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ABSTRACT

Gauge Models with Right-handed Currents

by

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We discuss gauge models of weak interactions with the new right-handed charmed current (R.H.C.C.) and $\Delta I = \frac{1}{2}$ rule, based on SU(2) x U (1) or $SU(2)_A \times SU(2)_B \times SU(4)$ gauge groups.

We verify the compatibility of these models with constraints derived from low energy weak interactions, and analyse some of the predictions of the R.H.C.C. concerning new phenomena. In particular, it is shown that the smallness of the K_L-K_S mass difference can be accomodated within these models and that a good fit for both s and p wave hyperon decays can be obtained. The structure of the neutral currents is analysed in detail and we point out the most important features of the various models, by contrasting them with the Weinberg-Salam model. The various decays of the "charmed" pseudo-scalar mesons to normal hadrons are analysed in models with the most general structure for the right-handed currents. Implications of the R.H.C.C. for the Spear Charm Search are analysed.

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Chapter I

Introduction

1.1 The Left-Right Symmetry

In the original Weinberg-Salam[1] unified theory of weak and electromagnetic interactions there is a striking asymmetry in the treatment of left-handed and right-handed fermions. While all the left-handed fermions are placed in SU(2) doublets and thus couple to the charged vector meson, the right-handed ones are put in singlets and therefore do not participate in the charged weak interactions.

However, it remains to be seen if this asymmetry corresponds to a fundamental law of nature, valid at all energies or if it is, on the contrary, just a prejudice of the low and medium energy interactions (or perhaps of the "light fermions"). It is then clear that in order to recover the left-right symmetry one has to introduce right-handed charged currents. These currents were in fact introduced in the theory of weak interactions in two distinct ways:

1) One was through the works of Pati, Salam and Mohapatra[2] who conjectured that nature is intrinsically left-right symmetric, but that in order to see this symmetry one has to go to very high energies. This idea can be implemented in the framework of unified models by postulating that the right-handed fermions also couple to charged vector mesons. More specifically, in the context of the group $SU(2)_L \times SU(2)_R \times U(1)$ (or a larger group containing this one as a subgroup), the left-right symmetry can be introduced by requiring equal bare coupling constants g_L, g_R. The observed left-right asymmetry would then be attributed to the heavier masses of the right-handed gauge mesons $W_R^{\pm}$, as compared to the left-handed gauge mesons $W_L^{\pm}$. One expects then the predominance of V–A over V+A couplings to disappear at energies

$E >> m_{W_R}$. It has been further shown that this left-right discrete symmetry can be maintained as a natural symmetry, provided that it is not broken in the Lagrangian except through Higgs meson masses.

Apart from their aesthetical beauty, left-right symmetric theories have the advantage of providing some understanding of the origin of the breakdown of P and CP invariance.[3] Namely, in left-right symmetric theories it is possible to construct models in which P and CP are exact symmetries of the Lagrangian, only violated through spontaneous symmetry breaking.

2) Another development was originated in the work of Mohapatra[4] who has shown that it is possible to introduce CP violation in the context of a simple SU(2) x U(1) gauge group and without proliferation of new particles, provided that a new right-handed charmed current (R.H.C.C.) is introduced. As far as CP violation is concerned, this R.H.C.C. could be of either $(\bar{\chi} n)_R$ type or $(\bar{\chi} \lambda)_R$ type, χ being the charmed quark. Note that this R.H.C.C. is coupled to the same charged vector boson as the usual weak currents and therefore it manifests itself with strength equal to the conventional V-A currents, even at low energy.

It has been recently pointed out by De Rujula, Giorgi and Glashow[5] that the Mohapatra charmed current has the virtue of providing an elegant explanation for the well extablished $\Delta I = \frac{1}{2}$ rule, together with some interesting predictions such as production of dilepton final states of the type $\mu^- \mu^-$ along with $\mu^+ \mu^-$ in deep inelastic neutrino scattering, equal fractions of strange and non-strange final states above charm threshold in $e^+ e^-$ etc. All these predictions seem qualitatively in accord with present experiments and have stimulated the construction of new models[6] with right-handed charged currents.

These models consisted essentially in generalizing Mohapatra's idea by allowing the right-handed parts of all the quarks to participate in the charged weak interactions. Since constraints on low energy weak interactions forbid two right-handed light fermions to sit on the same doublet, one is at once led to the introduction of new heavy quarks, and then heavy leptons by lepton-quark symmetry.

1.2 Gauge Models with Right-handed Currents

For future reference, we will present here some of the gauge models with right-handed currents that have been proposed.

Pati-Salam Model

This was the first gauge model[2] in which all the elementary fermions (i.e. quarks and leptons) were treated on an equal basis, having four colors and four flavors and the lepton number being the fourth color. Their original model is based on $G \equiv SU(4)_L \times SU(4)_R \times SU(4')$ and the fermions are put in two 16-plets:

$$F_{L,R} = \begin{bmatrix} p_a & p_b & p_c & p_d \equiv \nu \\ n_a & n_b & n_c & n_d \equiv e^- \\ \lambda_a & \lambda_b & \lambda_c & \lambda_d \equiv \mu^- \\ \chi_a & \chi_b & \chi_c & \chi_d \equiv \nu_\mu \end{bmatrix}$$

transforming as $F_L \equiv (4, 1, \bar{4})$, $F_R \equiv (1, 4, \bar{4})$, and where a, b, c, d, represent the four colors.

The phenomenological implications of this model have been extensively analyzed. Most of the essential features of the model remain unchanged if one works with a subgroup of G, like $SU(2)_L \times SU(2)_R \times SU(4')$, or $SU(2)_L \times SU(2)_R \times U(1)$.

In principle, the left-right symmetry manifests itself in these models only at very high-energy (i.e. $E \gg M_{W_R}$). However, we will point out in the next chapter that within the framework of $SU(2)_L \times SU(2)_R \times U(1)$ model the left-right symmetry may become apparent, even at present energies through the structure of the neutral currents.

Mohapatra Model

This was the first model[4] in which R.H.C.C. were introduced and it was motivated by the need to incorporate the observed CP violation in the framework of a unified gauge model. The assignment of fermions is as follows:

$$\begin{pmatrix} p_L \\ n_L(\theta_c) \end{pmatrix} \quad \begin{pmatrix} \chi_L \\ \lambda_L(\theta_c) \end{pmatrix} \quad \begin{pmatrix} \chi_R \\ n_R(\phi) \end{pmatrix} \qquad (\tfrac{1}{2}\ 1)$$

$$\begin{pmatrix} \nu_L \\ e^-_L \end{pmatrix} \quad \begin{pmatrix} \nu_{\mu L} \\ \mu^-_L \end{pmatrix} \quad \begin{pmatrix} E^\circ_R \\ e^-_R \end{pmatrix} \qquad (\tfrac{1}{2}\ {-1})$$

where

$$\begin{aligned} n_R(\phi) &= n_R \cos\phi + i\lambda_R \sin\phi \\ \lambda_R(\phi) &= i n_R \sin\phi + \lambda_R \cos\phi \end{aligned} \qquad (1.1)$$

and ϕ is chosen to be of the order of 10^{-4}, to reproduce the observed magnitude of CP violation. The remaining fermions, $E^\circ_L \cdot \mu^-_R \cdot p_R\ \lambda_R(\phi)$ are SU(2) singlets. The new heavy lepton E° is introduced in order to make the theory free of triangle anomalies.

As it has been pointed out in ref. 4, CP violation and low energy phenomenology are correctly reproduced if we interchange $n_R \leftrightarrow \lambda_R(\phi)$.

placing n_R in a singlet and $\lambda_R(\phi)$ in a doublet. However, the two alternatives lead to quite different structures for the neutral currents. Namely when the doublet is

$$\begin{pmatrix} \chi_R \\ \lambda_R(\phi) \end{pmatrix}$$

the structure of the hadronic neutral currents involving p and n quarks is the same as in the Weinberg model. On the other hand, when n_R is placed in a doublet, the structure of the neutral current changes drastically. In particular this choice leads to an asymmetry in the assignment of p_R, n_R quarks, and this asymmetry manifests itself through the appearance of an axial isoscalar current, which is a unique feature of this model.

Harvard Model

This model has been proposed by De Rujula et al[5] and suggests an interesting way of incorporating the Mohapatra R.H.C.C. in a unified gauge model and providing thus a new and simple explanation for the observed enhancement of $\Delta I=\frac{1}{2}$ non-leptonic weak processes. There is in this model a proliferation of quarks. motivated in part by the experimental value of $R=\sigma(\text{hadronic})/\sigma(\mu^+\mu^-)$. Recall that R equals approximately 5 at center of mass energies of 5 or 6 GeV while one expects $R=10/3$ in the standard model with 4 quarks and no extra leptons.

Among the various theoretical possibilities, De Rujula et al have opted for a vector-like theory, in which all fermion fields are placed in doublet representations of SU(2). A theoretical motivation for this choice comes from the fact that vector-like theories are always free of amomalies.

Obviously, a model like this, requires at least six quarks, the two extra quarks being the partners of p_R and λ_R in the two new doublets. The assignment of quarks is as follows:

$$\begin{pmatrix} p \\ \\ n_\theta \end{pmatrix}_L ; \begin{pmatrix} c \\ \\ \lambda_\theta \end{pmatrix}_L ; \begin{pmatrix} t \\ \\ b \end{pmatrix}_L ; \begin{pmatrix} p \\ \\ b \end{pmatrix}_R ; \begin{pmatrix} c \\ \\ n \end{pmatrix}_R ; \begin{pmatrix} t \\ \\ \lambda \end{pmatrix}_R$$

where c is the charmed quark, used to implement the GIM mechanism.

Princeton-Caltech Models[8]

This model is very similar to the Harvard Model, being also a vector like theory with 6 quarks. The most important difference is that these authors introduce the doublet:

$$\begin{pmatrix} c \\ \\ \lambda \end{pmatrix}_R \quad \text{instead of} \quad \begin{pmatrix} c \\ \\ n \end{pmatrix}_R$$

As far as the neutral currents are concerned, these models lead to the same structure as the Harvard model, namely, a purely vector neutral current. As we will discuss in chapter IV the best way to distinguish these two models is through the study of charmed hadrons decays. In these models, the $\Delta I=\frac{1}{2}$ rule is not automatically satisfied, and one has to rely on a possible dynamical enhancement.

C.C.N.Y. Models[9]

These models are based on the gauge group $SU(2)_A \times SU(2)_B \times SU(4)'$ and use 8 quarks and 8 leptons. If we take into account the color degree of freedom, one has altogether two sixteenplets of elementary fermions:

$$\begin{bmatrix} p_a & p_b & p_c & p_d \equiv \nu \\ n_a & n_b & n_c & n_d \equiv e^- \\ \lambda_a & \lambda_b & \lambda_c & \lambda_d \equiv \mu^- \\ \chi_a & \chi_b & \chi_c & \chi_d \equiv \nu' \end{bmatrix}$$

$$\begin{bmatrix} p'_a & p'_b & p'_c & p'_d \equiv E^{\circ} \\ n'_a & n'_b & n'_c & n'_d \equiv E^- \\ \lambda'_a & \lambda'_b & \lambda'_c & \lambda'_d \equiv M^- \\ \chi'_a & \chi'_b & \chi'_c & \chi'_d \equiv M^{\circ} \end{bmatrix}$$

where $p'_a, p'_b, p'_c, \ldots$ are the heavy (mirror) counterparts of $p_a \cdot p_b \cdot p_c \cdot$ and $E^{\circ}, E^-, M^{\circ}, M^-$ are heavy leptons. The assignment of fermions is as follows:

Model I

Quarks:

$$\begin{pmatrix} p_L \\ n_{L(\theta)} \end{pmatrix} \begin{pmatrix} \chi_L \\ \lambda_{L(\theta)} \end{pmatrix} \begin{pmatrix} \chi_R \\ n_R \end{pmatrix} \quad \begin{pmatrix} p'_R \\ \lambda_R \end{pmatrix} \qquad (\tfrac{1}{2}, 0, 4)$$

$$\begin{pmatrix} p_R \\ n'_R \end{pmatrix} \begin{pmatrix} \chi'_R \\ \lambda'_R \end{pmatrix} \begin{pmatrix} p'_L \\ n'_L \end{pmatrix} \quad \begin{pmatrix} \chi'_L \\ \lambda'_L \end{pmatrix} \qquad (0, \tfrac{1}{2}, 4)$$

$$\begin{pmatrix} \nu_L \\ e^-_L \end{pmatrix} \quad \begin{pmatrix} \nu_\mu(a)_L \\ \mu^-_L \end{pmatrix} \quad \begin{pmatrix} E^o_R \\ e^-_R \end{pmatrix} \quad \begin{pmatrix} M^o_R \\ M^-_R \end{pmatrix} \qquad (\tfrac{1}{2}\ 0.\ 4)$$

Leptons:

$$\begin{pmatrix} \nu_R \\ E^-_R \end{pmatrix} \quad \begin{pmatrix} M^o_L(a) \\ M^-_L \end{pmatrix} \quad \begin{pmatrix} \nu_{\mu R} \\ \mu^-_R \end{pmatrix} \quad \begin{pmatrix} E^o_L \\ E^-_L \end{pmatrix} \qquad (0, \tfrac{1}{2}\ 4)$$

This 32-plet of fermions can be used to unify all non-gravitational forces (weak, electromagnetic and strong) by using a larger gauge group, such as SU(4) x SU(4'). One obtains then only one coupling constant by imposing the discrete symmetry prime $\longleftrightarrow$ unprime. The introduction of the heavy fermions finds then a theoretical justification[2], since this unification is anomaly free only if the new 16-plet of mirror fermions is incorporated in the unified model. CP violation can be accounted for by introducing a complex phase in n_R, λ_R, as it is expressed in (1.1). It is then possible to establish a connection between CP violation and the magnitude of the $\chi - \chi'$ mass difference. In this framework, CP would be conserved in the symmetry limit $m_\chi = m_{\chi'}$ and the interactions which violated this symmetry would be also responsible for the observed CP violation. It will be shown in chapter III that the effective Δ S=2 Lagrangian receives a new contribution from the right-handed currents, which looks like:

$$\left[\mathcal{L}_{eff}\right]^{\Delta S=2}_{LR} = G^{\Delta S=2}_{LR} \left\{ 4\bar{\lambda}\,\tfrac{1}{2}(1-\gamma_5)n\,\bar{\lambda}\,\tfrac{1}{2}(1-\gamma_5)n \; + \right.$$
$$\left. + \bar{\lambda}\,\sigma_{\mu\nu}\,\tfrac{1}{2}(1-\gamma_5)n\,\bar{\lambda}\,\sigma^{\mu\nu}\tfrac{1}{2}(1-\gamma_5)n + h.c. \right\}$$

It turns out that in this model $G_{LR}^{\Delta S=2}$ is of order $G_F^2 m_\chi^2 \cos^2\theta_c$ and thus quite larger than in the standard model, due to the absence of Cabbibo suppression. This has led some authors to suggest[9] that these models are in contradiction with the observed value of K_L-K_S mass difference. It will be shown in chapter III that this conclusion is based on a unjustified approximation and therefore should not be taken seriously. However, we will present an alternative model, based also on the $SU(2)_A \times SU(2)_B \times SU(4)'$ gauge group. in which this question does not even arise. The lepton sector is unchanged and the assignment of quarks is as follows:

Model II

Quarks:

$$\begin{pmatrix} p_L \\ n_L(\theta_c) \end{pmatrix} \begin{pmatrix} \chi_L\cos\alpha + p'_L\sin\alpha \\ \lambda_L(\theta_c) \end{pmatrix} \qquad \begin{pmatrix} \chi_R\cos\beta + p'_R\sin\beta \\ n_R(\phi) \end{pmatrix} \begin{pmatrix} \chi'_R \\ \lambda_R(\phi) \end{pmatrix} \quad (\tfrac{1}{2}, 0, 4)$$

$$\begin{pmatrix} p_R \\ n'_R \end{pmatrix} \begin{pmatrix} -\chi_R\sin\beta + p'_R\cos\beta \\ \lambda'_R \end{pmatrix} \qquad \begin{pmatrix} -p'_L\cos\alpha + \chi_L\sin\alpha \\ n'_L \end{pmatrix} \begin{pmatrix} \chi'_L \\ \lambda'_L \end{pmatrix} \quad (0\ \tfrac{1}{2}\ 4)$$

It is clear that when the mixing angles α, β are small, the $\Delta I=\frac{1}{2}$ rule will be satisfied. On the other hand, one obtains for $G_{LR}^{\Delta S=2}$:

$$G_{LR}^{\Delta S=2} \simeq \frac{G_F^2}{4\pi^2} \left(m_\chi\cos\alpha\cos\beta + m_{p'}\sin\alpha\sin\beta \right)^2$$

One can then reduce the strength of $\left[\mathcal{L}_{\text{eff}}\right]_{LR}^{\Delta S=2}$, by choosing $(m_\chi\cos\alpha\cos\beta + m_{p'}\sin\alpha\sin\beta)$ small enough.

Finally, we note that in the leptonic sector, a mixing of ν_μ, M_L^o

was introduced by defining:

$$\nu_\mu(\alpha) = \cos\alpha\, \nu_\mu + \sin\alpha\, M_L^0$$

$$M_L^0(\alpha) = -\sin\alpha\, \nu_\mu + \cos\alpha\, M_L^0$$

This was done in order to provide a possible explanation for the Kolar-mine events.[11] The observed tracks would be the result of electromagnetic production of $M^\pm$ pairs by cosmic rays, followed by decay into M° which in turn decays into μ^+ + hadrons, with a long life-time.

$$\gamma \longrightarrow M^+ M^-$$

$$M^- \xrightarrow{G_F} M^0 \mu^- \bar{\nu}_\mu$$

$$M^0 \xrightarrow{G_F \sin\alpha} \mu^- + \text{hadrons}$$

The mixing angle α is of order 10^{-3} to avoid conflict with μ-decay data and at the same time give a long life-time to E°.

1.3 <u>Viability of Gauge Models with R.H.C.C.</u>

In the next chapters we will verify the compatability of the new R.H.C.C. with the observed weak interactions and at the same time investigate some of its consequences concerning new phenomena. In chapter II we analyze the structure of the neutral current in the various gauge models and look for constraints imposed on them by present experimental data. In chapter III, we show how the smallness of K_L-K_S mass difference can be accomodated in the new models. Chapter IV will be devoted to the study of the non-leptonic decays of the new "charmed hadrons," that are predicted by these gauge models. In particular, we will analyze the consequences of the new charmed current to the Spear Charm Search, and we will show that it is easier to understand the present non-observation of particles carrying net charm, in the framework of models with R.H.C.C. In chapter V we analyze the changes introduced by the new current in the structure of the $\Delta S=1$ non-leptonic Hamiltonian and we discuss in detail its consequences for hyperon and K decays. Finally in chapter VI our conclusions are drawn.

Chapter II

Neutral Currents

2.1 Introductory Remarks

Neutral currents have been a subject of great interest in the theory of weak interactions, long before unified gauge models of weak and electromagnetic interactions were proposed by Weinberg and Salam. However, these theoretical developments played a crucial role in stimulating the various experimental investigations which ultimately led to the discovery of the neutral currents. Although the existence of neutral currects in itself does not constitute a proof of the correctness of the unified gauge models, the knowledge of the detailed structure of the neutral current is of great interest in distinguishing between different models and thus in narrowing down the number of realistic candidates for unified gauge models.

In this chapter we will analyze the structure of the neutral currents in various models with R.H.C.C. and we will discuss some of the experiments that can test them.

In analogy with the effective current-current interaction for charged currents, we define the effective neutral current interaction as

$$\mathcal{L}_{eff} = \frac{G_F}{\sqrt{2}} J_\mu J^\mu \quad , \tag{2.1}$$

where

$$J_\lambda = \frac{1}{2}\bar{\nu}\gamma_\mu(1+\gamma_5)\nu + \bar{e}\gamma_\mu(g_V+g_A) + v_3 V^3_\mu + a_3 A^3_\mu + \\ + v_0 V^0_\mu + a_0 A^0_\mu . \tag{2.2}$$

2.2 Structure of the neutral currents in models with right-handed charged currents

2.2.1 Left-right symmetric model based on $SU(2)_L \times SU(2)_R \times U(1)$

We illustrate here, in the framework of a simple model, a way in which the left-right symmetry may manifest itself at low energies. For simplicity, we will work with a $SU(2)_L \times SU(2)_R \times U(1)$ model. The fermions are put in left-right symmetric doublets:

Leptons	Representation
$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L \quad \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$	$(\frac{1}{2}, 0, -1)$
Quarks	
$\begin{pmatrix} p \\ n_c \end{pmatrix}_L \quad \begin{pmatrix} \chi \\ \lambda_c \end{pmatrix}_L$	$(\frac{1}{2}, 0, 1)$

And similarly for the right-handed sector with $L \to R$ and $(\frac{1}{2}, 0, Y) \to (0, \frac{1}{2}, Y)$. In order to preserve the left-right symmetry, we put $g_L = g_R = g$ and define the electric charge by:

$$Q = T_{3L} + T_{3R} + \frac{Y}{2} \tag{2.3}$$

The electromagnetic field is then fixed and given by:

$$A_\mu = \cos\theta B_\mu + \frac{\sin\theta}{\sqrt{2}} W_L^3 + \frac{\sin\theta}{\sqrt{2}} W_R^3 , \tag{2.4}$$

where $\theta = \tan^{-1} \frac{\sqrt{2}g'}{g}$. The electromagnetic field defines a direction in B_μ, W_L^3, W_R^3 space. The two neutral gauge bosons will be in a plane orthogonal to this direction. If we define

$$T_1 = \frac{1}{\sqrt{2}}(W_L^3 - W_R^3) \qquad T_2 = A_\mu \times T_1 \tag{2.5}$$

Then the two neutral gauge bosons can in general be expressed as:

$$\begin{aligned} Z_\mu &= \cos\gamma T_1 + \sin\gamma T_2 \\ U_\mu &= -\sin\gamma T_1 + \cos\gamma T_2 \end{aligned} \tag{2.6}$$

The structure of the 2 currents coupled to Z_μ, U_μ is then a function of θ, γ and their relative strength depends on the mass ratio M_{Z_μ}/M_{U_μ}. In general, there is of course still another parameter, ϵ, which specifies the over-all strength of the neutral interactions. We will have to choose the Higg's system in such a way that $M_{W_R} \gg M_{W_L}$, for obvious reasons. However, no such constraint exist for the neutral sector and one may very well have $M_{Z_\mu} \simeq M_{U_\mu}$. The point we wish to emphasize is that the left-right symmetry may manifest itself through the neutral sector, even at low energies. For definiteness and in order to illustrate our point, we will now specialize to the case in which the neutral current interactions conserve parity. This can be achieved by choosing $\gamma = \pi/2$ in (2.6). The two eigenstates of the neutral boson mass matrix are then:

$$\begin{aligned} Z_\mu &= -\sin\theta B_\mu + \frac{\cos\theta}{\sqrt{2}} W_L^3 + \frac{\cos\theta}{\sqrt{2}} W_R^3 \\ U_\mu &= -\frac{W_L^3}{\sqrt{2}} + \frac{W_R^3}{\sqrt{2}} \end{aligned} \tag{2.7}$$

From (2.7) one sees that Z_μ couples only to vector currents while U_μ couples only to axial vector currents. Note that the neutral current efefective interaction will receive contributions from Z_μ and U_μ and thus in neutrino-electron interactions for example both g_V and g_A will have non-vanishing values. The fact that parity is conserved distinguishes this model from others with the same value for g_V, g_A. We will have now to choose a Higg's system which is consistent with (2.7). It can be readily seen that a possible choice is:

Higgs's boson	Representation	Vacuum expectation value
$\chi_L \equiv \begin{pmatrix} \chi_L^+ \\ \chi_L^0 \end{pmatrix}$	$\left(\frac{1}{2}, 0, 1\right)$	$\langle \chi_L \rangle = \begin{pmatrix} 0 \\ u \end{pmatrix}$
$\chi_R \equiv \begin{pmatrix} \chi_R^+ \\ \chi_R^0 \end{pmatrix}$	$\left(0, \frac{1}{2}, 1\right)$	$\langle \chi_R \rangle = \begin{pmatrix} 0 \\ u \end{pmatrix}$
$\sigma \equiv \begin{pmatrix} \sigma_1^+ & \sigma_2^+ \\ \sigma_1^0 & \sigma_2^0 \end{pmatrix}$	$\left(\frac{1}{2}, \frac{1}{2}, 0\right)$	$\langle \sigma \rangle = \begin{pmatrix} k & 0 \\ 0 & k' \end{pmatrix}$
$\pi \equiv \begin{pmatrix} \pi^+ \\ \pi^0 \\ \pi^- \end{pmatrix}$	$(0, 1, 0)$	$\langle \pi \rangle = \begin{pmatrix} 0 \\ V \\ 0 \end{pmatrix}$

We see that the left-right symmetry is only broken through π, whose role is to suppress the charged right-handed interactions by choosing V sufficiently large. The neutral sector acquires masses only through χ_L, χ_R, σ

(which respect the left-right symmetry) and this is the reason why the neutral current interactions conserve parity. One can easily verify that Z_μ, U_μ defined in (2.7) are, together with A_μ, the eigenstates of the mass matrix and the mass spectrum is given by:

$$M^2_{Z_\mu} = \frac{1}{4}\left(g^2+2g'^2\right)u^2 \qquad M^2_{W_R^+} = g^2\left(v^2+u^2+k^2+k'^2\right)$$

$$M^2_{U_\mu} = \frac{g^2}{4}\left(u^2+2\left(k^2+k'^2\right)\right) \qquad M^2_{W_L^+} = \frac{g^2}{4}\left(u^2+k^2+k'^2\right) \tag{2.8}$$

from where it follows $M^2_{Z_\mu} = \dfrac{M^2_{W_L}}{\cos^2\theta\,(1+\alpha)}$ and $M^2_{U_\mu} = M^2_{W_L}\left(\dfrac{1+2\alpha}{1+\alpha}\right)$, with $\alpha = \dfrac{k^2+k'^2}{u^2}$. The Fermi constant is related to $M^2_{W_L}$ by $\dfrac{G_F}{\sqrt{2}} = \dfrac{g^2}{8M^2_{W_L}}$.

We can now write the structure of the neutral current as manifested through the scattering of left-handed neutrinos on electron and nucleon targets. The effective interaction can still be written as in (2.1) and one obtains:

$$\begin{aligned} g_V &= -\frac{1}{2}(1+\alpha)\left(1-2\sin^2\theta\right) &;\quad g_A &= -\frac{1}{2}\,\frac{1+\alpha}{1+2\alpha} \\ v_3 &= (1+\alpha)\cos^2\theta &;\quad a_3 &= +\frac{1+\alpha}{1+2\alpha} \\ v_0 V^0_\mu &= -(1+\alpha)\sqrt{3}\sin^2\theta\, V^8_\mu + \tilde{V}^0_\mu &;\quad a_0 A^0 &= \tilde{A}^0 \end{aligned} \tag{2.9}$$

where $\tilde{V}^0$, $\tilde{A}^0$ are isoscalar currents not containing p, n quarks. For $\alpha=0$ one sees that this model coincides with the Weinberg-Salam model, with the correspondence $\sin^2\theta \to 2\sin^2\theta_W$.

2.2.2 Models with R.H.C.C. based on SU(2) x U(1).

We will examine here the structure of the neutral currents in a general class of models having the left-handed sector coinciding with that of Weinberg-Salam model and in which both the right-handed fermion sector and the Higgs system are arbitrary. Our only assumption is that the right-handed fermions have definite weak isotopic spin. The structure of the neutral currents depends only on the value of the third component of weak isotopic spin of the fermions. Furthermore, for most practical purposes (i.e. scattering of neutrinos off electron and nucleon targets) only t^3_e, t^3_p, t^3_n matter. One readily obtains:

$$
\begin{aligned}
g_V &= \varepsilon\left[-\tfrac{1}{2} + 2\sin^2\theta + t^3_{e_R}\right] \quad ; \quad & g_A &= \varepsilon\left[-\tfrac{1}{2} - t^3_{e_R}\right] \\
v_3 &= \varepsilon\left[\left(1 - 2\sin^2\theta\right) + \left(t^3_{p_R} - t^3_{n_R}\right)\right] \quad ; \quad & a_3 &= \varepsilon\left[t^3_{p_R} + t^3_{n_R} + 1\right] \qquad (2.10) \\
v_0 &= \varepsilon\left[-2\sin^2\theta + \left(t^3_{p_R} + t^3_{n_R}\right)\right] \quad ; \quad & a_0 &= \varepsilon\left[t^3_{p_R} + t^3_{n_R}\right]
\end{aligned}
$$

where ε is a parameter which measures the over-all strength of the neutral current interactions, and depends only on the Higgs system we choose. One has in general $\varepsilon \geqslant 1$ with the equal sign corresponding to the simplest Higg's system. The Weinberg model chooses $t^3_{e_R} = t^3_{p_R} = t^3_{n_R} = 0$. The Mohapatra model is the minimal model with R.H.C.C. and corresponds to the choice $t^3_{e_R} = -1/2$, $t^3_{n_R} = -1/2$, $t^3_{p_R} = 0$. If one requires that all fermions should be in doublets, then $t^3_{e_R} = t^3_{n_R} = -1/2$, $t^3_{p_R} = 1/2$, and one is led to the vector-like theory presented in the previous chapter.

From (2.10) one sees that the values of the axial vector parts of the neutral current do not depend on the parameter θ . This makes the experimental determination of g_A, a_3, a_0 extremely important in providing clues for model building. Thus the knowledge of g_A, a_3, a_0 is sufficient to determine uniquely the "correct model". or rule out all gauge models based on SU(2) x U(1). It is also apparent that a_0 is non-vanishing only in models where $t^3_{p_R} + t^3_{n_R} \neq 0$. Besides the Mohapatra model, another possible model possessing a axial isoscalar current corresponds to the choice $t_{p_R} = 1/2$ $t_{n_R} = 0$. This would of course imply the introduction of a new heavy quark and a new doublet:

$$\begin{pmatrix} p_R \\ b_R \end{pmatrix}$$

2.2.3 C.C.N.Y. Models[9]

Here we will analyze the neutral currents in the C.C.N.Y. models described in 1.2. As far as the p-n sector is concerned, the neutral currents in models I, II coincide. For simplicity, we work with the group $SU(2)_A \times SU(2)_B \times U(1)$ and define electric charge by:

$$Q = T_{3A} + T_{3B} + \frac{Y}{2} \, . \qquad (2.11)$$

Obviously, this model shares some of the features of the left-right symmetric model described in 2.2.1. There are however some important differences, that we will point out. In general, the neutral current interactions receive contributions of the two neutral gauge bosons Z_μ, U_μ, defined in (2.6). and the structure of the neutral current depends in general on three parameters $\theta, \varphi, \lambda \equiv \frac{M^2_{Z_\mu}}{M^2_{\mu}}$.

One obtains in general:

$$g_V = \frac{\varepsilon}{2}\Big[\left(\cos\varphi\cos\theta + \sin\varphi\right)\left(-\cos\varphi\cos\theta + \sin\varphi\,(2\sin^2\theta - 1)\right) + $$
$$+ \lambda\left(\cos\varphi - \sin\varphi\cos\theta\right)\left(\cos\theta\sin\varphi + \cos\varphi\,(-1 + 2\sin^2\theta)\right)\Big] ,$$

$$g_A = 0 \quad , \tag{2.12}$$

$$v_3 = \frac{\varepsilon}{2}\Big[\left(\cos\varphi\cos\theta + \sin\varphi\right)\left(\cos\varphi + 2\sin\varphi\cos\theta\right) +$$
$$+ \lambda\left(\cos\varphi - \sin\varphi\cos\theta\right)\left(-\sin\varphi + 2\cos\varphi\cos\theta\right)\Big]\cos\theta ,$$

$$v_0 = -\frac{\varepsilon}{2}\Big[\left(\cos\varphi\cos\theta + 2\sin\varphi\sin^2\theta\right)\left(\cos\varphi\cos\theta + \sin\varphi\right) + \lambda\left(\sin\varphi\cos\theta - 2\cos\varphi\sin^2\theta\right)\cdot$$
$$\cdot\left(\cos\varphi - \sin\varphi\cos\theta\right)\Big]\cos\theta ,$$

$$a_0 = a_3 = \frac{\varepsilon}{2}\Big[\left(\cos^2\varphi\cos\theta + \sin\varphi\cos\varphi\right) + \lambda\left(\sin^2\varphi\cos\theta - \sin\varphi\cos\varphi\right)\Big]\cos\theta .$$

Note that this model has also a nonvanishing axial isoscalar current. This is due to the fact that we placed n_R in a $SU(2)_A$ doublet while p_R is in a $SU(2)_B$ doublet. Had we put both p_R and n_R in $SU(2)_A$ doublets, the axial isoscalar current would not be present. A distinctive feature of this model is the fact that $a_0 = a_3$, independently of any parameter of the theory. Since we placed both e_L and e_R as $SU(2)_A$ doublets, we obtained $g_A = 0$ and the leptonic neutral current is purely vector. If we choose $\varphi=0$ and $\lambda \simeq 0$ one obtains the interesting result:

$$-g_V = -v_0 = v_3 = a_3 = a_0 = \frac{\varepsilon}{2}\left(1-\sin^2\theta\right) \qquad (2.13)$$

and all the neutral currents enter with equal strength. For $\varphi=\pi/2$ and $\lambda \simeq 0$ one gets a purely vector hadronic neutral current. However, it can be easily seen that in this model, it is not possible to couple the two vector bosons respectively to a purely vector and axial vector currents as it is the case in the left-right symmetric model described in 2.2.1.

2.3 Physical Constraints From Neutral Current Experiments

2.3.1 Purely leptonic interactions

The structure of the leptonic neutral currents can be best tested by studying neutrino-electron scattering. We have four reactions of interest:

$$\begin{matrix}\nu_e \\ \bar{\nu}_e\end{matrix} + e \rightarrow \begin{matrix}\nu_e \\ \bar{\nu}_e\end{matrix} + e \qquad (2.14)$$

$$\begin{matrix}\nu_\mu \\ \bar{\nu}_\mu\end{matrix} + e \rightarrow \begin{matrix}\nu_\mu \\ \bar{\nu}_\mu\end{matrix} + e \qquad (2.15)$$

Reactions (2.15) are due to neutral currents, while (2.14) receive contribution of both neutral and charged currents. A computation of these cross-sections gives:

$$\frac{d\sigma}{dy} = \frac{2G^2 m E_\nu}{\pi}\left[A + B(1-y)^2 - C\left(\frac{m_e\, y}{E_\nu}\right)\right] \qquad (2.16)$$

Table of neutral-current coupling constants

Table 2.1.

	g_V	g_A	v_3	a_3	$v_o V_o$	$a_o A_o$
Weinberg	$-\frac{1}{2}(1-4\sin^2\theta_W)$	$-\frac{1}{2}$	$1-2\sin^2\theta_W$	1	$v_0 V_0 = -2\sqrt{3}\sin^2\theta_W V_8 + \tilde{V}_0$	$a_0 A_0 = \tilde{A}_0$
Mohapatra	$-1+2\sin^2\theta_W$	0	$\frac{3}{2}-2\sin^2\theta_W$	$\frac{1}{2}$	$v_0 V_0 = (-2\sin^2\theta_W - \frac{1}{2})V_8 + \tilde{V}_0$	$a_0 A^0 = -\frac{\sqrt{3}}{2}A_8 + \tilde{A}_0$
Vector-like	$-1+2\sin^2\theta_W$	0	$2(1-2\sin^2\theta_W)$	0	$v_0 V_0 = -2\sqrt{3}\sin^2\theta_W V_8 + \tilde{V}_0$	$a_0 A_0 = \tilde{A}_0$
C.C.N.Y. ($g \simeq 0$, $\lambda \simeq 0$)	$\frac{1}{2}(-1+\sin^2\theta)$	0	$\frac{1}{2}(1-\sin^2\theta)$	$-\frac{1}{2}(1-\sin^2\theta)$	$v_0 V^0 = \frac{\sqrt{3}}{2}(-1+\sin^2\theta)V_8 + \tilde{V}_0$	$a_0 A_0 = \frac{\sqrt{3}}{2}(1-\sin^2\theta)V_8 + \tilde{A}_0$

where $y = \frac{E_e}{E_\nu}$, E_e and E_ν being the recoiling electron and incident neutrino energies, respectively. If we assume electron-muon universality, A, B, C, can be expressed in terms of only g_V, g_A defined in (2.2). We drop hereafter the last last term in (2.16), assuming $E_\nu \gg m_e$. The values of A, B for the four reactions are given in the table below, where we introduced $G_- = \frac{1}{2}(g_V - g_A)$ $G_+ = \frac{1}{2}(g_V + g_A)$. Integrating (2.16) one gets:

$$\sigma = 2\beta \left[A + \frac{1}{3} B \right] \tag{2.17}$$

with $\beta = \frac{G^2 m E_\nu}{\pi} = \frac{G^2 s}{2\pi}$. Since we have four reactions expressed in terms of only two coupling constants, constraint relations can be deduced. Namely, from table 2.2 and (2.17) one obtains:

$$\begin{gathered} \sigma_{\nu_e e} - 3\,\sigma_{\bar{\nu}_e e} = \sigma_{\nu_\mu e} - 3\,\sigma_{\bar{\nu}_\mu e} \\ \left| \left[\sigma_{\nu_e e} - \frac{1}{3}\,\sigma_{\bar{\nu}_e e} \right]^{\frac{1}{2}} \pm \left[\sigma_{\nu_\mu e} - \frac{1}{3}\,\sigma_{\bar{\nu}_\mu e} \right]^{\frac{1}{2}} \right| = \frac{4}{3}\beta \end{gathered} \tag{2.18}$$

where the $\pm$ sign means that one of the two relations should be satisfied. These relations, first derived by Sehgal,[10] are quite general and test effective locality, muon-electron universality and the V, A character of the interactions.

Table 2.2

Reactions \ Coefficients	A	B
$\nu_e\, e^-$	$(1+G_+)^2$	G_-^2
$\bar{\nu}_e\, e^-$	G_-^2	$(1+G_+)^2$
$\nu_\mu\, e^-$	G_+^2	G_-^2
$\bar{\nu}_\mu\, e^-$	G_-^2	G_+^2

We are interested here in finding constraints implied by the class of models under consideration. These can be found by looking at the domain in $G_+ G_-$ space covered by the different models. We will first consider models based on SU(2) x U(1). In vector-like models, $G_+ = G_-$, which implies, among other relations:

$$\sigma_{\bar{\nu}_\mu e^-} = \sigma_{\nu_\mu e^-}$$

$$\left(\sigma_{\nu_e e^-} - \sigma_{\bar{\nu}_e e^-} - \frac{4\beta}{3}\right)^2 = \frac{4\beta}{3}\left(3\,\sigma_{\bar{\nu}_\mu e^-} - \sigma_{\nu_\mu e^-}\right) \tag{2.19}$$

$$\left.\frac{d\sigma_{\nu_\mu e^-}}{dy}\right|_{y=1} = \left.\frac{d\sigma_{\bar{\nu}_\mu e^-}}{dy}\right|_{y=1} = \left.\frac{d\sigma_{\bar{\nu}_e e^-}}{dy}\right|_{y=1} = \frac{1}{2}\left.\frac{d\sigma_{\nu_\mu e^-}}{dy}\right|_{y=0} = \frac{1}{2}\left.\frac{d\sigma_{\bar{\nu}_\mu e^-}}{dy}\right|_{y=0} \qquad (2.20)$$

In non vector theories (i.e. $t^3_{e_R} \neq -\frac{1}{2}$), we note that

$$D_{-1} \supset \left(D_{-3/2} \cup D_{-2} \cup D_{-5/2} \cdots\right) \qquad (2.21)$$
$$D_0 \supset \left(D_{1/2} \cup D_1 \cup D_{3/2} \cup D_2 \cdots\right)$$

where D_j is the domain of the model with $t_3^{e_R} = j$ and arbitrary $\mathcal{E}$, θ. Using (2.21) one sees that either $G_+ \lesssim \left(-\frac{1}{2} + G_-\right)$ or $G_+ \gtrsim \left(\frac{1}{2} + G_-\right)$ which implies that one of the two following inequalities should be satisfied independently of θ, $\mathcal{E}$, $t^3_{e_R}$ (but $t^3_{e_R} \neq -\frac{1}{2}$):

$$E - M \lesssim \frac{7}{3}\beta + \frac{3}{2}\sqrt{3\beta}\ \bar{E} \quad \text{, or} \qquad (2.22)$$
$$E - M \gtrsim \frac{25}{3}\beta + \frac{3}{2}\sqrt{3\beta}\ \bar{E}$$

where $E = (3\sigma_{\nu_\mu e^-} - \sigma_{\bar{\nu}_\mu e^-})$ and $\bar{E} = (3\sigma_{\bar{\nu}_\mu e^-} - \sigma_{\nu_\mu e^-})$. and similarly for $M, \bar{M}$. In the minimal Weinberg model the following relation shold be satisfied:

$$(E - M)^2 = 8\bar{E}\left(E + M - 2\bar{E}\right) \qquad (2.23)$$

Note that (2.23) is independent of θ and it is thus a crucial test on the Weinberg-Salam model, especially when more accurate experimental data becomes available. Another point worth emphasizing is the fact that the class models based on SU(2) x U(1) and definite isotopic spin for e_R do not cover all the plane G_+G_- (obviously, this is why constraints can be derived), as illustrated in fig. (2.1).

Domain in G_+, G_-, for models based on $SU(2) \times U(1)$ and ~~with definite~~ e_R ~~spin~~

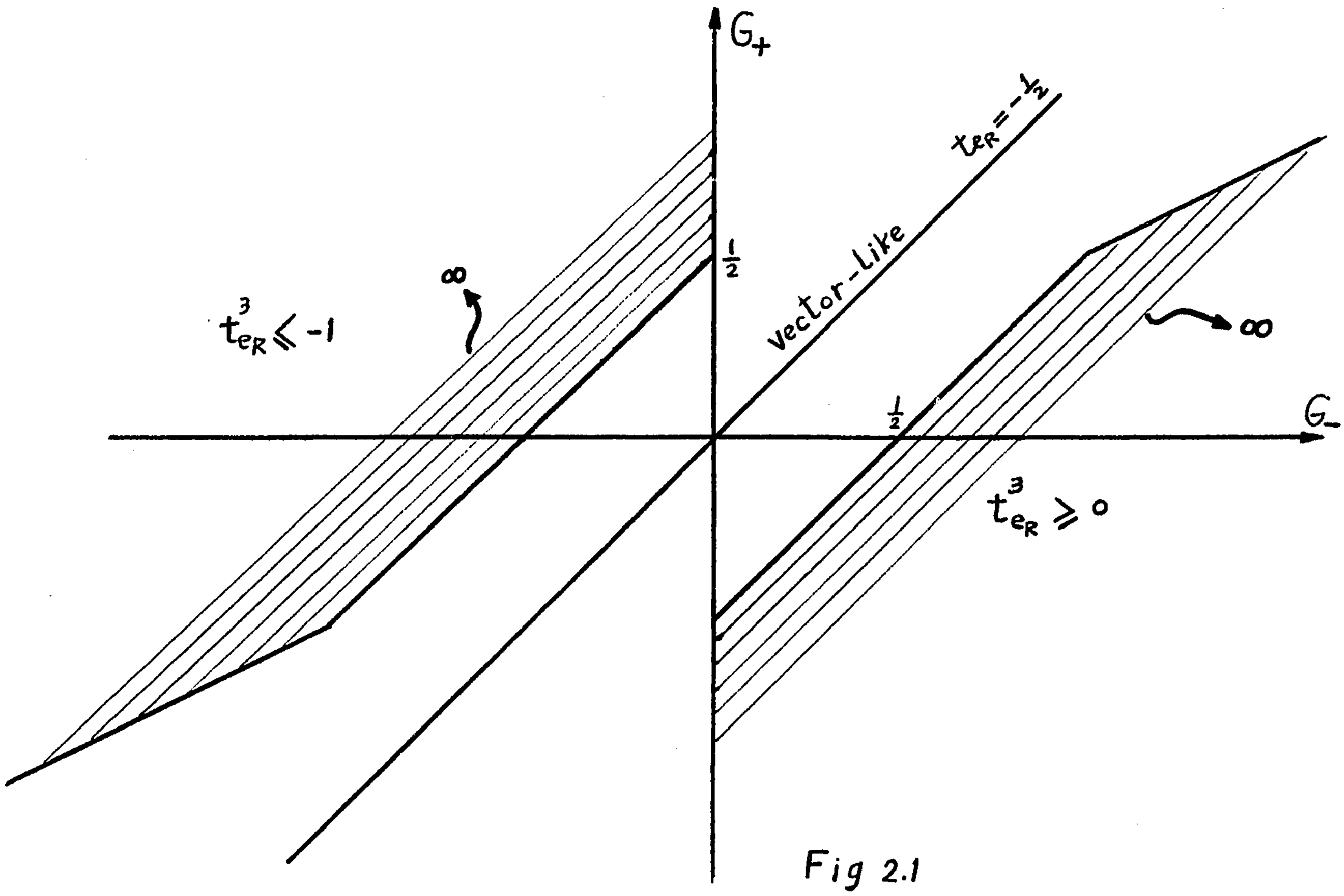


Fig 2.1

Domain in G_+G_- for left-right symmetric models

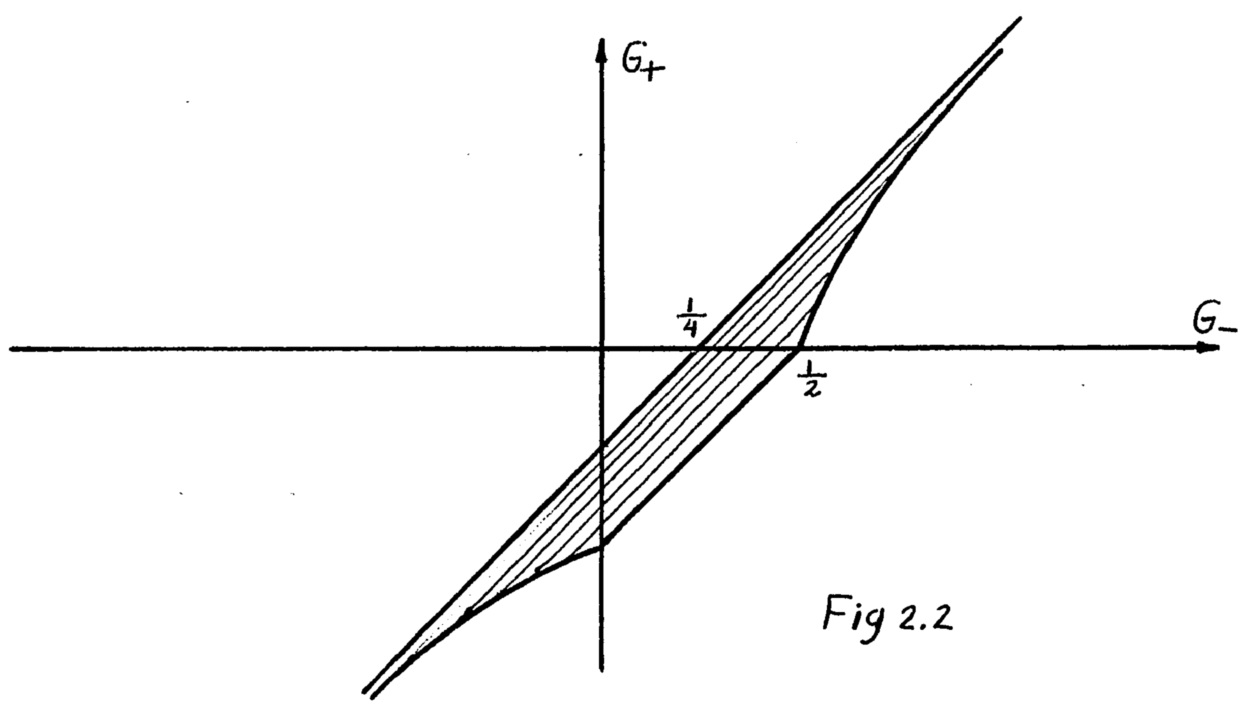


Fig 2.2

Wefinallyanalyse the left- right symmetric model considerd in 2.2.1.

Noting from (2.9), that:

$$G_+ = G_- - \frac{1}{2}\left(\frac{1+\alpha}{1+2\alpha}\right) \qquad (2.24)$$

one can easily obtain the allowed domain, shown in fig. 2.2. It is apparent that this model covers some of the domain excluded by the SU(2) x U(1) which may be of interest, if the previously derived constraints are not satisfied. However, the most interesting feature of the left-right symmetric model stems from the fact that it is parity conserving, with non-vanishing g_V, g_A. Parity violations can be detected by studying electron-positron annihilation into muon pairs:

$$e^+ + e^- \longrightarrow \mu^+ + \mu^- \qquad (2.25)$$

In general the effective neutral current interaction Lagrangian is:

$$\mathcal{L}_{eff} = \frac{G_F}{\sqrt{2}}\left[g_{VV}\, V_\mu V^\mu + 2 g_{VA}\, V_\mu A^\mu + g_{AA}\, A_\mu A^\mu\right] \qquad (2.26)$$

with

$$\begin{aligned} V_\mu &= \bar{e}\gamma_\mu e + \bar{\mu}\gamma_\mu \mu \\ A_\mu &= \bar{e}\gamma_\mu\gamma_5 e + \bar{\mu}\gamma_\mu\gamma_5 \mu \end{aligned} \qquad (2.27)$$

In models with only one neutral vector boson one necessarily has $g_{VA}^2 = g_{VV}\, g_{AA}$, but if there is more than one vector meson this relation does not hold in general. For example,in the L $\longleftrightarrow$ R symmetric model g_{VA} vanishes while g_{VV}, g_{AA} are non-vanishing. The parity violating effects are due to the term in g_{VA} and can be detected by evaluating the longitudinal muon polarization. A clear signal favoring the L-R symmetric model would be the confirmation of non-vanishing g_V, g_A from $\frac{\nu}{\bar{\nu}}$ e scattering, together

with non-observation of parity violation effects in $e^+e^- \to \mu^+\mu^-$.
This is a way in which the left-right symmetry proposed by Pati-Salam-Mohapatra could manifest itself even at present energies.

Finally, we analyze the implications of the presently available data for the various models considered. The data is summarized in fig. 2.3 where the ellipses denote the bounds on $C_\pm$ from Gargamelle neutrino experiments, and Reines reactor experiments. We note that among the SU(2) x U(1) models only the ones corresponding to t_{e_R} equal to $0, -\frac{1}{2}, -1$ are not excluded. These experimental results imply bounds on the parameters of the various models, and we obtain:

Weinberg Model: $\quad .1 < \varepsilon \sin^2\theta_W < .3$

Mohapatra model, vector-like Models: $\quad .11 < \varepsilon\left(-\frac{1}{2} + \sin^2\theta\right) < .23$

L-R symmetric Model
$$\begin{cases} .2 < \sin^2\theta < .6 & \text{for } \alpha = 0 \\ .26 < \sin^2\theta < .6 & \text{for } \alpha = 1 \end{cases} \tag{2.28}$$

C.C.N.Y. Model
(with $\varphi = 0$, $\lambda \simeq 0$) $\quad .22 < \varepsilon\left(1 - \sin^2\theta\right) < .46$

2.3.2 Semi-leptonic interactions

One of the most important questions concerning the hadronic neutral currents, is to know whether V^3_μ, A^3_μ, belong to the same isomultiplet as the charged hadronic currents. If one assumes that this is the case, one can deduce some interesting constraints on the various parameters of

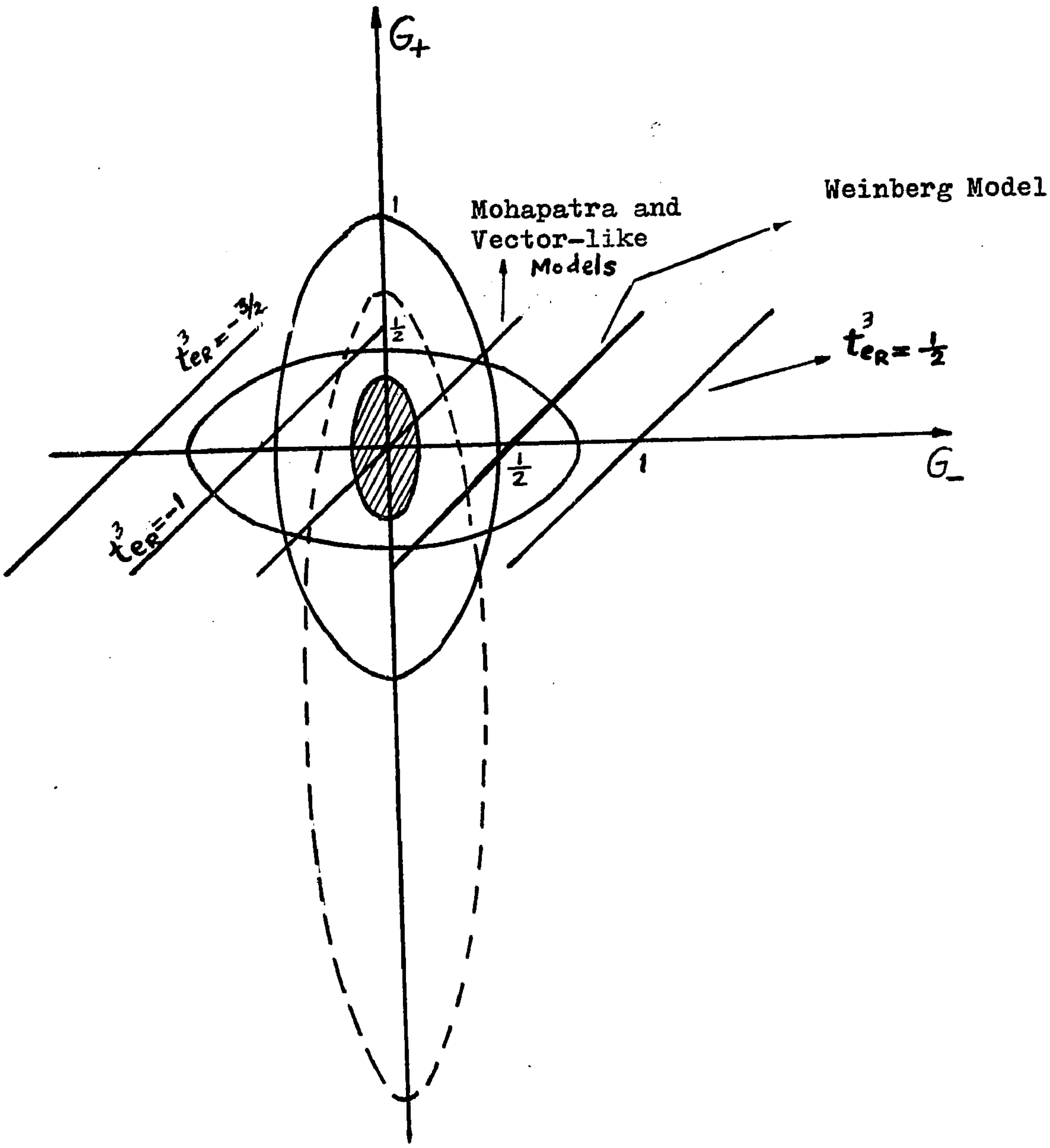


fig 2.3

The physical region determined by present experiments is the inside of the three ellipses, with the dashed ellipse excluded. The various solid lines are the loci of the various models, with $\varepsilon = 1$ and $\sin^2\theta$ varying from zero to one, from left to right.

the hadronic neutral current. Consider the cross sections:

$$\sigma^N \equiv \sigma(\nu p \to \nu + \cdots) + \sigma(\nu n \to \nu + \cdots)$$
$$\sigma^C \equiv \sigma(\nu p \to \mu^- + \cdots) + \sigma(\nu n \to \mu^- + \cdots) \tag{2.29}$$

with the corresponding antineutrino cross sections denoted by $\bar{\sigma}^N$, $\bar{\sigma}^C$. If V_μ^3, A_μ^3 belong to the same isomultiplet as the charged weak currents then σ^N and σ^C are related, and Paschos and Wolfenstein have used this fact to deduce the following inequality:

$$R = \frac{\sigma^N + \bar{\sigma}^N}{\sigma^C + \bar{\sigma}^C} \geq \frac{1}{2}\,\frac{v_3^2\,V + a_3^2\,A}{V + A} \tag{2.30}$$

where V, A are the contributions of V_μ^3, A_μ^3 to the cross section. Assuming that $V \approx A$, one obtains:

$$R \geq \frac{1}{4}\left(v_3^2 + a_3^2\right) \tag{2.31}$$

From the experimental value for $R = .27 \pm .09$ one constraints v_3, a_3 to lie within the circle represented in Fig. 2.4. Using Schwarz inequalities Rajasekaran and Sarma[14] were able to further restrict the values of v_3, a_3 to lie within the octogon of Fig. 2.4. We now analyze the implications of these restrictions for the various gauge models. In the case of the Weinberg model one obtains:

$$\sin^2\theta_W \simeq .36 \tag{2.32}$$

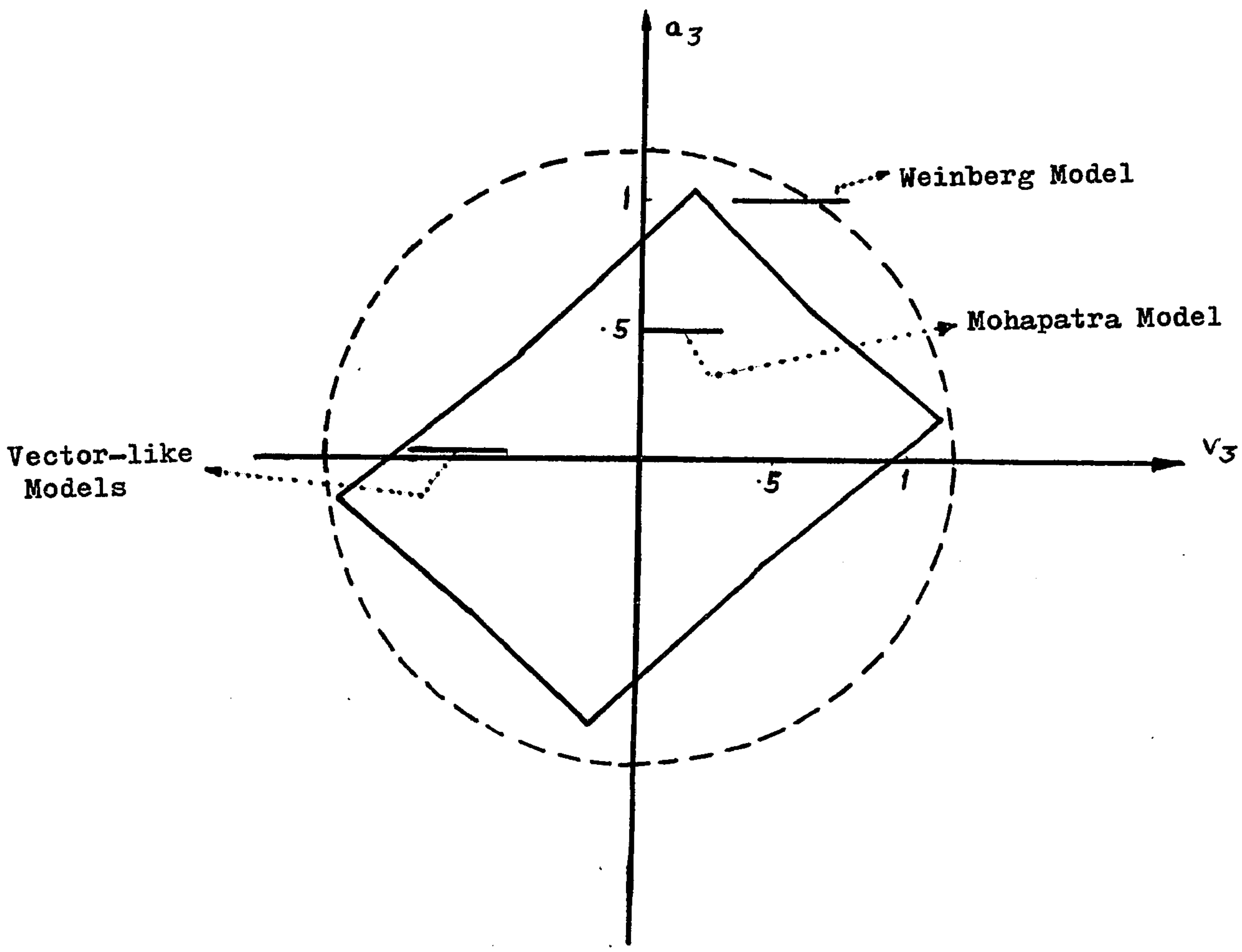


fig 2.4

Allowed ranges for the isovector coupling constants v_3, a_3.

If one takes into account experimental errors, (which were not included in Fig. 2.4) this is in agreement with the constraints (2.28) derived from the leptonic sector. Note that only the value $\varepsilon \simeq 1$ is consistent with experiment and therefore only the simplest Higgs system is allowed. Let us consider now the Mohapatra model. Since in this model $a_3 = \frac{1}{2}$. in order to be within the octogonon in Fig. 2.4 we need:

$$V_3 < .7$$

while using constraint (2.28), we obtain (putting $\varepsilon = 1$):

$$.05 \lesssim V_3 < .27$$

The Mohapatra model is thus in agreement with both leptonic and hadronic data when $.61 < \sin^2\theta_W < .73$. We choose $\varepsilon = 1$ for simplicity, but it is clear that other values of ε are also allowed. A similar analysis can be done for other models and we present in Fig. 2.4 the locus of V_3, a_3 for various theories by letting θ_W to vary within the range consistent with the leptonic sector. We will now point out the most important features of the hadronic neutral currents in models with R.H.C.C., as compared to the standard Weinberg-Salam model. For the standard model, one obtains using (2.32) and table 2.1 :

Weinberg-Salam:

$$V_3 \simeq .30 \;;\quad a_3 = 1 \;;\quad (V_3^2 + a_3^2) \simeq 1.1$$

$$V_0 \simeq -.23 \;;\quad a_0 = 0 \;;\quad V_0^2 \simeq .05$$

For models with R.H.C.C. we obtain on the other hand:

Mohapatra:

$$.05 < V_3 < .27 \quad ; \quad a_3 = \frac{1}{2} \quad ; \; .25 < V_3^2 + a_3^2 < .33$$

$$1.25 < -V_0 < 1.56 \quad ; \quad a_0 = -\frac{1}{2} \quad ; \qquad < V_0^2 + a_0^2 <$$

Harvard-Caltech-Princeton:

$$.5 < -V_3 < .9 \quad ; \quad a_3 = 0 \qquad .25 < V_3^2 + a_3^2 < .81$$

$$.75 < -V_0 < 1 \quad ; \quad a_0 = 0 \qquad .56 < V_0^2 + a_0^2 < 1$$

For the C.C.N.Y. models, the structure depends very much on the various parameters of the theory. For the special choice $\wp \simeq 0$, $\lambda \simeq 0$, $\varepsilon = 1$ we obtain:

$$.11 < -g_V = -V_0 = +V_3 = +a_3 = +a_0 < .23$$

It is clear from this analysis that the standard model has a dominant isovector component, and a very small isoscalar component. In the models with R.H.C.C. there is a general suppression of the isovector component and a significant contribution from the isoscalar component. A distinctive feature of the Mohapatra and C.C.N.Y. models is, of course, the existence of an axial isoscalar current. All this indicates the great importance of devising experiments that would allow a separate determination of a_3 a_0, V_3, V_0. The axial-vector parts have special interest since in the case of SU(2) x U(1) models, they do not depend on the Weinberg angle.

Various suggestions have been made for experiments which would enable the separate determination of the various couplings of the hadronic neutral current. For example, the reaction:

$$\bar{\nu}_e + d \rightarrow \bar{\nu}_e + n + p \qquad (2.33)$$

allows a separate determination of a_3 and is being studied by Gurr, Reines and Sobel. Unfortunately, at present, the experiment gives only an upperbound:

$$a_3^2 < 6$$

As far as the axial-isoscalar part of the neutral current is concerned, it can be detected by studying elastic scattering of neutrinos and anti-neutrinos off deuterium:

$$\begin{aligned} \nu + d &\rightarrow \nu + d \\ \bar{\nu} + d &\rightarrow \bar{\nu} + d \end{aligned} \qquad (2.34)$$

Only the isoscalar parts of the neutral currents (both V and A) contribute. If only V_0 is present, the two cross sections (2.34) are equal and thus a non-vanishing $\sigma^{\nu d} - \sigma^{\bar{\nu} d}$ is a clear signal for the presence of an axial isoscalar current. The determination of a_0 is of great importance for model building, because a non-vanishing a_0 would eliminate all models so far proposed with the exception of the Mohapatra and C.C.N.Y. models. Independently of gauge models considerations, the existence of an axial isoscalar current is also of special interest, dueto the fact that while V_0, V_3, A_3 can in principle be related through isotopic spin to the electromagnetic and charged weak currents, A_0 is necessarily a new current.

We will now consider the following ratios for the neutral and charged current interactions:

$$R^N = \frac{\bar{\sigma}^N}{\sigma^N} \quad ; \quad R^C = \frac{\bar{\sigma}^C}{\sigma^C}$$

In the framework of a naive quark parton-model one readily obtains for the neutral current ratio:

Weinberg model
$$R^N = \frac{\frac{1}{2} - \sin^2\theta_W + \frac{20}{9}\sin^4\theta_W}{\frac{3}{2} - 3\sin^2\theta_W + \frac{20}{9}\sin^4\theta_W}$$

Mohapatra model
$$R^N = \frac{5/4 - 2\sin^2\theta_W + 20/9\sin^4\theta_W}{7/4 - 10/3\sin^2\theta_W + 20/9\sin^4\theta_W}$$

Vector-like models
$$R^N = 1 .$$

Note that the vector-like models predict $R^N = 1$ independently of any parameter of the theory . In table 2.2 we show the present experimental data. One sees that the various experimental values are in contradiction with each other. Namely, the Gargamelle results predict

$$R^N = .675 \pm .25$$

While HPW-FNAL give:

$$R^N = 1 \pm .3$$

It is apparent that the H.P.W-FNAL results are consistent with the vector-like models, while Gargamelle rules out these models. Very recently, the H.P.W.-FNAL group has changed their initial data and their new results also rule out the vector-like models. Both Weinberg and Mohapatra models are consistent with the Gargamelle data, within experimental errors.

The predictions of the various models for R^C depend crucially on our assumptions for the masses of the new quarks. In table 2.1 we give the predictions of the naive quark model, in the various models with R.H.C.C., assuming that the energy is high enough for the production of all new flavors of quarks.

Table 2.1

Models	R^C
Mohapatra	$\frac{1}{4}$
Harvard, Princeton, Call-Tech	1
C.C.N.Y.	$\frac{1}{4}$ for $Q^2 \ll m_B^2$ 1 for $Q^2 \gg m_B^2$

Table 2.2

Ratio	Experiment	Energy range	Group
$\frac{\sigma(\nu N \to \nu + \cdots)}{\sigma(\nu N \to \mu^- + \cdots)}$	$.217 \pm .026$ $.11 \pm .05$	1-10 GeV 5-150 GeV	Gargamelle H.P.W.-FNAL
$\frac{\sigma(\bar{\nu} N \to \bar{\nu} + \cdots)}{\sigma(\bar{\nu} N \to \mu^+ + \cdots)}$	$.43 \pm .12$ $.32 \pm .09$	1-10 GeV 5-150 GeV	Gargamelle H.P.W.-FNAL

CHAPTER III

$K_L - K_S$ MASS DIFFERENCE

3.1 Introductory Remarks

A unified gauge theory with only the Cabbibo charged current, necessarily includes a neutral current containing both $\Delta S = 0$ and $\Delta S = 1$ parts. This can be easily seen by computing the commutator

$$\frac{1}{2}\left[J_\mu^+(x), J_0^-(y)\right] \delta(x^0 - y^0) = J_\mu^3(x)\, \delta^{(4)}(x-y) + \text{Schwinger terms}$$

If J_μ^+ is the usual Cabbibo current

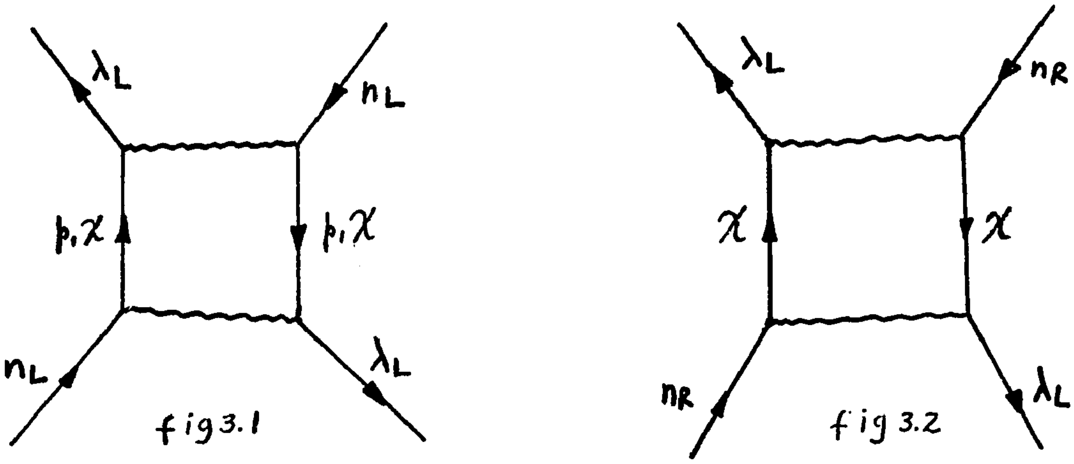
$$J_\mu^+ = \bar{p}\, \gamma_\mu (1+\gamma_5)\left[n \cos\theta + \lambda \sin\theta\right] ,$$

J_μ^3 will include strangeness-charging neutral currents which by gauge invariance should couple to the gauge bosons with the same strength as the charged currents. This in turn has disastrous consequences such as comparable rates for decays like $K_L \rightarrow \mu^+\mu^-$ and $K^+ \rightarrow \mu^+\nu$. A possible remedy consists of adding an extra piece to J_μ^+, containing new quarks, new heavy leptons, or both. Perhaps the most elegant scheme is the one suggested by Glashow, Iliopoulos and Maiani (GIM) who added an extra piece $\bar{\chi}$ $(-n \sin\theta_C + \lambda \cos\theta_C)$ to the charged current. Although this serves the purpose of eliminating first order $\Delta S = 1$ transitions, it turns out that one has also to analyze higher order transitions. This is due to the fact that second order processes are, in general, of order $G_F\,\alpha$ while the experimental value for $K_L - K_S$ mass difference or $K_L \rightarrow \mu^+\mu^-$ decay rate indicates a strength about three orders of magnitude smaller. One thus needs an extra suppression. In this respect, the smallness of the $K_L - K_S$ mass difference played a crucial role and enabled Gaillard and Lee to place a rather restrictive upper-bound in the mass of the charmed quark. The way this upper-bound on m_χ can be found is as follows: in the limit of exact of SU(4) symmetry the four diagrams in fig. 3.1 cancel each other and give zero net contribution to the K_L-K_S mass difference. In the real world where SU(4) is badly broken, the amplitude is proportional to $(m_\chi - m_p)^2/(m_W^2 \sin^2\theta_W)$ and thus leads to an upper-bound on m_χ.

In this chapter we analyze the K_L-K_S mass difference in the context of gauge models with the R.H.C.C. $\bar{\chi}_R \gamma_\mu n_R$. We comment also on the implications of our calculation for the standard model.

3.2 Evaluation of the effective $\Delta S=2$ weak Lagrangian

In this section, we evaluate the effective $\Delta S=2$ Lagrangian in the free quark model. The justification for the neglect of strong interaction effects resides in the belief that quarks are confined in a finite region of space and that within the "confinement" they behave as if they were free. In models with the current $\bar{\chi}_R \gamma_\mu n_R$, the following diagrams contribute to the elementary process $\bar{\lambda} + n \to \bar{n} + \lambda$:



The diagrams of fig.3.1 are already present in the standard model with only left-handed currents. Their contribution has been evaluated by Gaillard and Lee who obtained (assuming $m_W^2 \gg m_{p'}^2 \gg m_p^2$):

$$\left[\mathcal{L}_{eff}\right]_{LL}^{\Delta S=2} = G_{LL}^{\Delta S=2} \left\{ \bar{\lambda} \gamma_\mu \tfrac{1}{2}(1+\gamma_5) n \, \bar{\lambda} \gamma_\mu \tfrac{1}{2}(1+\gamma_5) n + h.c. \right\} \qquad (3.1)$$

where

$$G_{LL}^{\Delta S=2} = -\frac{G_F}{\sqrt{2}} \frac{\alpha}{4\pi} \left(\frac{m_\chi}{m_W \sin\theta_W} \right)^2 \sin^2\theta_c \cos^2\theta_c \qquad (3.2)$$

We will concentrate our attention here on the diagram of fig.3.2, which is the leading new contribution to $\mathcal{L}_{eff}^{\Delta S=2}$. One finds in the free quark model (see details in Appendix I):

$$\left[\mathcal{L}_{eff}\right]_{LR}^{\Delta S=2} = G_{LR}^{\Delta S=2} \left\{ 4 \, \bar{\lambda} \tfrac{1}{2}(1-\gamma_5) n \;\; \bar{\lambda} \tfrac{1}{2}(1-\gamma_5) n \; + \right. \qquad (3.3)$$
$$\left. + \bar{\lambda} \sigma_{\mu\nu} \tfrac{1}{2}(1-\gamma_5) n \, \bar{\lambda} \sigma^{\mu\nu} \tfrac{1}{2}(1-\gamma_5) n + h.c. \right\}$$

where

$$G_{LR}^{\Delta S=2} = -\frac{G_F}{\sqrt{2}} \frac{\alpha}{4\pi} \left(\frac{m_\chi}{m_W \sin\theta_W} \right)^2 2 \cos^2\theta_c \left(\ln \frac{m_W}{m_\chi} - 1 \right) \qquad (3.4)$$

3.3 Evaluation of $M^{K^\circ - \bar{K}^\circ}$

In the evaluation of $M^{K^\circ - \bar{K}^\circ}$, the $K^\circ - \bar{K}^\circ$ matrix element of $\left[\mathcal{L}_{eff}\right]_{LR}^{\Delta S=2}$, we assume that the contribution of the tensor term is negligible. Next we demonstrate that the contribution of the scalar and pseudo-scalar terms oppose each other in sign for all intermediate states with same charge conjugation parity. To show this, we evaluate $M^{K^\circ - \bar{K}^\circ}$ by inserting a complete set of intermediate states:

$$M^{K^0-\bar{K}^0} = G_{LR}^{\Delta S=2} \sum_m \left\{ \langle K^0|\bar{\lambda} n|m\rangle\langle m|\bar{\lambda} n|\bar{K}^0\rangle + \langle K^0|\bar{\lambda}\gamma_5 n|m\rangle\langle m|\bar{\lambda}\gamma_5 n|\bar{K}^0\rangle \right\} \tag{3.5}$$

Noting now that:

$$\begin{aligned} \bar{\lambda} n &= S_6 - i S_7 \\ \bar{\lambda}\gamma_5 n &= i\left(P_6 - i P_7\right) \end{aligned} \tag{3.6}$$

we write:

$$\begin{aligned} \langle K^0|\bar{\lambda} n|m\rangle &= \langle K^0|\, S_6 - i S_7 |m\rangle \\ &= \langle K^0|\text{б}^{-1}\text{б}\,(S_6 - iS_7)\,\text{б}^{-1}\text{б}|m\rangle \\ &= C_m \langle \bar{K}^0|\, S_6 + i S_7|m\rangle \\ &= C_m \langle m|\, S_6 - i S_7|\bar{K}^0\rangle^* \end{aligned} \tag{3.7}$$

and similarly

$$\langle K^0|\bar{\lambda}\gamma_5 n|m\rangle = i\, C_m \langle m| P_6 - i P_7|\bar{K}^0\rangle^* \tag{3.8}$$

It follows then from (3.5) → (3.8) that

$$M^{K^0-\bar{K}^0} = G_{LR}^{\Delta S=2} \sum_m \left\{ C_m \left|\langle m|S_6 - i S_7|\bar{K}^0\rangle\right|^2 - C_m \left|\langle m|P_6 - iP_7|\bar{K}^0\rangle\right|^2 \right\} \tag{3.9}$$

We are thus in the presence of an alternating series in which CP-odd and CP-even intermediate states contribute with opposite sign to $M^{K^0-\bar{K}^0}$. Next we evaluate the contribution of the vacuum (CP-even) and π^0 (CP-odd) intermediate states.

a) Vacuum contribution:

Noting that:

$$P_6 - i P_7 = \frac{-1}{m_n + m_\lambda} \partial_\mu \left(A_\mu^6 - i A_\mu^7 \right) \qquad (3.10)$$

we obtain

$$\langle 0 | P_6 - i P_7 | \overline{K^0}(q) \rangle = - i \sqrt{2} \frac{f_K m_K^2}{(m_n + m_\lambda)} \qquad (3.11)$$

Thus we find for the vacuum contribution:

$$M_{Vac.}^{K^0 - \overline{K^0}} = - G_{LR}^{\Delta S = 2} \frac{2 f_K^2 m_K^4}{(m_n + m_\lambda)^2} \qquad (3.12)$$

b) π^0 contribution:

The π^0 intermediate state will contribute to the scalar terms in (3.9).

We first observe that:

$$\bar{\lambda} n = \frac{1}{m_n - m_\lambda} \partial_\mu \left(V_\mu^7 + i V_\mu^6 \right) \qquad (3.13)$$

In general, one has for the matrix element of a vector current between pseudoscalar meson states:

$$\langle P^k(p) | V_\mu^i | P^j(k) \rangle = i f_{ijk} \left[f_+(q^2)(p+k)_\mu + f_-(q^2)(p-k)_\mu \right], \qquad (3.14)$$

where $q=(p-k)$. In the SU(3) limit one has $f_-(q^2)=0$. From (3.13), (3.14), one obtains

$$\langle \overline{K^0} | \bar{\lambda} n | \pi^0(k) \rangle = \frac{1}{\sqrt{2}} \frac{1}{m_n - m_\lambda} \left(m_K^2 - m_\pi^2 \right) f_+(q^2) \qquad (3.15)$$

In order to estimate $M_{\pi^0}^{K^0 - \overline{K^0}}$, one has to know the q^2 dependence of

$f_+(q^2)$. A parametrization[18] of the $K_{\ell 3}$ form factor in the accessible range of q^2, which is consistent with experiment, is given by:

$$f_+(q^2) = \left(1 + \frac{\lambda_0}{m_\pi^2} q^2\right) f_+(0) \tag{3.16}$$

where

$$f_+(0) \approx .97 \quad \text{and} \quad \lambda_0 \cong .02 \pm .003 \tag{3.17}$$

For the present purpose, we extrapolate this information for the whole range of q^2 by using a dipole fit:

$$f_+(q^2) = \frac{1}{\left(1 + \frac{q^2}{M_0^2}\right)^2} \tag{3.18}$$

Expanding the dipole fit for low q^2, one gets:

$$M_0^2 \cong \frac{2 m_\pi^2}{\lambda_0} \cong (1.8 \pm .3)\ \text{Gev}^2 \tag{3.19}$$

Using (3.15) and (3.18) and integrating over the π^0 momenta, we finally obtain:

$$M_{\pi^0}^{K^0-\bar{K}^0} \cong G_{LR}^{\Delta S=2} \frac{1}{(2\pi)^2} \frac{(m_K^2 - m_\pi^2)^2}{(m_n - m_\lambda)^2} \frac{M_0^8}{12 m_K^2 (M_0^2 - m_K^2)^2} \tag{3.20}$$

We now compare the magnitude of these two contributions:

$$\left| \frac{M_{vac.}^{K^0-\bar{K}^0}}{M_{\pi^0}^{K^0-\bar{K}^0}} \right| \cong 96\pi^2 \left(1 - \frac{m_\pi^2}{m_K^2}\right)\left(1 - \frac{m_K^2}{M_0^2}\right)^2 \left(\frac{m_K^2\, m_\pi^2}{M_0^4}\right)\left(\frac{m - m_\lambda}{m_n + m_\lambda}\right)^2 \tag{3.21}$$

If we choose $M_o^2 \approx 1.8\ (\mathrm{GeV})^2$ and $m_n \simeq \frac{m\lambda}{20}$, this ratio is .8 whereas for $M_o^2 \simeq 1.6\ (\mathrm{GeV})^2$ is 1. We have thus shown that the vacuum and π^o contribution to $M^{K^o-\overline{K^o}}$ are of the same order of magnitude and of opposite sign.

3.4 Concluding remarks

We have illustrated the inadequacy of the convention approximation of keeping only the vacuum intermediate state in the evaluation of $M^{K^o-\overline{K^o}}$. Although we have worked in the framework of models with the Mohapatra right-handed current $\bar{\chi}_R \gamma_\mu n_R$, it can be easily shown that our conclusion also holds for the standard model with only left-handed currents.

Our analysis has important consequences for models with R.H.C.C.: comparing (3.2) and (3.4) one notices the lack of Cabbibo suppression in $G_{LR}^{\Delta S=2}$. This has led Kingsley et al to conclude that the new charm-changing current is inconsistent with the observed K_L-K_S mass difference. However, their conclusion is based on the approximation of keeping only the vacuum contribution to $M^{K^o-\overline{K^o}}$. Since we have shown that this approximation is unjustified, we conclude that there is no conflict between the right-handed current $\bar{\chi}_R \gamma_\mu n_R$ and the value of the K_L-K_S mass difference.

Our analysis may also have some interesting implications for the standard Weinberg-Salam model, with only left-handed currents. The introduction of a fourth quark, as implied by the GIM mechanism, has interesting phenomenological implications due to the fact that the mass of the fourth quark is not arbitrary, and one can place an upper-bound on its value. Otherwise, one could assume the charmed quark to be massive enough to make its existence unimportant at present energies. As we have remarked in 3.1, constraints on the mass of the charmed quark come in principle from the study of the $\Delta S=2$ transitions, such as $K^+ \to \pi^+ \nu\bar{\nu}$, $K_L \to \mu^+\mu^-$ and the $K_L - K_S$ mass difference. In the case of the Weinberg-Salam model, Gaillard and Lee have shown that no constraints can be derived the decay $K_L \to \mu^+\mu^-$. Namely, these authors have shown that the two contributions to $K_L \to \mu^+\mu^-$, coming from $K_L \to W^+W^- \to \mu^+\mu^-$ and $K_L \to Z \to \mu^+\mu^-$, exactly cancel to order $G_F \propto \frac{m_\chi^2}{m_W^2}$, making thus the amplitude for $K_L \to \mu^+\mu^-$ of order $G_F \alpha^2$, independently of the mass of the charmed quark, m_χ. Furthermore, the present upperbound on $\Gamma(K^+ \to \pi^+ \nu\nu)/\Gamma(K^+ \to \text{all})$ does not lead to a very restrictive bound on m_χ. It is then only the value of the $K_L - K_S$ mass difference that leads to an upperbound on m_χ, of order 2 or 3 GeV. However, in their analysis Gaillard and Lee used the approximation of keeping only the vacuum contribution of $\mathcal{L}_{\text{eff.}}^{\Delta S=2}$. Since we have shown that this approximation is unjustified and perhaps misleading, the possibility of a charmed quark considerably heavier than what is usually assumed, remains open. The eventual connection between this possibility, and the present non-observation of particles carrying net charm is, of course, a matter of pure speculation.

CHAPTER IV

DECAYS OF CHARMED HADRONS

4.1 Introductory Remarks

Although many interpretations have been advanced for the newly discovered resonances, one of the most attractive suggestions is that they confirm the need for a new quantum number (charm) in the description of the strong interactions. In this scheme, the new particles are viewed as bound states of a charmed quark and its antiquark. At present, it is still an open question if this is indeed the case. Recently, this "charm interpretation" has been challenged by various authors,[20] who suggested instead that the new particles are bound states of unconfined integer charge heavy leptons. These heavy leptons would be the same objects which have been detected in the SLAC experiments. Even if this last view turns out to be correct, and the new particles do not have anything to do with "charm", the fact remains that "charm" seems to be a necessity,[16] if one wants to build realistic unified gauge models of weak and electromagnetic interactions. Hence, the great interest in finding a direct confirmation of the charm scheme, through the discovery of particles carrying net charm. A detailed study of[21,22,23] the properties of these particles has been done, based on the standard model of weak and electromagnetic interactions in which the four-quark model is incorporated into the Weinberg-Salam gauge theory. As far as the weak interactions are concerned, of particular interest is the question of how they break the new symmetry of the strong interactions. The study of the decay characteristics of the charmed hadrons seems to provide the crucial test for the different theoretical possibilities. Since the new quantum number is supposed to be broken only by the weak interactions, the least massive of these charmed hadrons are expected to decay only weakly. Most of the previous analyses of these decays have been done assuming a particular structure for the weak interactions of these particles, namely, the one indicated by the GIM charmed current. However, as it has been emphasized recently,[24] the possibility of an SU (4) symmetry for the strong interactions, (with its implied mass spectrum) and the problem of its symmetry breakdown by the weak interactions, are two independent questions.

As we have seen in previous chapters, one alternative to the GIM current consists of adding a right-handed charm-changing current. This current will, in general, be of the following type:

$$K_\mu = \alpha \, \bar{x}_R \gamma_\mu n_R + \beta \, \bar{y}_R \gamma_\mu \lambda_R + \gamma \, \bar{p}_R \gamma_\mu z_R + \cdots \tag{4.1}$$

where $p_R, n_R \lambda_R$ are the right handed p, n, λ quarks, and x_R, y_R, z_R stand for right-handed quarks, each one of them carrying a new quantum number conserved by the strong interactions. The dots stand for other currents, not containing any of the quarks p, n, λ. Either x or y will coincide with the charmed quark χ, used to implement the GIM mechanism.

In this chapter, we examine the weak decays of the pseudo-scalar bound states of each one of the x, y, z quarks with the three SU (3) quarks p, n, λ. We call these bound states $D_x^o(x\bar{p})$, $D_x^+(x\bar{n})$, $F_x^+(x\bar{\lambda})$, and similarly for y, z. It has in fact been recently suggested that the x, y quarks may indeed be degenerate. It is clear that either D_x^o, D_x^+, F_x^+ or D_y^o, D_y^+, F_y^+ will coincide with the D's and F's of ref. 21 As far as the decays of the charmed pseudo-scalar mesons are concerned, one has two classes of models, corresponding to the choices a) $x \equiv \chi$ and b) $y \equiv \chi$. We will deal separately with the two cases, pointing out in detail their implications to the non-leptonic decays of the charmed mesons $D_\alpha^o, D_\alpha^+, F_\alpha^+$, (where α = x or y) and D_z^o, D_z^-, F_z^o, and contrasting them with the predictions of the GIM model.

4.2 Leptonic Decays

4.2.1 Two body decays

We will analyse here the two body decays of F^+ $(x\bar{\lambda})$, D^+ $(x\,\bar{n})$, D^0 (xF). For completeness, we will review the predictions of the GIM model and then deal separately with classes a) and b).

4.2.1.1 GIM model

These decays are similar to $K_{\ell 2}$ decays and one thus obtains

$$\Gamma(D^{\pm}_{\ell_2}) = \frac{G^2}{8\pi} f_D^2\, m_D\, m_\ell^2 \left(1 - \frac{m_\ell^2}{m_D^2}\right) \sin^2\theta_c$$
$$\Gamma(F^{\pm}_{\ell_2}) = \frac{G^2}{8\pi} f_F^2\, m_F\, m_\ell^2 \left(1 - \frac{m_\ell^2}{m_F^2}\right) \cos^2\theta_c \qquad (4.2)$$

where f_D, f_F are defined by:

$$(2\pi)^{3/2}\, 2q^0 \langle 0 | J_\mu^C | D^+(q)\rangle = i f_D\, q_\mu \sin\theta_c$$
$$(2\pi)^{3/2}\, 2q^0 \langle 0 | J_\mu^C | F^+(q)\rangle = -i f_F\, q_\mu \cos\theta_c \qquad (4.3)$$

assuming[25]

$$f_D \simeq f_F \simeq f_K \simeq f_\pi \qquad (4.4)$$

one obtains

$$\Gamma(D^+ \to \mu^+\nu) \simeq \frac{m_D}{m_K}\,\Gamma(K^+ \to \mu^+\nu) \simeq 2\times 10^8\ \mathrm{sec}^{-1}$$
$$\Gamma(F^+ \to \mu^+\nu) \simeq \frac{m_F}{m_K}\,\mathrm{cotg}^2\theta_c\,\Gamma(K^+ \to \mu^+\nu) \simeq 5\times 10^9\ \mathrm{sec}^{-1} \qquad (4.5)$$

4.2.1.2 Class a) models

Due to the absence of Cabbibo suppression in the right handed current $(\bar{\chi}\, n)_R$, one obtains now:

$$\Gamma(F^+ \to \mu^+\nu) \simeq \Gamma(D^+ \to \mu^+\nu) \simeq 5\times 10^9\ \mathrm{sec}^{-1} \qquad (4.6)$$

4.2.1.3 Class b) Models

Note that in this case one has

$$J_\mu^C = \overline{\chi}_L \gamma_\mu \left(-n_L \sin\theta + \lambda_L \cos\theta\right) + \overline{\chi}_R \gamma_\mu \lambda_R$$ (4.7)

$$= -\sin\theta\, \overline{\chi}_L \gamma_\mu n_L + (1+\cos\theta)\, \overline{\chi} \gamma_\mu \lambda + (1-\cos\theta)\, \overline{\chi} \gamma_\mu \gamma_5 \lambda .$$

Since only the axial-vector current contributes to the two-body leptonic decays of pseudo-scalar particles, it follows from (4.7) that the $F^+_{\ell 2}$ decays are strongly suppressed. One thus obtains:

$$\Gamma(D^+ \to \mu^+ \nu) \simeq 2 \times 10^8\ \mathrm{sec}^{-1}$$ (4.8)

$$\Gamma(F^+ \to \mu^+ \nu) \simeq 2 \times 10^6\ \mathrm{sec}^{-1}$$

It is perhaps worthwhile emphasizing that the knowledge of the decay rates of the F and D charmed mesons allows in principle for a clear-cut distinction between the three classes of models, as it is illustrated in the table below:

Models	GIM	Class a	Class b
$\dfrac{\Gamma(F^+ \to \mu^+ \nu)}{\Gamma(D^+ \to \mu^+ \nu)}$	25	1	10^{-2}

4.2.2 Three body decays

4.2.2.1 GIM model

For the standard model, there is the selection rule $\Delta C = \Delta S$ and the enhanced decays are:

$$F^+ \to \eta + \ell + \nu$$

$$F^+ \to \eta' + \ell^+ + \nu$$ (4.9)

$$D^+ \to \bar{K}^0 + \ell^+ + \nu$$

$$D^0 \to K^- + \ell^+ + \nu$$

where ℓ^+ stands for either e^+ or μ^+. The magnitude of the decay rates has been estimated in ref. 21 . Assuming the usual f_+ form factors to be constant and equal for $D_{\ell 3}$ and $K_{\ell 3}$ decays, the order of magnitude of the decay rates is predicted to be $\simeq 10^{11}\ \mathrm{sec}^{-1}$.

4.2.2.2. Class a) models.

It is obvious that in this class of models, besides (4.9),
the following decays are enhanced:

$$
\begin{aligned}
&F^+ \to K^\circ + \ell^+ + \nu \\
&D^+ \to \pi^\circ + \ell^+ + \nu \\
&D^+ \to \eta + \ell^+ + \nu \\
&D^\circ \to \pi^- + \ell^+ + \nu
\end{aligned}
\tag{4.10}
$$

4.2.2.3. Class b) models.

The predictions for this class coincide with those of the GIM model since it is the vector current which contributes to these decays.

4.2.3 Multipion decays

Due to the large mass of the charmed mesons, one in principle expects multipion decays of the type $D^\circ \to \bar{K}^\circ + n\pi + \ell^+ + \nu$ to constitute a substantial part of the total semi-leptonic rate. However, this simple conclusion is doubtful for the GIM model, as well as for class b). This is due to the fact that for these models, the amplitude vanishes in the soft pion limit of any of the final state pions:

$$
\begin{aligned}
&\lim_{q \to 0} \langle \bar{K} \ldots \pi^i(q) \,|\, \bar{\chi}_L \gamma_\mu \chi_L \,|\, D^\circ \rangle = \\
&= \frac{i}{f_\pi} \langle \bar{K} \ldots | \left[F_5^i , \bar{\chi} \gamma_\mu \chi_L \right] | D^\circ \rangle = 0
\end{aligned}
\tag{4.11}
$$

where F_5^i is the axial charge. This argument is no longer true in class a) models and one thus expects a significant fraction of multipion decays.

4.3 Non-leptonic Decays

4.3.1. Properties of the non-leptonic Hamiltonian

We will assume that the non-leptonic weak Hamiltonian has the usual current-current form:

$$\mathcal{H}_W = \frac{G_F}{\sqrt{2}} \left(H_\mu H^{+\mu} + H_\mu^+ H_\mu \right) \tag{4.12}$$

The hadronic current H_μ contains both the usual left-handed J_μ and the new right-handed K_μ:

$$\begin{aligned} H_\mu &= J_\mu + K_\mu \\ J_\mu &= \bar{p}_L \gamma_\mu n_L(\theta_c) + \bar{\chi}_L \gamma_\mu \lambda_L(\theta_c) \end{aligned} \tag{4.13}$$

where K_μ is given by (4.1) and

$$\begin{pmatrix} n(\theta_c) \\ \lambda(\theta_c) \end{pmatrix} = \begin{pmatrix} \cos\theta_c & \sin\theta_c \\ -\sin\theta_c & \cos\theta_c \end{pmatrix} \begin{pmatrix} n \\ \lambda \end{pmatrix} \tag{4.14}$$

with θ_c the Cabbibo angle. Recall that χ denotes the charm quark used to implement the GIM mechanism.

As we show in the Appendix, this Hamiltonian decomposes with respect to SU(4) as:

$$\mathcal{H}_W = \underline{1} \oplus \underline{15}_S + \underline{20} + \underline{84} \tag{4.15}$$

Note that, with the modified current, the $\underline{15}$ is now present in H_W, contrary to what happens in the standard GIM model.

We will analyze now the part of H_W contributing to the charm changing non-leptonic decays. We will designate collectively by "charm" the new quantum numbers of the x, y, z quarks and distinguish among them by calling C_x, C_y, C_z. We will analyze separately the two classes of models mentioned before.

a) Class of models with $x \equiv \chi$

We recall that in this class of models, the $\Delta C_\chi = 0$, $\Delta S = \pm 1$ Hamiltonian is given effectively by

$$\mathcal{H}_W^S = \frac{G_F}{\sqrt{2}} \left\{ \cos\theta_c \, (\bar{\lambda}\chi)_L (\bar{\chi} n)_R + \cos\theta_c \sin\theta_c \left[(\bar{\lambda} p)_L (\bar{p} n)_L - (\bar{\lambda}\chi)_L (\bar{\chi} n)_L \right] \right\} \tag{4.16}$$

Where for convenience we have omitted Lorentz indices, and L, R stand for left and right-handed currents, respectively. The first term[5] contributes to $\Delta I = 1/2$ transitions only and is enhanced with respect to the second term (which contributes both to $\Delta I = 1/2$ and $\Delta I = 3/2$ transitions) by the Cabbibo factor $(\sin\theta)^{-1}$. The first term is further enhanced by the large mass of the charmed quark χ.

For the study of the decays of the SU(3) triplet of charmed pseudoscalar mesons $D_\chi^+, D_\chi^o, F_\chi^+$, we are interested in the $\Delta C_\chi = 1$ part of the Hamiltonian which is:

$$\mathcal{H}_W^{C\chi} = \mathcal{H}_{W(LL)}^{C\chi} + \mathcal{H}_{W(RL)}^{C\chi}$$

$$\begin{aligned} \mathcal{H}_{W(LL)}^{C\chi} = {} & \cos^2\theta_c \, (\bar{\chi}\lambda)_L (\bar{n} p)_L - \sin^2\theta_c \, (\bar{\chi} n)_L (\bar{\lambda} p)_L + \\ & + \cos\theta_c \sin\theta_c \left[(\bar{\chi}\lambda)_L (\bar{\lambda} p)_L - (\bar{\chi} n)_L (\bar{n} p)_L \right] \end{aligned} \tag{4.17}$$

$$\mathcal{H}_{W(RL)}^{C\chi} = \cos\theta_c \, (\bar{\chi} n)_R (\bar{n} p)_L + \sin\theta_c \, (\bar{\chi} n)_R (\bar{\lambda} p)_L$$

$H_{W(LL)}^{C\chi}$ is present in the standard gauge models with the GIM current. Its properties were fully discussed in ref.21,22,23. We will direct our attention to the new term $H_{W(RL)}^{C\chi}$ which appears in the right-handed current models of class a).

With respect to SU(3), H_W^C transforms as $\underline{8} \times \underline{3} = 3^* \times \underline{6} \times \underline{15}_M^*$ The states of these representations can be expressed in the basis of quark current-current products, as follows:

$$
\begin{aligned}
[\underline{3}^*]_i &= H_{ik}^{k} \\
[\underline{6}]^{kl} &= \varepsilon^{kij} H_{ij}^{l} + \varepsilon^{lij} H_{ij}^{k} - \frac{1}{4}\left(\delta_i^k H_{jl}^{l} + \delta_j^k H_{il}^{l}\right) \\
[\underline{15}_M]_{ij}^{k} &= H_{ij}^{k} + H_{ji}^{k} - \frac{1}{4}\left(\delta_i^k H_{jl}^{l} + \delta_j^k H_{il}^{l}\right)
\end{aligned}
\tag{4.18}
$$

where

$$
H_{ij}^{k} = (\bar{q}^k q_j)(\bar{\chi} q_i) - \frac{1}{3}\delta_i^k (\bar{q}^l q_l)(\bar{\chi} q_j) \tag{4.19}
$$

with

$$
(\chi, q_1, q_2, q_3) = (\chi, p, n, \lambda)
$$

Thus we find that $H_W^{C\chi}$ can be decomposed as,

$$
\mathcal{H}_{W(LL)}^{C_x} = \frac{1}{4}\left\{\cos^2\theta_c [\underline{6}]^{22} + \cos\theta_c \sin\theta_c [\underline{6}]^{23} + \sin^2\theta_c [\underline{6}]^{33}\right\}_{LL} \tag{4.20}
$$

+ terms belonging to the $\underline{15}_M^*$ representation

$$\mathcal{H}^{C\chi}_{W}(RL) = \frac{1}{4}\left\{\cos\theta_c\,[\underline{6}]^{23} + \sin\theta_c\,[\underline{6}]^{33} + \frac{3}{2}\cos\theta_c\,[3^*]_1\right\}_{RL} + \qquad (4.21)$$

+ terms belonging to the $\underline{15}^*_M$ representation.

Terms in the SU(3) $\underline{15}^*_M$ representation will be discarded hereafter, since they come from the SU(4) $\underline{84}$. We recall that the $\underline{84}$ contains the SU(3) $\underline{27}$, which is known to be suppressed in strangeness - changing decays. Note that the $\underline{3}^*$ is absent in $H^C_{W(LL)}$, but contributes to $H^C_{W(RL)}$.

We analyze now the part of the Hamiltonian contributing to the decay of the y-type and z-type "charmed" mesons. We obtain:

$$\mathcal{H}^{Cy}_{W(RL)} = (\bar{y}\,\lambda)_R(\bar{n}\,p)_L\cos\theta_c + (\bar{y}\,\lambda)_R(\bar{\lambda}\,p)_L\sin\theta_c$$

$$\mathcal{H}^{Cz}_{W(RL)} = (\bar{z}\,p)_R(\bar{p}\,n)_L\cos\theta_c + (\bar{z}\,p)_R(\bar{p}\,\lambda)_L\sin\theta_c \qquad (4.22)$$

These expressions can in turn be decomposed as:

$$\mathcal{H}^{Cy}_{W(RL)} = \frac{1}{4}\left\{\cos\theta_c\,[\underline{6}]^{22} + \sin\theta_c\,[\underline{6}]^{23} + \frac{3}{2}[3^*]_1\right\}_{RL} + \cdots$$

$$\mathcal{H}^{Cz}_{W(RL)} = \frac{1}{4}\left\{\cos\theta_c\,[\underline{6}]^{13} + \frac{3}{2}\cos\theta_c\,[\underline{3}^*]_2 - \sin\theta_c\,[\underline{6}]^{12} + \frac{3}{2}\sin\theta_c\,[3^*]_3\right\}_{RL} + \cdots \qquad (4.23)$$

b) Class of models with y≡χ

It is apparent that models of this type conform to the constraints derived in ref. 29 . The $\Delta S=\pm1$ Hamiltonian is now given by

$$\mathcal{H}_W^S = \frac{G}{\sqrt{2}} \cos\theta_c \sin\theta_c \Big\{ (\bar\lambda \chi)_R (\bar\chi n)_L + (\bar\lambda p)_L (\bar p n)_L - (\bar\lambda \chi)_L (\bar\chi n)_L \Big\} + h.c. \tag{4.24}$$

We see that the (RL) term, which contributes only to the $\Delta I=1/2$ transitions. is also suppressed by sin θ_c. Thus one has to rely only on a dynamical enhancement (which, for example, may arise either due to renormalization group arguments or otherwise) to explain the $\Delta I=1/2$ rule. Note that this right-handed charmed current may appear in a variety of gauge models, as for example in those of ref. 6 (in some of these models one would have to do the interchange $n_R \leftrightarrow \lambda_R$). We are interested here in the charm-changing part of the Hamiltonian. As explained in eq.4.17, there is a LL and a RL part in the Hamiltonian, the (LL) term remains unchanged (See eq.4.17) and the (RL) term is given by:

$$\mathcal{H}_{W(RL)}^{C\chi} = \frac{1}{4} \Big\{ \cos\theta_c \, [\underline{6}]^{22} + \sin\theta_c \, [\underline{6}]^{23} + \frac{3}{2} \sin\theta_c \, [3^*]_1 \Big\}_{RL} \tag{4.25}$$

Obviously, this coincides with (4.23) when y is identified with χ. The $\Delta C_\chi=1$ part of the Hamiltonian is now given by:

$$\mathcal{H}_{W(RL)}^{C\chi} = (\bar\chi n)_R (\bar n p)_L \cos\theta_c + (\bar\chi n)_R (\bar\lambda p)_L \sin\theta_c \tag{4.26}$$

which can in turn be decomposed as:

$$\mathcal{H}^{Cx}_{W(RL)} = \frac{1}{4}\left\{\cos\theta_c\,[\underline{6}]^{23} + \sin\theta_c\,[\underline{6}]^{33} + \frac{3}{2}\cos\theta_c\,[3^{*}]_1\right\}_{RL} \qquad (4.27)$$

Note that $\mathcal{H}^{Cz}_{W(RL)}$ is the same in both classes of models.

4.3.2 Analysis of non-leptonic decays.

4.3.2.1 Class of models with $x=\chi$

We will first analyze the various decay modes of $D^+_\chi, D^\circ_\chi, F^+_\chi$. From (9), it is apparent that:

(i) Right-handed currents contribute to $\Delta C=1$ transitions via both the $\underline{6}_{RL}$ and the $\underline{3}^*_{RL}$.

(ii) Decay modes which were suppressed by a $\sin\theta_c$ factor in the standard GIM model, can now occur through $[\underline{6}]^{23}_{RL}$ without Cabbibo suppression.

We present the tables with the relative rates (apart from phase space factors) for the decays of D^+, D°, F^+ into two pseudo-scalars, two vectors, a pseudo-scalar-vector pair and a baryon-antibaryon pair. Many of these results can be easily obtained from those of ref.23 by taking into account (i), (ii). However, for completeness we present all the tables. We will now point out some of their salient features.

1) Decays of F^+_χ, D^+_χ.

It is interesting to note that for the decays of F^+, D^+ into two pseudo-scalars, the relative rates can be expressed in terms of only two parameters P, (Q+R) which will be defined in Table 4.6 This is due to the fact that the contributions of $\underline{6}_{RL}$ and $\underline{3}_{RL}$ enter always in the same combination in F^+, D^+ decays. We will see that this is no longer true for D°. It is easy to see that the decay characteristics of D^+ are drastically changed. Recall that one of the salient features of the GIM current is that it predicts (if one assumes sextet dominance) no enhanced decays of D^+ into two pseudo-scalars. Based on this fact, it has been argued that for D^+, the semi-leptonic modes may compete favorably with the non-leptonic ones. It is apparent from Table (4.3) that this is no longer true for the new current, since there are decays of D^+ via the $\underline{6}_{RL}$ and $\underline{3}^*_{RL}$ with no Cabbibo suppression.

2) Decays of D°_{x}

For the decay of D° into two pseudo-scalars, one can measure separately the contributions of $\underline{6}_{RL}$ and $\underline{3}^*_{RL}$. Note for example that $D^{\circ} \to K^{\circ}\bar{K}^{\circ}$ goes only via $\underline{3}^*_{RL}$ and observation of this mode would thus provide a direct measure of its strength.

Of particular interest for the current search for charm, are the decays of D° into two charged particles. While in the GIM model, only π^+K^- is enhanced, now the $\pi^+\pi^-$, K^+K^- decays can go unsuppressed.

We will now analyze the decays of "y" and "z" type charmed mesons. From expression (4.23) for $H^{Cy}_{W(RL)}$, it is clear that the Cabibbo enhancement decays of $F^+_y, D^+_y, D^{\circ}_y$ will coincide with those of the F^+, D^+, D° in the GIM scheme. Since the SU(3) structure of the Cabibbo enhanced part for both $H_{W(LL)}$ and $H^{Cy}_{W(RL)}$ has the same SU(3) structure. We therefore do not present a separate table for these.

The decay of the z-type charmed mesons are given in Table (4.5).

4.3.2.2 Class of models with $y = \chi$

1) Decays of $F^+_\chi, D^+_\chi, D^\circ_\chi$

From 4.25 one sees that the Cabbibo enhanced part of $H^{C\chi}_{W(RL)}$ has the same form as $H^{C\chi}_{W(LL)}$ (recall that the LL part of H_W is the same in all models and coincides with that of the GIM scheme). The term in sin θ_C is however modified by the presence of $[3^*]_{RL}$. It follows then that as far as the enhanced non-leptonic decays of χ- charmed mesons are concerned the predictions of this model coincide with those of the standard model, except if the reduced matrix elements of the $[3^*]_{RL}$ are considerably larger than those of $[6]_{RL \cdot LL}$ (i.e. larger enough to compensate the Cabbibo suppression factor).

2) Decays of $F^+_\chi, D^+_\chi, D^\circ_\chi$

From (14) it is obvious that the decays of F^+_χ D^+_χ D°_χ can be obtained from those of Tables 4.2,3,4 by simply setting P=D=S=A=α=0.

4.3.3 Implication for Spear Charm Search.

We now analyze the impact of the prediction of right-handed current models on the Spear charm search. Boyarski et al [27] have looked for narrow peaks in the invariant mass distribution of various combination of particles, such as $K^\mp, \pi^\pm, \pi^+\pi^-, K^+K^-, K^\mp\pi^\pm\pi^\pm$ etc. No convincing narrow peaks were found in the invariant mass distributions and therefore it was possible to place upper limits for inclusive production cross section times branching ratio, for various decay modes and mass regions. In order to interpret these results, one has to make some assumptions about the magnitude of the charmed particle production cross sections, at $\sqrt{S}$ = 4.8 GeV. In making this estimation, one may take two different points of view:

1) The excess of hadronic cross section (σ hadronic) above

the value observed at $\sqrt{S} \sim 3.7$ GeV ($R \sim 2.5$) is all due to charmed particle production.

2) The observed excess of the σ hadronic is also due to the production of a heavy lepton, with a cross-section of about 2.8 nb.

Here, we will assume 1) but the analysis can be easily repeated for 2). Furthermore, we will assume that the charmed final states necessarily include one of the pairs D^+D^-, $D^o\overline{D^o}$, F^+F^- and that all these are produced with equal probability. From the experimental value for R ($R \simeq 4.83$ at $\sqrt{S}=4.8$ GeV) one deduces that each of these channels has a cross-section of about 3.14 nb., at $\sqrt{S}=4.8$GeV. Using table(4.1) one easily obtains the upper-bounds for the various branching ratios. We assume that the relevant mass region in table(4.1) is 1.85-2.40 GeV. The most serious challenge to the GIM current model seems to come from the following upper limit:

$$B(D^o \to K^-\pi^+) \leq 2.9\% \qquad (4.28)$$

In the standard GIM model (with the assumption of sextet dominance), the $K^-\pi^+$ rate is 1/3 the rate for $D^o \to$ PP. Hence one gets rather low bound

$$B(D^o \to pp) \leq 8.7\% \qquad (4.29)$$

The other upper bounds, relevant to our discussion, are known experimentally in the following channels:[27]

$$B(D^o \to \pi^+\pi^-) \leq 4.1\%$$
$$B(D^o \to K^+K^-) \leq 3.8\% \qquad (4.30)$$
$$B(D^o \to \pi^+\pi^- \text{ and } K^+K^-) \lesssim 5.1\%$$

These experimental bounds are not restrictive at all for GIM model. In the models of class a), however, the above decays are not Cabibbo suppressed, thus we have to examine what sort of bounds on branching ratios result from them.

First, we will show that the bound of (4.28) becomes less restrictive in the class a) models. For definiteness, let us assume that only the sextet is enhanced (the final conclusion does not depend on this assumption in any important way). We will further assume that $\langle 3||6_L||8_S\rangle \simeq \langle 3||6_{RL}||8_S\rangle$. One sees then from Table 4.4 that in the right-handed current model the $K^-\pi^+$ rate is about 1/10 of the rate for $D^0\to PP$. Thus one obtains the much less restrictive bound

$$B(D^0\to PP)\leq 30\% \tag{4.31}$$

We further notice that since $K^-\pi^+$ goes only via the 6_{LL} if one assumes $(\langle 3||6_{RL}||8_S\rangle)>(\langle 3||6_{LL}||8_S\rangle)$, following renormalization group arguments, then the bound becomes even less restrictive.

We analyze, now, the significance of the upper-limit (4.28) for the models of class b). As mentioned earlier, the situation will not differ significantly from the GIM scheme, unless if the $[3^*]_{RL}$ is substantially enhanced with respect to the $[6]_{RL,LL}$. If this enhancement does not take place, then one gets still a rather low upper-limit for $B(D^0\to PP)$.

Let us now examine the significance of the bounds given by (4.30) under various assumptions. First, in the case of 3*-dominance one obtains:

$$B(D^0\to\pi^+\pi^-)=B(D^0\to K^+K^-)\simeq 1/4\ B(D^0\to PP) \tag{4.32}$$

From (4.30), it follows that

$$B(D^0\to PP;\underline{3^*})\lesssim 10.2\% \tag{4.33}$$

On the other hand, if we assume sextet dominance with $\langle 3||6_{RL}||8_S\rangle \cong \langle 3||6_{LL}||8_S\rangle$, one gets:

$$B(D^0\to\pi^+\pi^-)=B(D^0\to K^+K^-)\approx 1/9\ B(D^0\to PP). \tag{4.34}$$

And we are led to the much less restrictive bound.

$$B(D^{o} \to PP;\underline{6}) \lesssim 23\% \qquad (4.35)$$

Finally, if we assume equal contributions of all reduced matrix elements we find,

$$B(D^{o} \to PP; \underline{3}^{*} \text{ and } \underline{6}) \lesssim 16\% \qquad (4.36)$$

From the experimental bound

$$B(D^{+} \to \bar{K}^{o}K^{+}) + B(F^{+} \to \bar{K}^{o}K^{+}) \leq 10\% \qquad (4.37)$$

and taking into account that (assuming again sextet dominance and $\langle 3 \| 6_{LL} \| 8_{S} \rangle \simeq \langle 3 \| 6_{RL} \| 8_{S} \rangle$:

$$B(D^{+} \to \bar{K}^{o}\, K^{+}) = 1/3\ B(D^{+} \to PP) \qquad (4.38)$$

$$B(F^{+} \to \bar{K}\ \ K^{+}) = 1/6\ B(D^{+} \to PP) \qquad (4.39)$$

one obtains

$$B(F^{+} \to PP) + 2B(D^{+} \to PP) \leq 60\% \qquad (4.40)$$

which is not a very restrictive bound.

The upper-limit obtained by Boyarski et al on $B(D^{+} \to \bar{K}^{o}\pi^{+}$ or $K^{-}\pi^{+}\pi^{+})$:

$$B(D^{+} \to \bar{K}^{o}\pi^{+} \text{ or } K^{-}\pi^{+}\pi^{+}) \leq 7.2\% \qquad (4.41)$$

does not present any serious difficulty. The point is that these decay modes are not supposed to be enhanced either in the GIM model or in the present right-handed current models.

Finally we note that if we assume that a heavy lepton is already contributing to σ(hadronic) all the bounds become somewhat less restrictive.

4.4 Concluding Remarks

We have alalyzed the various decays of the charmed mesons in the framework of R.H.C.C. For the case of purely leptonic decays we pointed out that they allow an easy distinction between the GIM model and each one of the models with R.H.C.C. Our discussion of purely leptonic decays is of little practical interest if $f_D=f_F=f_\pi$ since under this assumption these decays are expected to have much slower rate[21] than the non-leptonic ones However, if the leptonic decays are enhanced, as suggested in Ref. 25 our discussion becomes relevant.

For the non-leptonic decays we have pointed out that when the current $(\bar{\chi}n)_R$ is present there are decays of D^+ unsuppressed by the Cabbibo factor, and thus we expect the non-leptonic decays of D^+ to dominate over the semi-leptonic ones. With respect to the results of the Spear Charm search we have shown that the new current $(\bar{\chi}n)_R$ can accomodate the upper limit set for B $(D^o \to K^- \pi^+)$. In general, one notices that in models of class a) there are many more enhanced decays (each one with small branching ratio) than in the GIM case. This has implications for the current search for charm since it makes the detection of a charmed particle somewhat more difficult.

Decay mode	Mass Region (GeV/c^2) 1.50-1.85	1.85-2.40	2.40-4.00
$K^-\pi^+$ and $K^+\pi^-$	0.25	0.18	0.08
$K_S^0\pi^+\pi^-$	0.57	0.40	0.29
$\pi^+\pi^-$	0.13	0.13	0.09
K^+K^-	0.23	0.12	0.10
$K^-\pi^+\pi^+$ and $K^+\pi^+\pi^-$	0.51	0.49	0.19
$K_S^0\pi^+$ and $K_S^0\pi^-$	0.26	0.27	0.09
$K_S^0 K^+$ and $K_S^0 K^-$	0.54	0.33	0.09
$\pi^+\pi^-\pi^+$ and $\pi^+\pi^-\pi^-$	0.48	0.38	0.18
$K^\mp\pi^\pm$, $\bar{K}^0\pi^+\pi^-$, and $K^0\pi^+\pi^-$	1.16	0.90	0.58
K^+K^- and $\pi^+\pi^-$	0.23	0.16	0.15
$K^\mp\pi^\pm\pi^\pm$, $\bar{K}^0\pi^\pm$ and $K^0\pi^\pm$	0.64	0.51	0.30
$\bar{K}^0K^\pm$, $K^0K^\pm$, and $\pi^+\pi^-\pi^\pm$	1.10	0.76	0.29

TABLE 4.1. Largest upper limits at the 90% confidence level for inclusive production cross section times branching ratio, in nanobarns, for various mass regions. (From A. Boyarski et al Phys. Rev. Lett. 35 196 (1975).

TABLE 4.6

TABLE FOR REDUCED MATRIX ELEMENTS

$P,\alpha \equiv \langle 3 \| 6_{LL} \| 8_S \rangle_{\rho\rho,VV}$
$Q,\beta \equiv \langle 3 \| 6_{RL} \| 8_S \rangle_{\rho\rho,VV}$
$R,\gamma \equiv \langle 3 \| 3_{RL} \| 8_S \rangle_{\rho\rho,VV}$
$R',\gamma' \equiv \langle 3 \| 3_{RL} \| 1 \rangle_{\rho\rho,VV}$
$D,D' \equiv \langle 3 \| 6_{LL,RL} \| 10 \rangle_{\rho V,B\bar{B}}$
$S,S' \equiv \langle 3 \| 6_{LL,RL} \| 8_S \rangle_{\rho V,B\bar{B}}$
$S'' \equiv \langle 3 \| 3_{RL} \| 8_S \rangle_{\rho V,B\bar{B}}$
$A,A' \equiv \langle 3 \| 6_{LL,RL} \| 8_A \rangle_{\rho V,B\bar{B}}$
$A'' \equiv \langle 3 \| 3_{RL} \| 8_A \rangle_{\rho V,B\bar{B}}$
$S''' \equiv \langle 3 \| 3_{RL} \| 1 \rangle_{\rho V,B\bar{B}}$

TABLE 4.2

RELATIVE DECAY RATES OF F^+

$F^+ \to \mathcal{P}\mathcal{P}$		$F^+ \to VV$		$F^+ \to \mathcal{P} V$	$F^+ \to \bar{B}B$
$\cos^4\theta \; \lvert P \rvert^2 \; x$		$\cos^4\theta \; \lvert \alpha \rvert^2$		$\cos^4\theta \; x$	$\cos^4\theta \; x$
$\pi^+\eta'$	$\frac{4}{9}$	$\rho^+\omega$	$\frac{2}{3}$	$\pi^+\omega \; \lvert -\frac{D}{\sqrt{2}} - \sqrt{2}\, S \rvert^2$	$\Sigma^+\bar{\Lambda}$ $\frac{1}{6}\lvert 3D$
$K^+\bar{K}^0$	$\frac{1}{3}$	$K^{*+}\overline{K^{*0}}$	$\frac{1}{3}$	$\pi^+\rho^0 \; \lvert \frac{D}{\sqrt{2}} - \sqrt{2}\, A \rvert^2$	$\Sigma^+\bar{\Sigma}^0$ $\frac{1}{2}\lvert D-$
$\pi^+\eta$	$\frac{2}{9}$			$\pi^+\phi \; \lvert D \rvert^2$	$\Lambda\bar{\Sigma}^-$ $\frac{1}{6}\lvert 3D$
$\cos^2\theta \; \lvert Q+R \rvert^2 \; x$		$\cos^2\theta \; \lvert \beta+\gamma \rvert^2$		$\eta\rho^+ \; \frac{1}{6}\lvert (3D-2S) \rvert^2$	$\Sigma^0\bar{\Sigma}^-$ $\frac{1}{2}\lvert D-$
$K^+\eta'$	$\frac{4}{9}$	$K^{*+}\phi$	$\frac{1}{3}$	$\pi^0\rho^+ \; \frac{1}{2}\lvert (D-2A) \rvert^2$	$p\bar{n}$ $\lvert D-S$
π^+K^0	$\frac{1}{3}$	$K^{*0}\rho^+$	$\frac{1}{3}$	$K^+\overline{K^{*0}} \; \lvert D-S+A \rvert^2$	$\Xi^0\bar{\Xi}^-$ $\lvert D+S$
π^0K^+	$\frac{1}{6}$	$K^{*+}\rho^0$	$\frac{1}{6}$	$\bar{K}^0K^{*+} \; (D+S+A) \; 2$	$\cos^2\theta \; x$
$K^+\eta$	$\frac{1}{18}$	$K^{*+}\omega$	$\frac{1}{6}$	$\eta'\rho^+ \; \frac{4}{3}\lvert S \rvert^2$	$p\bar{\Sigma}^0$ $\frac{1}{2}\lvert 2D'-S'-S''-$
				$\cos^2\theta \; x$	$p\bar{\Lambda}$ $\frac{1}{6}\lvert 6D'-S'-S''+3$
				$K^+\omega \; \frac{1}{2}\lvert 2D'-S'-S''-A'-A'' \rvert^2$	$\Sigma^0\bar{\Xi}^-$ $\frac{1}{2}\lvert 2D'+S'+S''$
				$K^+\rho^0 \; \frac{1}{2}\lvert 2D'+S'+S''+A'+A'' \rvert^2$	$\Lambda\bar{\Xi}^-$ $\frac{1}{6}\lvert 6D'+S'+S''$
				$K^+\phi \; \lvert 2D'-S'-S''+A'+A'' \rvert^2$	$\Sigma^+\bar{\Xi}^0$ $\lvert -2D'-S'-S''+$
				$\pi^0K^{*+} \; \frac{1}{2}\lvert 2D'+S'+S''-A'-A'' \rvert^2$	$n\bar{\Sigma}^-$ $\lvert 2D'-S'-S''-A$
				$\eta K^{*+} \; \frac{1}{6}\lvert 6D'+S'+S''+3A'+3A'' \rvert^2$	
				$\pi^+K^{*0} \; \lvert -2D'-S'-S''+A'+A'' \rvert^2$	
				$K^0\rho^+ \; \lvert 2D'-S'-S''-A'-A'' \rvert^2$	
				$\eta'K^{*+} \; \frac{4}{3}\lvert S'+S'' \rvert^2$	

TABLE 4.2

RELATIVE DECAY RATES OF F^+

P	$F^+ \to VV$	$F^+ \to \mathcal{P}V$	$F^+ \to \bar{B}B$
^{2} x	$\cos^4\theta\ \lvert\alpha\rvert^2$	$\cos^4\theta$ x	$\cos^4\theta$ x
$\frac{4}{9}$	$\rho^+\omega$ $\frac{2}{3}$	$\pi^+\omega$ $\lvert -\frac{D}{\sqrt{2}} - \sqrt{2}\ S\rvert^2$	$\Sigma^+\bar{\Lambda}$ $\frac{1}{6}\lvert 3D+2S\rvert^2$
$\frac{1}{3}$	$K^{*+}\overline{K^{*o}}$ $\frac{1}{3}$	$\pi^+\rho^o$ $\lvert\frac{D}{\sqrt{2}} - \sqrt{2}\ A\rvert^2$	$\Sigma^+\bar{\Sigma}^o$ $\frac{1}{2}\lvert D-2A\rvert^2$
$\frac{2}{9}$		$\pi^+\phi$ $\lvert D\rvert^2$	$\Lambda\bar{\Sigma}^-$ $\frac{1}{6}\lvert 3D-2S\rvert^2$
$R\rvert^2$ x	$\cos^2\theta\ \lvert\beta+\gamma\rvert^2$	$\eta\rho^+$ $\frac{1}{6}\lvert(3D-2S)\rvert^2$	$\Sigma^o\bar{\Sigma}^-$ $\frac{1}{2}\lvert D-2A\rvert^2$
$\frac{4}{9}$	$K^{*+}\phi$ $\frac{1}{3}$	$\pi^o\rho^+$ $\frac{1}{2}\lvert(D-2A)\rvert^2$	$p\bar{n}$ $\lvert D-S+A\rvert^2$
$\frac{1}{3}$	$K^{*o}\rho^+$ $\frac{1}{3}$	$K^+\overline{K^{*o}}$ $\lvert D-S+A\rvert^2$	$\Xi^o\bar{\Xi}^-$ $\lvert D+S+A\rvert^2$
$\frac{1}{6}$	$K^{*+}\rho^o$ $\frac{1}{6}$	$\bar{K}^oK^{*+}$ $(D+S+A)^2$	$\cos^2\theta$ x
$\frac{1}{18}$	$K^{*+}\omega$ $\frac{1}{6}$	$\eta'\rho^+$ $\frac{4}{3}\lvert S\rvert^2$	$p\bar{\Sigma}^o$ $\frac{1}{2}\lvert 2D'-S'-S''-A'-A''\rvert^2$
		$\cos^2\theta$ x	$p\bar{\Lambda}$ $\frac{1}{6}\lvert 6D'-S'-S''+3A'+3A''\rvert^2$
		$K^+\omega$ $\frac{1}{2}\lvert 2D'-S'-S''-A'-A''\rvert^2$	$\Sigma^o\bar{\Xi}^-$ $\frac{1}{2}\lvert 2D'+S'+S''-A'-A''\rvert^2$
		$K^+\rho^o$ $\frac{1}{2}\lvert 2D'+S'+S''+A'+A''\rvert^2$	$\Lambda\bar{\Xi}^-$ $\frac{1}{6}\lvert 6D'+S'+S''+3A'+3A''\rvert^2$
		$K^+\phi$ $\lvert 2D'-S'-S''+A'+A''\rvert^2$	$\Sigma^+\bar{\Xi}^o$ $\lvert -2D'-S'-S''+A'+A''\rvert^2$
		π^oK^{*+} $\frac{1}{2}\lvert 2D'+S'+S''-A'-A''\rvert^2$	$n\bar{\Sigma}^-$ $\lvert 2D'-S'-S''-A'-A''\rvert^2$
		ηK^{*+} $\frac{1}{6}\lvert 6D'+S'+S''+3A'+3A''\rvert^2$	
		π^+K^{*o} $\lvert -2D'-S'-S''+A'+A''\rvert^2$	
		$K^o\rho^+$ $\lvert 2D'-S'-S''-A'-A''\rvert^2$	
		$\eta'K^{*+}$ $\frac{4}{3}\lvert S'+S''\rvert^2$	

TABLE 4.3

RELATIVE DECAY RATES OF D^+

$D^+ \to PP$		$D^+ \to VV$		$D^+ \to PV$		D^+	
$\cos^2\theta\,\lvert(Q+R)+P\sin\theta\rvert^2 \times$		$\cos^2\theta\,\lvert(\beta+\gamma)+\alpha\sin\theta\rvert^2 \times$		$\cos^4\theta \times$		co	
$\pi^+\eta'$	$\frac{4}{9}$	$\rho^+\omega$	$\frac{2}{3}$	$\bar{K}^o\rho^+$	$\lvert 3D\rvert^2$	$\Xi^o\bar{\Sigma}^-$	
$K^+\bar{K}^o$	$\frac{1}{3}$	$K^{*+}K^o$	$\frac{1}{3}$	$\pi^+\overline{K^{*o}}$	$\lvert 3D\rvert^2$	$\Sigma^+\bar{n}$	
$\pi^+\eta$	$\frac{2}{9}$			$\cos^2\theta \times$		co	
$\sin^2\theta\,\lvert Q+P\sin\theta\rvert^2 \times$		$\sin^2\theta\,\lvert\beta+\alpha\sin\theta\rvert^2 \times$		$\pi^+\phi$	$4\lvert D'\rvert^2$	$\Sigma^+\bar{\Sigma}^o$	
$K^+\eta'$	$\frac{4}{9}$	$K^{*+}\phi$	$\frac{1}{3}$	$\pi^+\rho^o$	$2\lvert D'+A'+A''\rvert^2$	$\Sigma^+\bar{\Lambda}$	$\frac{2}{3}\lvert$
$K^o\pi^+$	$\frac{1}{3}$	$K^{*o}\rho^+$	$\frac{1}{3}$	$\pi^+\omega$	$2\lvert -D'+S'+S''\rvert^2$	$\Lambda\bar{\Sigma}^-$	$\frac{2}{3}\lvert$
$K^+\pi^o$	$\frac{1}{6}$	$K^{*+}\rho^o$	$\frac{1}{6}$	$\eta\rho^+$	$\frac{2}{3}\lvert 3D'+S'+S''\rvert^2$	$\Sigma^o\bar{\Sigma}^-$	2
$K^+\eta$	$\frac{1}{18}$	$K^{*+}\omega$	$\frac{1}{6}$	$\pi^o\rho^+$	$2\lvert D'+A'+A''\rvert^2$	$p\bar{n}$	$\lvert 2D$
				$K^+\overline{K^{*o}}$	$\lvert 2D'+S'+S''-A'-A''\rvert^2$	$\Xi^o\bar{\Xi}^-$	$\lvert -2$
				$\overline{K^o}K^{*+}$	$\lvert -2D'+S'+S''+A'+A''\rvert^2$		
				$\eta'\rho^+$	$\frac{4}{3}\lvert S'+S''\rvert^2$		

TABLE 4.3

RELATIVE DECAY RATES OF D^+

$\to PP$	$D^+ \to VV$		$D^+ \to PV$		$D^+ \to \bar{B}B$	
$\text{-}R)+P\sin\theta\|^2$ x	$\cos^2\theta\|(\beta+\gamma)+\alpha\sin\theta\|^2$ x		$\cos^4\theta$ x		$\cos^4\theta$ x	
$\frac{4}{9}$	$\rho^+\omega$	$\frac{2}{3}$	$\bar{K}^o\rho^+$	$\|3D\|^2$	$\Xi^o\bar{\Sigma}^-$	$\|3D\|^2$
$\frac{1}{3}$	$K^{*+}K^o$	$\frac{1}{3}$	$\pi^+\overline{K^{*o}}$	$\|3D\|^2$	$\Sigma^+\bar{n}$	$\|3D\|^2$
$\frac{2}{9}$			$\cos^2\theta$ x		$\cos^2\theta$ x	
$\text{-}\sin\theta\|^2$ x	$\sin^2\theta\|\beta+\alpha\sin\theta\|^2$ x		$\pi^+\phi$	$4\|D'\|^2$	$\Sigma^+\bar{\Xi}^o$	$2\|D'+A'+A''\|^2$
$\frac{4}{9}$	$K^{*+}\phi$	$\frac{1}{3}$	$\pi^+\rho^o$	$2\|D'+A'+A''\|^2$	$\Sigma^+\bar{\Lambda}$	$\frac{2}{3}\|-3D'+S'+S''\|^2$
$\frac{1}{3}$	$K^{*o}\rho^+$	$\frac{1}{3}$	$\pi^+\omega$	$2\|-D'+S'+S''\|^2$	$\Lambda\bar{\Sigma}^-$	$\frac{2}{3}\|3D'+S'+S''\|^2$
$\frac{1}{6}$	$K^{*+}\rho^o$	$\frac{1}{6}$	$\eta\rho^+$	$\frac{2}{3}\|3D'+S'+S''\|^2$	$\Sigma^o\bar{\Sigma}^-$	$2\|D'+A'+A''\|^2$
$\frac{1}{18}$	$K^{*+}\omega$	$\frac{1}{6}$	$\pi^o\rho^+$	$2\|D'+A'+A''\|^2$	$p\bar{n}$	$\|2D'+S'+S''-A'-A''\|^2$
			$K^+\overline{K^{*o}}$	$\|2D'+S'+S''-A'-A''\|^2$	$\Xi^o\overline{\Xi^-}$	$\|-2D'+S'+S''+A'+A''\|^2$
			$\overline{K^o}K^{*+}$	$\|-2D'+S'+S''+A'+A''\|^2$		
			$\eta'\rho^+$	$\frac{4}{3}\|S'+S''\|^2$		

TABLE 4.4

RELATIVE DECAYS OF D^0

$D^0 \to \mathcal{P}\mathcal{P}$	$D^0 \to VV$	$D^0 \to \mathcal{P}V$	D
$\cos^4\theta\ \lvert P\rvert^2$ x	$\cos^4\theta$ x	$\cos^4\theta$	c
$\eta'\bar K^0\quad \frac{4}{3}$	$K^{*-}\rho^+\quad \lvert\alpha\rvert^2$	$\bar K^0\rho^0\ \frac{1}{2}\lvert 2D+S-A\rvert^2$	$\Xi^0\bar\Sigma^0\ \frac{1}{2}\lvert 2$
$\pi^+K^-\quad 1$	$\bar K^{*0}\phi\quad \lvert\alpha\rvert^2$	$\pi^0\bar K^{*0}\ \frac{1}{2}\lvert 2D-S-A\rvert^2$	$\Xi^0\bar\Lambda\ \frac{1}{6}\lvert S$
$\pi^0\bar K^0\quad \frac{1}{2}$	$\bar K^{*0}\rho^0\quad \frac{1}{2}\lvert\alpha\rvert^2$	$K^-\rho^+\ \lvert -D+S-A\rvert^2$	$\Sigma^0\bar n\ \frac{1}{2}\lvert 2$
$\eta\bar K^0\quad \frac{1}{6}$	$\bar K^{*0}\omega\quad \frac{1}{2}\lvert\alpha\rvert^2$	$\pi^+K^{*-}\lvert D+S+A\rvert^2$	$\Lambda\bar n\ \frac{1}{6}\lvert -$
$\cos^2\theta$ x	$\cos^2\theta$ x	$\eta\bar K^{*0}\ \frac{1}{6}\lvert -S+3A\rvert^2$	$\Xi^-\bar\Sigma^-\ \lvert -$
$\eta'\eta'\quad \lvert R'\rvert^2$	$\phi\phi\quad \lvert -\beta-\frac{1}{3}\gamma+\gamma'\rvert^2$	$\eta'\bar K^{*0}\ \frac{4}{3}\lvert S\rvert^2$	$\Sigma^+\bar p\ \lvert D$
$\pi^+\pi^-\ \lvert Q+\frac{1}{3}R+R'\rvert^2$	$\rho^+\rho^-\lvert \beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar K^0\phi\ \lvert S+A\rvert^2$	(
$K^+K^-\ \lvert -Q+\frac{1}{3}R+R'\rvert^2$	$K^{*+}K^{*-}\lvert -\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar K^0\omega\ \frac{1}{2}\lvert S-A\rvert^2$	$n\bar n\lvert -D'+2A'$
$\pi^0\pi^0\ \lvert \frac{1}{2}Q+\frac{1}{6}R+R'\rvert^2$	$\rho^0\rho^0\lvert \frac{1}{2}\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\cos^2\theta$ x	$\Xi^0\bar\Xi^0\lvert D'-2A'$
$\eta\eta\ \lvert \frac{1}{2}Q-\frac{1}{6}R+R'\rvert^2$	$\omega\omega\lvert \frac{1}{2}\beta+\frac{1}{6}\gamma+\gamma'\rvert^2$	$\pi^-\rho^+\lvert D'-S'+A'+\frac{1}{3}S''+A''+S'''\rvert^2$	$\Xi^-\bar\Xi^-\ \lvert -D'+S'$
$\eta\eta'\lvert \sqrt{2}Q+\frac{\sqrt{2}}{3}R\rvert^2$	$\omega\rho^0\lvert -\beta+\gamma\rvert^2$	$K^0\bar K^{*0}\lvert -D'+2A'-\frac{2}{3}S''+S'''\rvert^2$	$p\bar p\lvert D'+S'+A$
$\eta\pi^0\lvert -\frac{Q}{\sqrt{3}}+\frac{1}{3\sqrt{3}}R\rvert^2$	$K^{0*}\overline{K^{0*}}\lvert -\frac{2}{3}\gamma+\gamma'\rvert^2$	$K^+K^{*-}\lvert D'+S'+A'-\frac{1}{3}S''-A''+S'''\rvert^2$	$\Sigma^+\Sigma^+\lvert -D'$
$\eta'\pi^0\lvert -\frac{2}{\sqrt{6}}Q+\frac{2}{3\sqrt{6}}R\rvert^2$		$K^-K^{*+}\lvert -D'+S'-A'+\frac{1}{3}S''+A''+S'''\rvert^2$	$\Sigma^-\Sigma^-\lvert D'-S'$
$K^0\bar K^0\lvert -\frac{2}{3}R+R'\rvert^2$		$\bar K^0K^{*0}\lvert D'-2A'-\frac{2}{3}S''+S'''\rvert^2$	$\Sigma^0\bar\Lambda\ \frac{1}{3}\lvert -3D'$
		$\pi^+\rho^-\lvert -D'-S'-A'+\frac{1}{3}S''-A''+S'''\rvert^2$	$\Lambda\bar\Sigma^0\ \frac{1}{3}\lvert -D'+$
		$\pi^0\phi\quad 2\lvert D'\rvert^2$	$\Sigma^0\bar\Sigma^0\lvert -S'+\frac{1}{3}$
		$\eta\phi\ \frac{4}{3}\lvert \sqrt{2}S'+\frac{\sqrt{2}}{3}S''\frac{S'''}{\sqrt{2}}\rvert^2$	$\Lambda\bar\Lambda\ \lvert S'-\frac{1}{3}S'$
		$\eta'\phi\ \frac{1}{3}\lvert 2S'-\frac{2}{3}S''+S'''\rvert^2$	
		$\eta'\rho^0\ \frac{2}{3}\lvert -S'+S''\rvert^2$	
		$\eta'\omega\ \frac{2}{3}\lvert S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\eta\rho^0\ \frac{1}{3}\lvert -3D'+S'+S''\rvert^2$	
		$\eta\omega\ \frac{1}{3}\lvert -S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\pi^0\rho^0\ \lvert -S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\pi^0\omega\ \lvert -D'+S'+S''\rvert^2$	

TABLE 4.4

RELATIVE DECAYS OF D^o

	$D^o \to VV$	$D^o \to \mathcal{P}V$	$D^o \to B\bar{B}$
x	$\cos^4\theta$ x	$\cos^4\theta$	$\cos^4\theta$ x
	$K^{*-}\rho^+ \quad \lvert\alpha\rvert^2$	$\bar{K}^o\rho^o \ \frac{1}{2}\lvert 2D+S-A\rvert^2$	$\Xi^o\bar{\Sigma}^o \ \frac{1}{2}\lvert 2D+S-A\rvert^2$
	$\overline{K^{*o}}\phi \quad \lvert\alpha\rvert^2$	$\pi^o\overline{K^{*o}} \ \frac{1}{2}\lvert 2D-S-A\rvert^2$	$\Xi^o\bar{\Lambda} \ \frac{1}{6}\lvert S+3A\rvert^2$
	$\overline{K^{*o}}\rho^o \quad \frac{1}{2}\lvert\alpha\rvert^2$	$K^-\rho^+ \ \lvert -D+S-A\rvert^2$	$\Sigma^o\bar{n} \ \frac{1}{2}\lvert 2D-S-A\rvert^2$
	$\bar{K}^{*o}\omega \quad \frac{1}{2}\lvert\alpha\rvert^2$	$\pi^+K^{*-}\lvert D+S+A\rvert^2$	$\Lambda\bar{n} \ \frac{1}{6}\lvert -S+3A\rvert^2$
	$\cos^2\theta$ x	$\eta\overline{K^{*o}} \ \frac{1}{6}\lvert -S+3A\rvert^2$	$\Xi^-\overline{\Sigma^-} \ \lvert -D+S-A\rvert^2$
$\rvert^2$	$\phi\phi \ \lvert -\beta-\frac{1}{3}\gamma+\gamma'\rvert^2$	$\eta'\overline{K^{*o}} \ \frac{4}{3}\lvert S\rvert^2$	$\Sigma^+\bar{p} \ \lvert D+S+A\rvert^2$
$R'\rvert^2$	$\rho^+\rho^-\lvert\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar{K}^o\phi \ \lvert S+A\rvert^2$	$\cos^2\theta$ x
$+R'\rvert^2$	$K^{*+}K^{*-}\lvert -\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar{K}^o\omega \ \frac{1}{2}\lvert S-A\rvert^2$	$n\bar{n}\lvert -D'+2A'-\frac{2}{3}S''+S'''\rvert^2$
$+R'\rvert^2$	$\rho^o\rho^o\lvert\frac{1}{2}\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\cos^2\theta$ x	$\Xi^o\overline{\Xi^o}\lvert D'-2A'-\frac{2}{3}S''+S'''\rvert^2$
	$\omega\omega\lvert\frac{1}{2}\beta+\frac{1}{6}\gamma+\gamma'\rvert^2$	$\pi^-\rho^+\lvert D'-S'+A'+\frac{1}{3}S''+A''+S'''\rvert^2$	$\Xi^-\overline{\Xi^-}\lvert -D'+S'-A'+\frac{1}{3}S''+S'''\rvert^2$
	$\omega\rho^o\lvert -\beta+\gamma\rvert^2$	$K^o\overline{K^{*o}}\lvert -D'+2A'-\frac{2}{3}S''+S'''\rvert^2$	$p\bar{p}\lvert D'+S'+A'+\frac{1}{3}S''+A''+S'''\rvert^2$
$R\rvert^2$	$K^{o*}\overline{K^{o*}}\lvert -\frac{2}{3}\gamma+\gamma'\rvert^2$	$K^+K^{*-}\lvert D'+S'+A'-\frac{1}{3}S''-A''+S'''\rvert^2$	$\Sigma^+\Sigma^+\lvert -D' \quad -A'+\frac{1}{3}S''+A''+S'''\rvert^2$
$_6R\rvert^2$		$K^-K^{*+}\lvert -D'+S'-A'+\frac{1}{3}S''+A''+S'''\rvert^2$	$\Sigma^-\Sigma^-\lvert D'-S'+A'+\frac{1}{3}S''-A''+S'''\rvert^2$
		$\bar{K}^oK^{*o}\lvert D'-2A'-\frac{2}{3}S''+S'''\rvert^2$	$\Sigma^o\bar{\Lambda} \ \frac{1}{3}\lvert -3D'+S'+S''\rvert^2$
		$\pi^+\rho^-\lvert -D'-S'-A'+\frac{1}{3}S''-A''+S'''\rvert^2$	$\Lambda\bar{\Sigma}^o \ \frac{1}{3}\lvert -D'+S'+S''\rvert^2$
		$\pi^o\phi \quad 2\lvert D'\rvert^2$	$\Sigma^o\bar{\Sigma}^o\lvert -S'+\frac{1}{3}S''+S'''\rvert^2$
		$\eta\phi \ \frac{4}{3}\lvert\sqrt{2}\,S'+\frac{\sqrt{2}}{3}S''\frac{S'''}{\sqrt{2}}\rvert^2$	$\Lambda\bar{\Lambda} \ \lvert S'-\frac{1}{3}S''+S'''\rvert^2$
		$\eta'\phi \ \frac{1}{3}\lvert 2S'-\frac{2}{3}S''+S'''\rvert^2$	
		$\eta'\rho^o \ \frac{2}{3}\lvert -S'+S''\rvert^2$	
		$\eta'\omega \ \frac{2}{3}\lvert S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\eta\rho^o \ \frac{1}{3}\lvert -3D'+S'+S''\rvert^2$	
		$\eta\omega \ \frac{1}{3}\lvert -S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\pi^o\rho^o \ \lvert -S'+\frac{1}{3}S''+S'''\rvert^2$	
		$\pi^o\omega \ \lvert -D'+S'+S''\rvert^2$	

TABLE 4.5

RELATIVE DECAY RATES OF THE "z" TYPE HEAVY MESONS

$F_z^0 \to \mathcal{P}\mathcal{P}$		$D_z^- \to \mathcal{P}\mathcal{P}$		$D_z^0 \to \mathcal{P}\mathcal{P}$	
$\cos^2\theta\ \lvert Q+R\rvert^2$ x		$\cos^2\theta$ x $\lvert Q+R\rvert^2$		$\cos^2\theta$ x	
$K^+\pi^-$	$\frac{1}{3}$	$\pi^-\eta$	$\frac{2}{9}$	K^+K^-	$\lvert\frac{1}{3}Q-\frac{2}{3}R+R'\rvert^2$
$K^0\pi^0$	$\frac{1}{6}$	$\pi^-\eta'$	$\frac{4}{9}$	$K^0\bar{K}^0$	$\lvert\frac{1}{3}Q+\frac{1}{3}R+R'\rvert^2$
$K^0\eta$	$\frac{1}{18}$	K^-K^0	$\frac{1}{3}$	$\pi^+\pi^-$	$\lvert-\frac{2}{3}Q+\frac{1}{3}R+R'\rvert^2$
$K^0\eta'$	$\frac{4}{9}$			$\eta\eta$	$\lvert\frac{1}{3}Q-\frac{1}{6}R+R'\rvert^2$
				$\pi^0\pi^0$	$\lvert-\frac{1}{3}Q+\frac{1}{6}R+R'\rvert^2$
				$\eta\eta'$	$\lvert\frac{4}{3\sqrt{2}}Q+\frac{4}{9\sqrt{2}}R\rvert^2$
				$\eta'\eta'$	$\lvert R'\rvert^2$
				$\pi^0\eta$	$\frac{1}{3}\lvert R\rvert^2$
				$\pi^0\eta'$	$\frac{2}{3}\lvert R\rvert^2$

Chapter V

Structure of the $\Delta S=1$ non-leptonic Hamiltonian

5.1 Introductory Remarks

As we have seen the new right-handed currents couple the light quarks with the heavy quarks and this has interesting consequences for the semi-leptonic and non-leptonic decays of the new charmed particles. As far as the usual non-charmed particles are concerned, their semi-leptonic decays are obviously unaffected by the new currents. However, the structure of the $\Delta S=1$ non-leptonic Hamiltonian is drastically changed. It is then pertinent to ask if these changes are compatible with constraints derived from presently available data. At this point it is important to emphasize that in the standard model there are some experimental facts, for which one does not find any convincing justifications, such as the $\Delta I=\frac{1}{2}$ rule for non-leptonic decays, the relative size of the s and p waves in hyperon decays, etc. Constraints on the structure of the $\Delta S=1$ non-leptonic Hamiltonian, may come from various areas:

a) $\Delta I=\frac{1}{2}$ rule

As we have mentioned, the revival of interest in right-handed currents, was in great part due to the work of De Rujula, Giorgi and Glashow who have pointed out that the Mohapatra charmed current $(\bar{\chi}n)_R$ provides an elegant explanation for the observed enhancement of the non-leptonic decays and the well satisfied $\Delta I=\frac{1}{2}$ rule.

b) $K_S \to 2\pi$ decays

It is well known that in the standard model the decay $K_o^S \to 2\pi$, one of the fastest known weak decays, is forbidden in the limit of exact SU(3) symmetry. Recall that this is a consequence of the fact that the parity violating part of the $\Delta S=1$ effective Hamiltonian transforms as λ_6. This apparent puzzle is solved in models with R.H.C.C., since for these models:

$$\left[\mathcal{H}_{eff}^{p.v.}\right]^{\Delta S=1} \sim \lambda_7$$

and the decay $K\to 2\pi$ is no longer forbidden.

c) Radiative hyperon decays

The study of two body radiative decays such as $\Sigma^+\to P\gamma$, $\Xi^-\to\Sigma^-\gamma$ provide interesting information about the $\Delta S=1$ effective Hamiltonian ($\mathcal{H}_{eff}^{S=1}$). It can be easily shown that in the standard model the U spin transformation properties of $\mathcal{H}_{eff}^{\Delta S=1}$ are such that the parity violating part of the amplitudes for $\Sigma^+\to P\gamma$, $\Xi^-\to\Sigma^-\gamma$ should vanish in the limit of U spin symmetry. This implies a vanishing asymmetry parameter α, in disagreement with the experimental value $\alpha\simeq -1$. Even if one allows for U-spin breaking effects, a detailed analysis suggests a rather small asymmetry parameter, in the framework of the standard model. On the other hand, models with the right-handed current $(\bar{\chi}n)_R$ predict $\alpha = -1$ and are thus favored by present experiments.[28]

e) $\Delta I=3/2$ effects in K meson decays.

It has been pointed out by Golowich and Holstein that the study of the $\Delta I=3/2$ amplitude in $K\to 3\pi$ does not favor models with a right-handed current of the $(\bar{\chi}n)_R$ type. It must, however, be stressed that this is a small effect as compared to the enhanced $\Delta I=\frac{1}{2}$ amplitudes and that the conclusions of ref. 29 may be substantially changed if : a) the final state interactions in $K\to 3\pi$ turn out to be important. b)there is a significant part of the $\Delta I=3/2$ weak Hamiltonian coming from either (i) the electromagnetic interactions or (ii) from the Higg's sector of the complete gauge theory Lagrangian. In particular, one can show in the context of a simple SU(2) x U(1) model with right-handed currents, that it is possible to have a large $\Delta I=3/2$ piece in the effective weak Hamiltonian coming from the Higg's sector. This introduces mixed chiral structures into $\mathcal{H}^{3/2}$ and invalidates the conclusions of ref. 29.

f) Non-leptonic hyperon decays

The study of s and p wave amplitudes in hyperon decays provides another important test on the structure of $\mathcal{H}_{eff}^{\Delta S=1}$. In this chapter we will analyze this question in some detail and in particular we will show that there is no conflict between the new current $(\bar{\chi}n)_R$ and the experimental value for the hyperon s and p wave amplitudes.

5.2 Evaluation of the Non-leptonic Hyperon Amplitudes

There exist various competing schemes for the study of the non-leptonic decays of hyperons.[30] One scheme based on current algebra PCAC and soft-pion extrapolation provides a good description of the hyperon s-wave amplitudes,[31] but fares rather badly in predicting correct magnitudes when applied to p-wave amplitudes.[32,33] Another approach consists of assuming K* dominance and is able to give a good fit for the s-wave amplitude. Since the K* amplitude vanishes in the soft-pion limit, while it may give a significant contribution to a physical pion, one is then faced with a somewhat ambiguous situation, in which two independent terms provide separately a good fit for the s-wave amplitude. It is then not clear if one should take into account simultaneously these two contributions and if one did that what their relative magnitude should be.

Corrections to the standard current algebra contribution have been considered by various authors. In particular, Pati and Kumar have shown that by taking into account SU(3) breaking effects in mass one is able to improve the fit of the p-wave amplitudes. Using a different approach Okubo[35] has shown that corrections to the s-wave amplitude have the same form as the K* contribution.

In our analysis, we will take the point of view that one should consider both the current algebra and K* contributions (or K*-like corrections arising from Okubo's approach) in evaluating the hyperon s-wave amplitudes The relative magnitude of these two contributions will however be kept as a free parameter.

In models with the $(\bar{\chi}n)_R$ current the effective non-leptonic weak Hamiltonian can be written:

$$\mathcal{H}_{eff}^{\Delta S=1} = \left[\mathcal{H}_{eff}\right]_{LL}^{\Delta S=1} + \left[\mathcal{H}_{eff}\right]_{RL}^{\Delta S=1} \tag{5.1}$$

where

$$\left[\mathcal{H}_{eff}\right]_{LL}^{\Delta S=1} = \sin\theta_c \cos\theta_c \left[(\bar{n}_L \gamma_\mu p_L)(\bar{p}_L \gamma_\mu \lambda_L) - (\bar{n}_L \gamma_\mu \chi_L)(\bar{\chi}_L \gamma_\mu \lambda_L) + h.c. \right],$$

$$\left[\mathcal{H}_{eff}\right]_{RL}^{\Delta S=1} = \cos\theta_c \left[(\bar{n}_R \gamma_\mu \chi_R)(\bar{\chi}_L \gamma_\mu \lambda_L) + h.c. \right] . \tag{5.2}$$

The **LL** part of the Hamiltonian is already present in the standard model and contributes both to the $\Delta I=\frac{1}{2}$ and $\Delta I=3/2$ transitions. The RL part is the new term which is not Cabbibo supressed and it contributes only to $\Delta I=\frac{1}{2}$ transitions. It can be easily seen from ( ) that $\left[\mathcal{H}_{eff}^{p.c.}\right]_{RL}$ and $\left[\mathcal{H}_{eff}^{p.v}\right]_{RL}$ transform respectively as the 6th and 7th components of an octet. This is to be contrasted with the standard model, in which both the parity violating and parity conserving parts of $\left[\mathcal{H}_{eff}\right]_{LL}$ transform as λ_6.

The amplitude for the decay:

$$B_j(p) \longrightarrow B_k(p') + \pi_i(k) \tag{5.3}$$

can be written:

$$\begin{aligned} T^{ijk} &= (2k^0)^{\frac{1}{2}} \left\langle B_k(p')\, \pi_i(k) \right| H_w(0) \left| B_j(p) \right\rangle \\ &= \bar{u}(p')\left(A^{ijk} + B^{ijk}\gamma_5\right)u(p) \end{aligned} \tag{5.4}$$

where A and B are the S and P wave amplitudes, respectively. Using the low energy theorem for soft pions, one obtains:

$$\lim_{k\to 0} T^{ijk} = -\frac{\sqrt{2}}{f_\pi}\left\langle B_k\right|\left[F_i^5, H_w(0)\right]\left|B_j\right\rangle + \frac{\sqrt{2}}{f_\pi}\lim_{k\to 0} k^\mu M_\mu^{ijk} \tag{5.5}$$

where

$$M_\mu^{ijk} = \int d^4x\, e^{-ikx} \left\langle B_k\right|\left[A_\mu^i(x), H_w(0)\right]\left|B_j\right\rangle \theta(x^0) \tag{5.6}$$

Following the usual procedure, we separate the baryon-pole contribution from the rest of the amplitude:

$$T^{ijk}(k) = B^{ijk}(k) + R^{ijk}(k) \tag{5.7}$$

where B^{ijk} is the baryon pole term. It follows then that:

$$\lim_{k\to 0} R^{ijk} = -\frac{\sqrt{2}}{f_\pi}\langle B_k|[F_i^5, H_w]|B_j\rangle + \lim_{k\to 0}\frac{1}{f_\pi}\left[\sqrt{2}\,k^\mu M_\mu^{ijk} - f_\pi B^{ijk}\right] \tag{5.8}$$

The usual soft-pion approximation consists of assuming that $R^{ijk}(K)$ is a slowly varying function of K, so that its physical value may be replaced by $R^{ijk}(0)$. We will, however, keep the K* contribution since it vanishes in the soft-pion limit, while it can give a significant contribution for a physical pion. One thus obtains:

$$T^{ijk}(k\,;\,k^2 = m_\pi^2) \cong -\frac{\sqrt{2}}{f_\pi}\langle B_k|[F_i^5(0), H_w(0)]|B_j\rangle + k^{*ijk}(k^2 = m_\pi^2) +$$

$$+\ B^{ijk}(k^2 = m_\pi^2) + \lim_{k\to 0}\frac{1}{f_\pi}\left[\sqrt{2}\ k^\mu\ M_\mu^{ijk} - f_\pi\, B^{ijk}\right]. \tag{5.9}$$

The first two terms in (5.9) contribute to the s-wave amplitude, while the p-wave will be assumed to receive only contribution from the last two terms. Noting that $\mathcal{H}_{eff}^{\Delta S=1} \simeq [\mathcal{H}_{eff}]_{RL}^{\Delta S=1}$ the commutator term can be easily evaluated by using the relation:

$$\left[F_i^5, \mathcal{H}_{RL}\right] = -\left[F_i, \mathcal{H}_{RL}\right] \tag{5.10}$$

The K* contribution is easily evaluated and one finds

$$K^{d\,ijk} = \lambda\, f_{7i\ell}\left(i\, f_{\ell kj} + \frac{d'}{f'}\, d_{\ell kj}\right)\left(m_j - m_k\right)\bar{u}(p')\,u(p)\,, \tag{5.11}$$

where λ measures the overall strength of the $K^*-\pi$ transition amplitude and d', f' refer to the coupling of K* with the baryons ($f'+d'=1$). The $K^*-\pi$ transition Lagrangian is chosen to be of $\lambda f_{7ij}\, \partial_\mu P_i V_j$ type since the parity violating amplitudes transform like V_7 in models with R.H.C.C.

The last two terms in (5.9) can be evaluated by assuming SU(3) invariance for the meson-baryon vertex and the generalized Goldberg-Treinman relation:

$$\frac{f_\pi}{\sqrt{2}} = \frac{g_A^{ijk}}{g_{ijk}\, k(0)}\left(m_j + m_k\right) \tag{5.12}$$

The following results are then obtained for the s and p waves:

$$\begin{aligned} A(\Lambda^0_-) &= \frac{-1}{\sqrt{3}\, f_\pi}\left(D + 3F\right) - \frac{\sqrt{3}\,\lambda}{4}\left(\Lambda - N\right)\left(1 + \frac{d'}{3f'}\right) \\ A(\Sigma^+_0) &= \frac{1}{f_\pi}\left(F - D\right) + \frac{\lambda}{4}\left(\Sigma - N\right)\left(1 - \frac{d'}{f'}\right) \\ A(\Sigma^+_+) &= 0 \end{aligned} \tag{5.13}$$

$$A(\Sigma^-_-) = \frac{\sqrt{2}}{f_\pi}(D-F) - \frac{\sqrt{2}\,\lambda}{4}(\Sigma - N)\left(1 - \frac{d'}{f'}\right)$$

$$A(\Xi^-_-) = \frac{1}{\sqrt{3}\,f_\pi}(3F-D) + \frac{\sqrt{3}\,\lambda}{4}(\Xi - \Lambda)\left(1 - \frac{d'}{3f'}\right)$$

$$B(\Lambda^0_-) = \sqrt{2/3}\; g\,(\Lambda + N)\left[\frac{(3F+D)(f+d)}{2N\,(\Lambda - N)} + \frac{2d\,(D-F)}{(\Sigma - N)(\Sigma + \Lambda)}\right]$$

$$B(\Sigma^+_0) = g\,(N+\Sigma)\left[\frac{(F-D)(f+d)}{\sqrt{2}\,N\,(\Sigma - N)} + \frac{2f\,(D-F)}{(\Sigma - N)\,2\Sigma}\right]$$

$$B(\Sigma^+_+) = 2g\,(N+\Sigma)\left[\frac{(F-D)(f+d)}{(\Sigma - N)\,2N} - \frac{(3F+D)d}{3(\Lambda - N)(\Sigma + \Lambda)} - \frac{(F-D)f}{(\Sigma - N)\,2\Sigma}\right]$$

$$B(\Sigma^-_-) = 2g\,(N+\Sigma)\left[\frac{(F-D)f}{(\Sigma - N)\,2\Sigma} - \frac{(3F+D)d}{3(\Lambda - N)(\Sigma + \Lambda)}\right]$$

$$B(\Xi^-_-) = \sqrt{\tfrac{2}{3}}\; g\,(\Xi + \Lambda)\left[\frac{(3F-D)(d-f)}{(\Xi - \Lambda)\,2\Xi} - \frac{2\,(F+D)\,d}{(\Xi - \Sigma)(\Sigma + \Lambda)}\right],$$

where the particle symbols denote the corresponding masses. f and d refer to the meson-baryon coupling constants and F, D refer to the couplings of H_W to the baryons. The best fit to the experimental amplitudes is obtained for:

$$F = 3.3 \times 10^{-5}\ \text{MeV} \;;\quad \frac{d}{f} = 1.78 \;;\quad \frac{d'}{f'} = -.25$$

$$\frac{D}{F} = -.8 \;;\quad \lambda = -9.5 \times 10^{-9}\ \text{MeV}^{-1} \tag{5.14}$$

The agreement with experiment is very good as one can see from the results of the best fit presented in Table 5.1 .

Table 5.1

Best fit solution to s and p wave amplitudes

Decay process	A		B	
	theory	experiment	theory	experiment
Σ^-_-	2.1	1.89±.03	-0.8	-0.72±.01
Σ^+_0	-1.5	-1.53±.14	10.3	11.52±1.85
Σ^+_+	0	0	13.9	19.07±.34
Λ^0_-	1.6	1.53±.02	9.1	10.5 ±.33
Ξ^-_-	-1.8	-2.07±.02	4.4	6.68±.70

One sees that our best fit seems to require a value of λ which is about three times larger than what one usually assumes[34]. Note that λ is essentially an arbitrary parameter whose value can be estimated by assuming that the entire $K_S \to 2\pi$ amplitude arises out of it. In R.H.C.C. models, however, this assumption is no longer valid because the K^* contribution vanishes in the SU(3) limit, whereas $K_S \to 2\pi$ does not vanish in that limit. Thus λ is essentially unconstrained in our case.

5.3 The Possible Role of the Higgs Scalars

Note that in models with the new current $\bar{\chi}_R \gamma_\mu n_R$ the following commutation relations hold:

$$\left[F_i^5, \mathcal{H}_W\right] = -\left[F_i, \mathcal{H}_W^{\Delta I = \frac{1}{2}}\right] + \left[F_i, \mathcal{H}_W^{\Delta I = \frac{3}{2}}\right] \qquad (5.15)$$

Golowich and Holstein[29] have pointed out that these commutation relations are in conflict with the observed $\Delta I=3/2$ effects in $K \to 3\pi$. Their analysis is based on current algebra methods and on the assumption of the validity of a linear expansion in energy for $K \to 3\pi$ amplitudes.

We would like to point out here that since $\Delta I=3/2$ transitions are small, they can be easily affected by other small effects, such as the ones coming from the Higgs part of the Lagrangian. For definiteness, we will work in the framework of a simple SU(2) x U(1) model[4]. In order to give mass to all quarks, one needs the follwing Higgs scalars:

$$\begin{bmatrix} T^+ \\ T_1^0 + i T_2^0 \end{bmatrix} \quad ; \quad \begin{bmatrix} \pi^+ \\ \pi^0 \\ \pi^- \end{bmatrix}$$

The non-vanishing vacuum expectation values are:

$$\langle 0 | T_1^0 | 0 \rangle = k \quad ; \quad \langle 0 | \pi^0 | 0 \rangle = \rho$$

The most general Higgs potential, consistent with renormalizability and gauge invariance is:

$$\mathcal{L}_{Higgs} = \mu_1^2\, T^\dagger T - h_1^2 \left(T^\dagger T\right)^2 + \mu_2^2\, \vec{\pi}\cdot\vec{\pi} - h_2^2 \left(\vec{\pi}\cdot\vec{\pi}\right)^2 + \\ + f\, T^\dagger \vec{\tau}\, T\, \vec{\pi} + f' \left(T^\dagger T\right)\left(\vec{\pi}\cdot\vec{\pi}\right) \tag{5.16}$$

Diagonalizing the mass matrix, one obtains the following two eigenstates for the charged sector:

$$G^{\pm} = \frac{1}{\left(\rho^2 + k^2\right)^{\frac{1}{2}}} \left[k\, T^{\pm} - \rho\, \pi^{\pm} \right]$$

$$H^{\pm} = \frac{1}{\left(\rho^2 + k^2\right)^{\frac{1}{2}}} \left[\rho\, T^{\pm} + k\, \pi^{\pm} \right] \tag{5.17}$$

$G^{\pm}$ are Goldstone bosons and are absorbed as longitudinal components of $W^{\pm}$. $H^{\pm}$ are physical charged scalars whose mass is given by:

$$m^2_{H^{\pm}} = \sqrt{2}\, \frac{f}{\rho} \left(\rho^2 + k^2\right) \tag{5.18}$$

In the neutral sector, one can easily verify that T_2^0 is a Goldstone boson which will be absorbed by the longitudinal component of the neutral vector boson Z. There will be still two physical neutral Higgs bosons which can be expressed as:

$$N_1^0 = \cos\alpha\; t_1^0 + \sin\alpha\; \rho^0$$

$$N_2^0 = -\sin\alpha\; t_1^0 + \cos\alpha\; \rho^0 \tag{5.19}$$

where $\tan\alpha = \dfrac{\lambda_1 - 2h_1^2 k^2}{2f'k\varrho}$, and with t_1° , p° defined by:

$$T_1^\circ = K + t_1^\circ$$
$$\pi^\circ = \varrho + p^\circ$$

The neutral Higgs scalars couple to n,λ quarks, and may contribute to ΔS=2 transitions like:

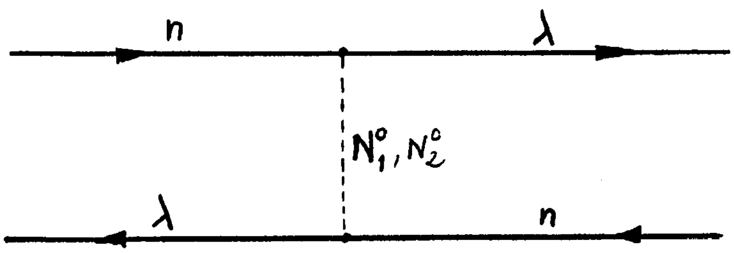


fig 5.1

One can avoid any possible conflict with the experimental size of Δ S=2 transitions either by postulating that the masses of N_1 , N_2 are very high or by eliminating these couplings by an appropriate choice of the parameters of the theory.

We are interested here in the contributions of the physical scalars $H^\pm$, to the effective ΔS=1 weak Lagrangian. If the masses of $H^\pm$ are of order of a few GeV. then the contribution of $H^\pm$ to $\mathcal{L}_{eff}^{\Delta S=1}$. illustrated in fig. 5.2

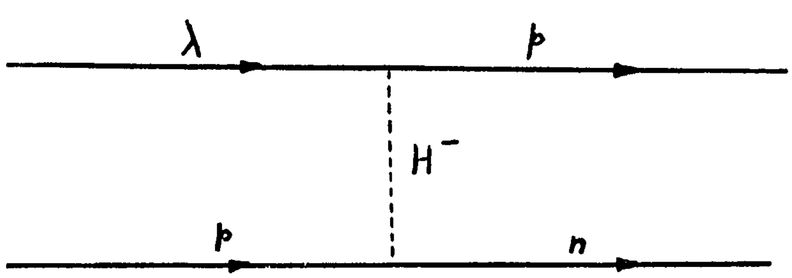


fig 5.2

can be of order

$$\left[\mathcal{L}_{eff}\right]^{\Delta S=1}_{Higgs} \simeq G_F \sin\theta \left[(\bar{n}_L p_R)(\bar{p}_R \lambda_L) + \cdots\right]$$

With the contribution of this new term, commutation relations (5.15) no longer hold and thus, the conclusions of Ref. 29 are no longer valid.

We do not wish to imply that the Higgs scalars $H^{\pm}$ necessarily exist. The point we want to emphasize is that $\Delta I=3/2$ transitions, being a small effect, can be easily affected by the contributions of other interactions, the Higgs scalars contribution being just an example.

Chapter VI

Summary and Conclusions

We have analysed the various implications of the introduction of R.H.C.C. in models of weak interactions. Particular attention was given to the compatibility of the new current with constraints derived from the known low energy weak interactions. We have also investigated the predictions of these models for new phenomena, contrasting them with the standard Weinberg model. In chapter III, it was shown that there is no conflict between the new current $(\bar{\chi}n)_R$ and the smallness of K_L-K_S mass difference. In chapter IV, the hyperon decays were analysed and a good fit was obtained for both s and p waves.

With respect to low energy phenomenology, one is confronted with the following situation:

One class of models (namely those with the Mohapatra charmed current $(\bar{\chi}n)_R$) is able to account forfirst order processes ($\Delta I = \frac{1}{2}$ non-leptonic decays) but fails (in the framework an unjustified approximation) to explain second order processes (namely the smallness of K_L-K_S mass difference).

On the other hand, another theory (the standard model) fails to explain first order phenomena (non-leptonic decays) but can (in an unjustified approximation) account for second order processes. Simple logic leads us to prefer the models with R.H.C.C.

As we have emphasized in chapter IV, the crucial test for models with R.H.C.C., will be the analysis of the various decays of particles carrying net charm. Until these particles are found, it would be of great interest to obtain more accurate measurements for the Σ^+ radiative decays, which at present rule out the standard model and favor models with the current $(\bar{\chi}\, n)_R$.

Appendix I

Evaluation of $\left[\mathcal{L}_{\text{eff}}\right]^{\Delta S=2}_{LR}$

We define the S matrix operator as:

$$S = T \exp\left[i \int d^4x \, \mathcal{L}_I(x)\right] \qquad (A1.1)$$

where

$$\mathcal{L}_I = \frac{g}{\sqrt{2}}\left\{\bar{\lambda}\,\gamma_\mu \frac{1}{2}(1+\gamma_5)\chi \cos\theta + \bar{n}\,\gamma_\mu \frac{1}{2}(1-\gamma_5)\chi\right\} W^\mu + h.c. \qquad (A1.2)$$

The contribution of fig. 2 is then:

$$S_4 = \frac{g^4}{8}\cos^2\theta_c \int d^4x_1\, d^4x_2\, d^4x_3\, d^4x_4\; \bar{\lambda}(x_1)\gamma^\mu\left[\frac{1}{2}(1+\gamma_5)\, i\, S_F(x_1-x_2, m_\chi)\gamma^\alpha \frac{1}{2}(1-\gamma_5)\right] n(x_2)$$

$$\bar{\lambda}(x_4)\gamma^\beta\left[\frac{1}{2}(1+\gamma_5)\, i\, S_F(x_4-x_3, m_\chi)\gamma^\nu \frac{1}{2}(1-\gamma_5)\right] n(x_3) \qquad (A\,1.3)$$

This integral can be easily evaluated in momentum space. Neglecting the external momenta with respect to the internal momentum, one is led to the following integral:

$$\frac{1}{(2\pi)^4}\int d^4k \,\frac{1}{\left(k^2-m_\chi^2+i\epsilon\right)^2\left(k^2-m_W^2+i\epsilon\right)^2} \cong \frac{i}{16\pi^2}\,\frac{1}{m_W^4}\left[\log\frac{m_W^2}{m_\chi^2}-2\right] \qquad (A\,1.4)$$

where we assumed $m_W^2 \gg m_\chi^2$. Going then back to configuration space and using the relation:

$$\gamma^\mu\gamma^\alpha = g^{\mu\alpha} - i\,\sigma^{\mu\alpha} \qquad A(1.5)$$

one finally obtains

$$S_4 \cong \frac{i g^4 \cos^2\theta_c}{128\pi^2} \frac{m_\chi^2}{m_W^2} \ln\left(\frac{M_W^2}{m_\chi^2} - 2\right) I$$

$$\simeq \frac{i G_F}{\sqrt{2}} \frac{\alpha}{4\pi} \left(\frac{m_\chi}{m_W \sin\theta_W}\right)^2 2\cos^2\theta_c \ln\left(\frac{M_W}{m_\chi} - 1\right) I \qquad (A1.6)$$

where

$$I = \int d^4x \left\{ 4\bar{\lambda} \tfrac{1}{2}(1-\gamma_5) n \bar{\lambda} \tfrac{1}{2}(1-\gamma_5) n + \right. \qquad (A1.7)$$

$$\left. + \bar{\lambda} \sigma_{\mu\nu} \tfrac{1}{2}(1-\gamma_5) n \bar{\lambda} \sigma^{\mu\nu} \tfrac{1}{2}(1-\gamma_5) n + h.c. \right\} .$$

AppendixII

We give here the decomposition of $\mathcal{H}_W$ into irreducible representations of SU(4). Let q_α be the basis for the fundamental representation of SU(4) with the identification:

$$q_\alpha \equiv (\chi, p, n, \lambda) \qquad (A\,2.1)$$

One can then construct the second-rank tensor:

$$T_\alpha^\beta = q^\alpha q_\beta - \frac{1}{4}\,\delta_\beta^\alpha\,(q^\gamma q_\gamma) \qquad (A\,2.2)$$

The non-leptonic weak Hamiltonian is given by:

$$\mathcal{H}_W = \frac{G_F}{\sqrt{2}}\left(H_\mu H^{\dagger\mu} + H_\mu^\dagger H_\mu\right) \qquad (A\,2.3)$$

where

$$H_\mu = J_\mu + k_\mu \qquad (A\,2.4)$$

with J_μ given by (4.13). We will consider here the case in which $k_\mu = \bar{\chi}_R \gamma_\mu n_R$

One can then write:

$$H_\mu = A_{\alpha\beta}\, T_{L\alpha}^{\beta} + A'_{\alpha\beta}\, T_{R\alpha}^{\beta} \qquad (A\,2.5)$$

where

$$A = \begin{pmatrix} 0 & C \\ 0 & 0 \end{pmatrix}, \quad \text{with} \quad C = \begin{pmatrix} -\sin\theta_c & \cos\theta_c \\ \cos\theta_c & \sin\theta_c \end{pmatrix}$$

and $\qquad A\,(2.6)$

$$A' = \begin{pmatrix} 0 & M \\ 0 & 0 \end{pmatrix}, \quad \text{with} \quad M = \begin{pmatrix} 1 & 0 \\ 0 & 0 \end{pmatrix}$$

The non-leptonic weak Hamiltonian can then be written:

$$\mathcal{H}_W = \Big[A_{\alpha\beta} A^{\dagger}_{\gamma\delta} \left\{ T^{\beta}_{L\alpha} , T^{\delta}_{L\gamma} \right\} + A'_{\alpha\beta} A'^{\dagger}_{\gamma\delta} \left\{ T^{\beta}_{R\alpha} , T^{\delta}_{R\gamma} \right\} + \qquad (A\,2.7)$$

$$+ A_{\alpha\beta} A'^{\dagger}_{\gamma\delta} \left\{ T^{\beta}_{L\alpha} , T^{\delta}_{R\gamma} \right\} + A'_{\alpha\beta} A^{\dagger}_{\gamma\delta} \left\{ T^{\beta}_{R\alpha} , T^{\delta}_{L\gamma} \right\} \Big]$$

From (A2.5) we see that the current H_μ is a sum of elements of the 15-dimensional representation of SU(4). The symmetric and antisymmetric 15 representation are expressed in terms of T^β_α as:

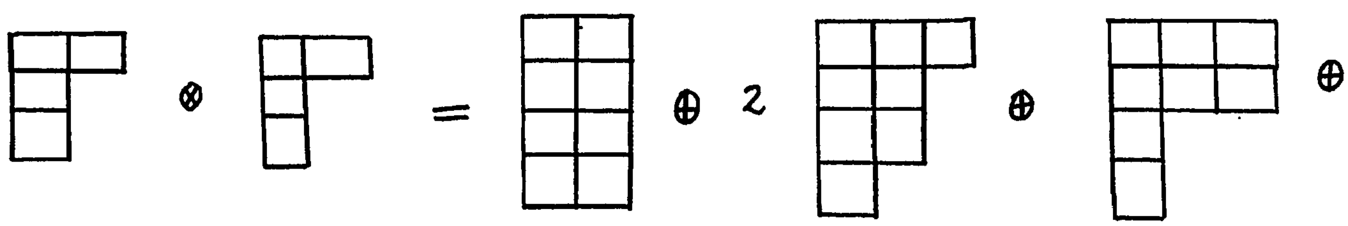
$$[15_A]^{\beta}_{\alpha} = T^{\beta}_{\gamma} T^{\gamma}_{\alpha} - T^{\gamma}_{\alpha} T^{\beta}_{\gamma} \qquad (A\,2.7)$$

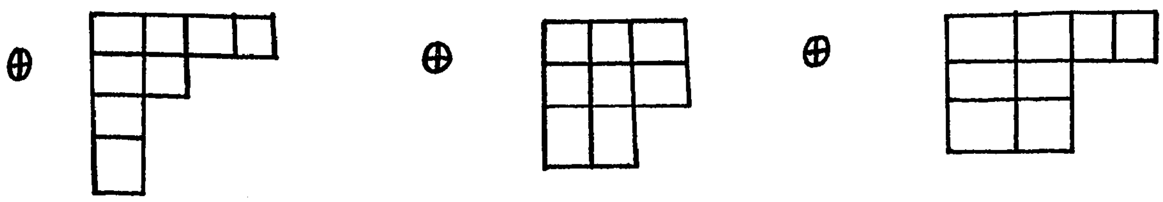
$$[15_S]^{\beta}_{\alpha} = T^{\beta}_{\gamma} T^{\gamma}_{\alpha} + T^{\gamma}_{\alpha} T^{\beta}_{\gamma} - \frac{1}{2} \delta^{\beta}_{\alpha} \left(T^{\gamma}_{\delta} T^{\delta}_{\gamma} \right)$$

In order to obtain the decomposition of $\mathcal{H}_W$, one has thus to evaluate the direct product of $\underline{15} \times \underline{15}$. One obtains in general:

$$\underline{15} \otimes \underline{15} = \underline{1} \oplus \underline{15}_S \oplus \underline{15}_A \oplus \underline{20} \oplus \underline{45} \oplus \underline{45}^* \oplus \underline{84} \qquad (A\,2.8)$$

Or in terms of Young diagrams:





One has now to express $\left\{T_\alpha^\beta , T_\gamma^\delta\right\}$ in terms of irreducible representations of SU(4). Using (A2.2), (A2.8) and taking into account that:

$$\left[\underline{20}\right]^{\beta\gamma}_{\alpha\delta} = T_\alpha^\beta T_\delta^\gamma - T_\alpha^\gamma T_\delta^\beta + T_\delta^\gamma T_\alpha^\beta - T_\delta^\beta T_\alpha^\gamma - \frac{1}{2}\left\{\delta_\delta^\beta \left[\underline{15}_s\right]_\alpha^\gamma - \delta_\delta^\gamma \left[\underline{15}_s\right]_\alpha^\beta - \delta_\alpha^\beta \left[\underline{15}_s\right]_\delta^\gamma + \delta_\alpha^\gamma \left[\underline{15}_s\right]_\delta^\beta \right\} + \frac{1}{6}\left\{\delta_\delta^\beta \delta_\alpha^\gamma - \delta_\delta^\gamma \delta_\alpha^\beta\right\}\left(T_\varepsilon^\lambda T_\lambda^\varepsilon\right) \tag{A2.9}$$

$$\left[\underline{84}\right]^{\beta\gamma}_{\alpha\delta} = \left(T_\alpha^\beta T_\delta^\gamma + T_\alpha^\gamma T_\delta^\beta + T_\delta^\gamma T_\alpha^\beta + T_\delta^\beta T_\alpha^\gamma\right) - \frac{1}{6}\left\{\delta_\alpha^\beta \left[\underline{15}_s\right]_\delta^\gamma + \delta_\delta^\beta \left[\underline{15}_s\right]_\alpha^\gamma + \delta_\delta^\gamma \left[\underline{15}_s\right]_\alpha^\beta + \delta_\alpha^\gamma \left[\underline{15}_s\right]_\delta^\beta \right\} - \frac{1}{10}\left\{\delta_\alpha^\beta \delta_\delta^\gamma + \delta_\delta^\beta \delta_\alpha^\gamma\right\}\left(T_\varepsilon^\lambda T_\lambda^\varepsilon\right)$$

one easily arrives at:

$$\left\{T_\alpha^\beta \; T_\delta^\gamma\right\} = \frac{2}{15}\left\{\delta_\delta^\beta \delta_\alpha^\gamma - \frac{1}{4}\delta_\alpha^\beta \delta_\delta^\gamma\right\}[1] + \frac{1}{2}\left[20\right]^{\beta\gamma}_{\alpha\delta} + \frac{1}{2}\left[84\right]^{\beta\gamma}_{\alpha\delta} + \frac{1}{3}\left\{\delta_\delta^\beta \left[\underline{15}_s\right]_\alpha^\gamma - \frac{1}{2}\delta_\alpha^\beta \left[\underline{15}_s\right]_\delta^\gamma + \delta_\alpha^\gamma \left[\underline{15}_s\right]_\delta^\beta - \frac{1}{2}\delta_\delta^\gamma \left[\underline{15}_s\right]_\alpha^\beta\right\} \tag{A2.10}$$

Finally, the decomposition of $\mathcal{H}_W$ is obtained by substituting (A2.10) in (A2.7). We will first consider the contribution of the $\underline{15}$s since it is in this respect that the models with R.H.C.C. differ drastically from the GIM model. We show next that the $\underline{15}$s is absent in the standard model. From (A.2.10) and (A2.7) one obtains:

$$\frac{1}{3}\left[\underline{15}_s\right]^{\beta}\left[-\frac{1}{2}\left(A_{\alpha\beta}\,\mathrm{tr}\,A^{+}+A^{+}_{\alpha\beta}\,\mathrm{tr}\,A\right)+\left\{A,A^{+}\right\}_{\alpha\beta}\right]. \qquad (A2.11)$$

The first term vanishes, since A is traceless and the second term also vanishes due to the fact that $\left\{A,\ A^{+}\right\}_{\alpha\beta}=\delta_{\alpha\beta}$. Thus one sees that the $\underline{15}$s is not contained in weak non-leptonic Hamiltonian of the standard model. However, when the R.H.C.C. $\bar{\chi}_R\gamma_\mu n_R$ is present this is no longer true and one obtains:

$$\begin{aligned}\mathcal{H}_W = {} & \frac{2}{15}\left[\left[\underline{1}\right]_{LL}+\left[\underline{1}\right]_{RR}-\sin\theta\left[\underline{1}\right]_{LR}-\sin\theta\left[\underline{1}\right]_{RL}\right]+\\ & +\frac{1}{2}\left[A_{\alpha\beta}A^{+}_{\gamma\delta}\left[\underline{20}\right]^{\beta\delta}_{LL\,\alpha\gamma}+A'_{\alpha\beta}A^{+'}_{\gamma\delta}\left[\underline{20}\right]^{\beta\delta}_{RR\,\alpha\gamma}+A_{\alpha\beta}A'_{\gamma\delta}\left[\underline{20}\right]^{\beta\delta}_{LR\alpha\gamma}+\right.\\ & \left.+A'_{\alpha\beta}A^{+}_{\gamma\delta}\left[20\right]^{\beta\delta}_{RL\alpha\gamma}\right]+\left[\text{same with } 20\leftrightarrow 84\right]+\\ & +\frac{1}{3}\left[\left\{A',A'^{+}\right\}_{\alpha\beta}\left[\underline{15}_s\right]^{\beta}_{RR\alpha}+\left\{A',A^{+}\right\}_{\alpha\beta}\left[\underline{15}_s\right]^{\beta}_{RL\alpha}+\left\{A,A'^{+}\right\}_{\alpha\beta}\left[15_s\right]^{\beta}_{LR\alpha}\right]\end{aligned} \qquad (A2.12)$$

The derivation can be repeated for models of class b) with the same final result, apart from an obvious modification in M. In both class of models the contribution of the 15s is non-vanishing, essentialy due to the fact that

$$\left\{ A', A'^{+} \right\} \neq \delta_{\alpha\beta} \, .$$

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$$\frac{f_D}{f_\pi} \simeq 5.95 \qquad \frac{f_F}{f_\pi} \simeq 6.31$$

This would change our predictions drastically and make the decay rates about 30 times larger.

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