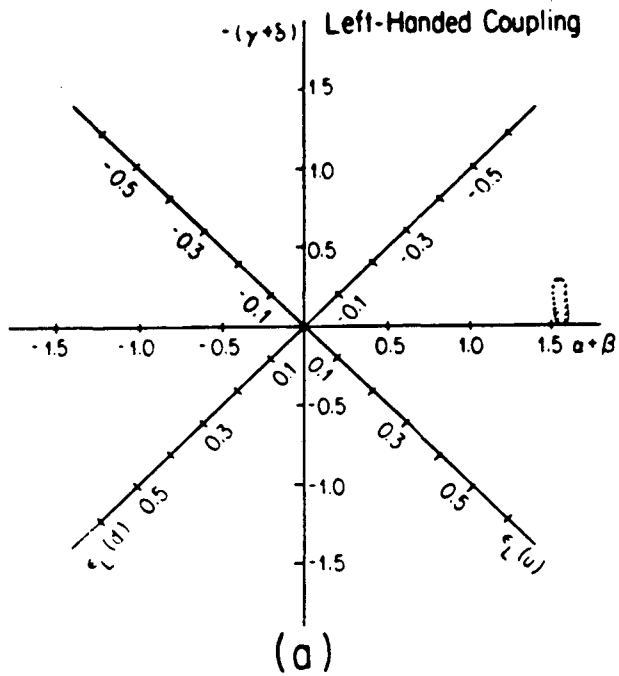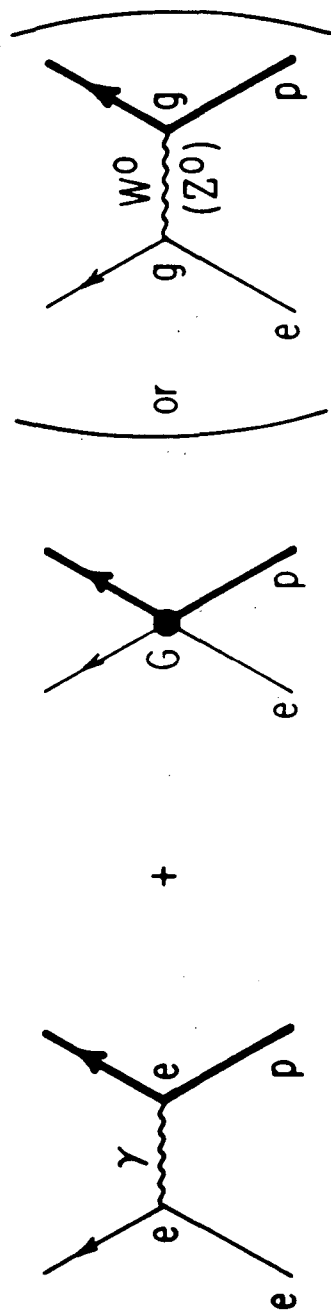
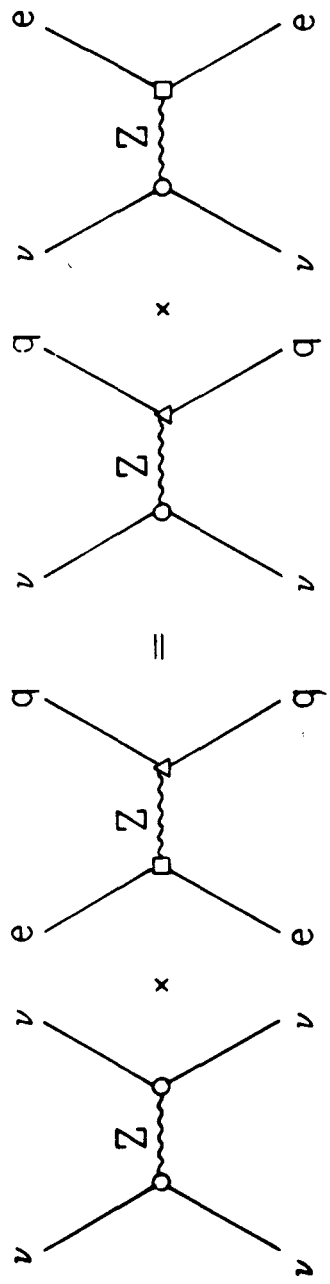Figure 4.5

Figure 5.1
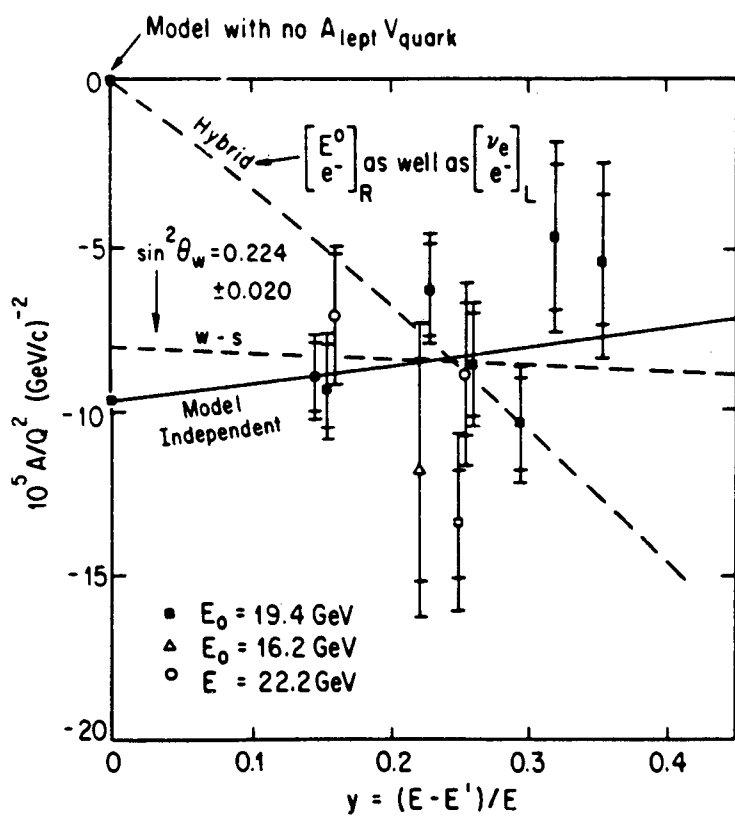
Figure 5.2

Figure 5.4

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Figure 5.5

Figure 6.1

Figure 6.2

Figure 7.1