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# Study of the *V*, *A* Structure of the Charged Weak Currents of Quarks

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**Abstract.** The $V, A$ structure of charged weak currents of quarks is investigated in the light of the recent results from the high-energy $\overset{(-)}{\nu}_\mu N$ scattering. The bounds for deviations of the quark couplings from the $V-A$ values are determined in three cases: in a general one-boson-exchange parametrization, in the fermion/mirror-fermion mixing parametrization, and in the general left-right symmetric $SU(2)_L \times SU(2)_R \times U(1)$ model. The possibility of a light mirror neutrino is discussed and it is shown that this is not ruled out by the experimental results considered.

## 1. Introduction

The conventional picture of charged weak interactions is based on the $V-A$ structure of fermion currents. In the phenomenological Fermi theory this form of coupling, implying maximal violation of parity, was an *ad hoc* assumption, which in the case of leptons could be justified by assuming that neutrinos are strictly massless and enter the picture with one helicity component, left-handed, only. In the standard $SU(2) \times U(1)$ electroweak model and in the $SU(5)$ grand unified model the $V-A$ structure of charged currents is a basic property of a Lagrangian and is due to the chosen fermion representations under the respective gauge groups. The form of interaction is then not directly connected with any particular property of leptons, e.g. neutrino mass, and is similar for both leptons and quarks.

In all these models charged weak currents are *exactly* $V-A$. Observations do not disagree with this but, on the other hand, large errors of experimental results still admit sizeable modifications of the theory. There exist, indeed, models which may manifest themselves as small deviations from the purely left-handed charged currents. The present authors [1, 2] and some other groups [3, 4] have recently analysed the $V-A$ structure of charged weak couplings of leptons in three models of that sort: a general (factorized) one-boson-exchange parametrization [1, 2], the fermion/mirror-fermion mixing model [2, 3] and a general (non-factorized) left-right symmetric model [2, 4]*. It has been found out, for example, that in the one-boson-exchange parametrization deviations up to a level of 15% from the $V-A$ value of lepton couplings are still allowed.

In the present paper we will extend our study to the quark sector. We analyse the recent high-accuracy charged-current data obtained by the CDHS Collaboration from deep-inelastic $\overset{(-)}{\nu}_\mu N$ scattering [6, 7]. We will work within the frame of the three models mentioned above and make use of our previous results on the muonic current part of the processes. For a detailed description of and motivations for the models we consider, we refer the reader to [2].

The paper is organized as follows: In Sect. 2 we briefly recall the difficulty one meets when trying to determine weak couplings of quarks. We also present our analysis for the general one-boson-exchange parametrization in this section. Sections 3 and 4

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* For a more general analysis of charged leptonic weak interactions including also $S$, $P$, and $T$ couplings, see [5]

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deal with the mixing model and the left-right symmetric model, respectively. We summarize our results in Sect. 5.

## 2. The General One-Boson-Exchange Parametrization

Weak interaction properties of quarks can, in principle, be studied by examining static and dynamic properties of hadrons. In practice, however, these tests are always complicated by the confining colour force experienced by quarks. What one is usually able to do in the case of charged weak currents is just to test exact $V-A$ against exact $V+A$. Surveying the information obtained in various experimental studies it can be concluded (see [8] for a review and references) that the $V-A$ structure is favoured by all charged weak processes of quarks explored so far.

There is, however, a conspicuous scarcity of reliable upper bounds on the possible small deviations from the pure $V-A$ form of quark currents. It is particularly difficult to obtain any such estimates from the static properties of hadrons, since the non-perturbative methods to calculate these properties starting from the basic weak couplings of individual quarks are too inaccurate to be useful. High-energy neutrino-nucleon scattering is expected to be more suitable for this kind of study since there, owing to asymptotic freedom, a perturbative treatment of the strong interactions in the impulse approximation is justified. In the present work we will use the data on the $\overset{(-)}{\nu}_\mu N$ scattering at high energies provided recently by the CDHS Collaboration [6, 7]. There is slight arbitrariness also in the interpretation of these results due, for example, to the non-leading QCD effects and higher twist corrections. However, we are confident of finding results with the right order of magnitude even when neglecting these corrections.

Let us consider the $\overset{(-)}{\nu}_\mu N$ scattering first in the general one-boson-exchange model. We parametrize the charged weak currents of leptons and quarks as follows:

$$\mathscr{J}_\ell^\alpha = \bar{\ell}\gamma^\alpha (V_\ell - A_\ell \gamma_5)\nu_\ell, \qquad \mathscr{J}_q^\alpha = \bar{q}'\gamma^\alpha (V_q - A_q\gamma_5) q. \tag{1}$$

The differential cross-sections for $\overset{(-)}{\nu}_\mu q$ scatterings are then given by

$$\frac{d^2\sigma^{\nu_\mu q}}{dx\,dy} = Q(x)[C_1 + P_{\nu_\mu} C_4 + (1-y)^2 (C_2 + P_{\nu_\mu} C_3)],$$
$$\frac{d^2\sigma^{\bar{\nu}_\mu q}}{dx\,dy} = Q(x)[C_2 - P_{\bar{\nu}_\mu} C_3 + (1-y)^2 (C_1 - P_{\bar{\nu}_\mu} C_4)], \tag{2}$$

where $x$ and $y$ are the standard scaling variables, $Q(x)$ stands for kinematical factors and respective quark structure functions, and the constants $C_i$ have the following expressions ($CP$ invariance assumed)

$$C_{1(2)} = (V_\mu^2 + A_\mu^2)(V_q^2 + A_q^2) \underset{(-)}{+} 4 V_\mu A_\mu V_q A_q,$$
$$C_{3(4)} = \underset{(-)}{+} 2 V_q A_q (V_\mu^2 + A_\mu^2) - 2 V_\mu A_\mu (V_q^2 + A_q^2). \tag{3}$$

We have assumed here that muon neutrinos are Dirac particles. Since also the production vertex of neutrinos has the general structure of the form given in (1), the longitudinal polarization of the beam is non-trivial:

$$P_{\nu_\mu} = -P_{\bar{\nu}_\mu} = -\frac{2 V_\mu A_\mu}{V_\mu^2 + A_\mu^2}. \tag{4}$$

The CDHS Collaboration has used the $y$-distributions, which they found to be flat in accordance with left-handed currents and quark dominance, to determine upper bounds on the cross-sections associated with right-handed currents. In the case of dimuon production they give the limit [7]

$$\rho_{2\mu}^2 \equiv \frac{\sigma_R^{2\mu}}{\sigma_L^{2\mu}} < 0.07 \qquad (95\,\%\ \text{C.L.}). \tag{5}$$

For one-muon production their result is [6]

$$\rho_{1\mu}^2 \equiv \frac{\sigma_R^{1\mu}}{\sigma_L^{1\mu}} = 0.000 \pm 0.005 \qquad (<0.009,\ 90\,\%\ \text{C.L.}). \tag{6}$$

The corresponding theoretical quantity in the case of the general one-boson-exchange model is obtained from (2):

$$\rho_{\text{th}}^2(\lambda_\mu, \lambda_q) = \frac{C_2 + P_{\nu_\mu} C_3}{C_1 + P_{\nu_\mu} C_4} = \frac{(1+\lambda_\mu^2)^2(1+\lambda_q^2) - 8\lambda_\mu\lambda_q(1+\lambda_\mu^2) + 4\lambda_\mu^2(1+\lambda_q^2)}{(1+\lambda_\mu^2)^2(1+\lambda_q^2) + 8\lambda_\mu\lambda_q(1+\lambda_\mu^2) + 4\lambda_\mu^2(1+\lambda_q^2)}, \tag{7}$$

where we have introduced the notation

$$\lambda_i = \frac{A_i}{V_i}. \tag{8}$$

The experimental constraints given in (5) and (6) determine a domain of allowed values in the $(\lambda_\mu, \lambda_q)$ plane. However, we can make better use of the data if we take into account the bounds we have got earlier from other experimental constraints for the muonic current [2] (we give all our results at the 68 % confidence level):

$$|\lambda_\mu - 1| < 0.11. \tag{9}$$

It turns out in fact that bounds on the quark-current parameter $\lambda_q$ are, to the precision used, independent of $\lambda_\mu$ within this range.

The dimuon production can be understood to arise from the excitation and subsequent decay of a charm quark. Measuring the $y$-distributions of these processes thus provides information about the structure of $cq$ currents, $q$ being mostly $d$ and $s$. More specifically, one can turn the result given in [7] into limits on the value of the parameter $\lambda_c$:

$$0.69 \leqq \lambda_c \leqq 1.46. \tag{10}$$

The more accurate one-muon data measure a weak charged current which dominantly is the $ud$ current. In this case the data [6] are far more restrictive and lead to the following limits for the parameter $\lambda_q$ ($\simeq \lambda_u$):

$$0.868 \leqq \lambda_q \leqq 1.152. \tag{11}$$

We see that according to this one experiment the charged weak current of light quarks is limited to being as close to the $V-A$ value as what we have found for the leptonic currents ($|\lambda_e - 1| < 0.15$, $|\lambda_\mu - 1| < 0.11$, see [2]).

We want to make a few remarks here. Because of the non-trivial form (4) of the longitudinal neutrino polarization, the expression (7) for $\rho_{\text{th}}^2$ is asymmetric with respect to $\lambda_\mu$ and $\lambda_q$. In particular, $\rho_{\text{th}}^2$ is much less sensitive to the small deviations of $\lambda_\mu$ from its $V-A$ value than to those of $\lambda_q$. This can easily be seen by setting one of the two equal to unity:

$$\begin{aligned} \rho_{\text{th}}^2(1, \lambda_q) &= \left(\frac{1-\lambda_q}{1+\lambda_q}\right)^2, \\ \rho_{\text{th}}^2(\lambda_\mu, 1) &= \left(\frac{1-\lambda_\mu}{1+\lambda_\mu}\right)^4. \end{aligned} \tag{12}$$

If we assumed that the quark current is purely $V-A$, i.e. $\lambda_q = 1$, the one-muon data would give the limits

$$0.580 \leqq \lambda_\mu \leqq 1.724 \tag{13}$$

for the muon coupling. These bounds are much less restrictive than those we have obtained from other experiments, given in (9), and the limits we obtained for $\lambda_q$, given in (11). On the other hand, adopting a "universality" relation $\lambda_q = \lambda_\mu$ would yield the same bounds as given for $\lambda_q$ in (11).

In the next sections we will consider the two other models mentioned in the Introduction. We shall limit our analysis there to the one-muon data.

## 3. The Fermion/Mirror-Fermion Mixing Model

The fermion/mirror-fermion mixing model [2–4] is based on the gauge structure of the $SU(2) \times U(1)$ electroweak model* but differs from the standard model in having the particle content extended by mirror fermions. Mirror fermions are like ordinary fermions except that they have the opposite chirality (the mirror weak eigenstates couple to the charged gauge boson in $V+A$ currents) and different masses, of course. What makes this model interesting is that ordinary fermions and mirror fermions may be mixed by mass terms. Mass eigenstates will then have charged weak couplings** which are not purely chiral, i.e. deviations from the exact $V-A$ form of currents may appear. Since the possibility of light mirror neutrinos is not ruled out, we will consider two cases separately: the muonic mirror neutrino $N_\mu$ is heavy, $m_{N_\mu} > m_K$ (models b and c of [2]), and $N_\mu$ is light, $m_{N_\mu} \ll m_e$ (models a and d).

### *Case 1: $N_\mu$ Heavy*

The results of Sect. 2 are directly applicable in this case since the only effect of mirror fermions is to shift the couplings of ordinary leptons and quarks from the pure $V-A$ to the forms given in (1). Instead of the parameters $V_i$ and $A_i$ it is now, however, more instructive to use mixing angles [2, 3] related to $V_i$ and $A_i$ as follows:

$$\begin{aligned} V_\mu &= \cos(\Theta_\mu - \Phi_\mu), \\ A_\mu &= \cos(\Theta_\mu + \Phi_\mu), \\ V_q &= \cos(\Theta_d - \Theta_u), \\ A_q &= \cos(\Theta_d + \Theta_u), \end{aligned} \tag{14}$$

where $\Theta_\mu$, $\Phi_\mu$, $\Theta_u$, and $\Theta_d$ are, respectively, mixing angles of muon, muon neutrino, $u$ quark and $d$ quark, with the corresponding mirror particles (the $ud$ current dominance is assumed here explicitly). The range for $\lambda_q$ which we have given in (11) corresponds now to a domain in the $(\Theta_u, \Theta_d)$ plane. This domain is symmetrical in the two axes; the quadrant $\Theta_u, \Theta_d \geqq 0$ is depicted in Fig. 1. We present in the same figure also the slightly tighter boundary for $\Theta_\mu$ and $\Phi_\mu$ obtained from the limit in (9).

We notice two facts. Firstly, if one of the angles $\Theta_d$, $\Theta_u$ (similarly $\Theta_\mu$ and $\Phi_\mu$) is zero, the other one is unrestricted and the parameter $\lambda_q$ ($\lambda_\mu$) is equal to one, trivially. This multivaluedness of the $V-A$ limit, obvious on the grounds of the relations (14), has

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* Some grand unified models in which mirror fermions appear have $SU(2)_L \times SU(2)_R \times U(1)$ electroweak symmetry, and have therefore two sources for $V+A$ couplings. For simplicity, we will not consider this possibility here

** The neutral current couplings in the mixing model are discussed in [9]

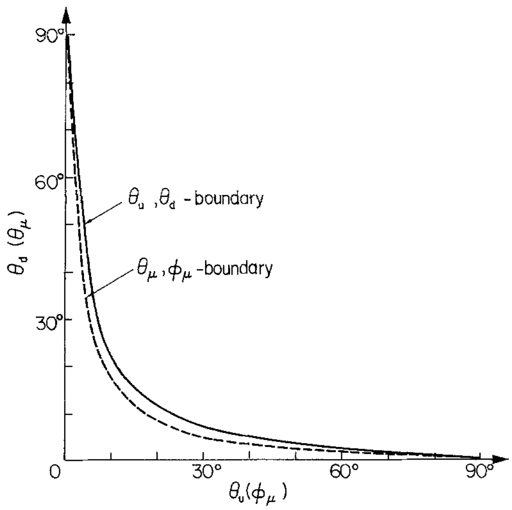


**Fig. 1.** Boundaries giving the upper limits for the fermion/mirror-fermion mixing angles (68%C.L.). The boundaries are symmetric with respect to the two axes, only one quadrant being presented in the figure

been discussed in connection with other processes in [2] (model c) and in [3]. Secondly, imposing the condition $\Theta_d = \Theta_u$ we find $|\Theta_d| = |\Theta_u| \leqq 15°$. This is very close to the corresponding limit on the leptonic side: if $\Theta_\mu = \Phi_\mu$, then $|\Theta_u| = |\Phi_u| \leqq 13.6°$.

*Case 2: $N_\mu$ Light*

This case is more complicated, in principle, since neutrino beams now consist of two components, $\nu_\mu$ and $N_\mu$, which are coherently produced in pion and kaon decay. The probability of finding either of these neutrinos interacting is an oscillatory function depending, for example, on the squared mass difference $\Delta m^2 = |m_{N_\mu}^2 - m_{\nu_\mu}^2|$. However, we will see below that the quantity $\rho_{\text{th}}^2$ does not depend on this oscillation and consequently there is no arbitrariness due to the unknown neutrino and mirror-neutrino masses. Furthermore, it turns out that bounds on the mixing angles $\Theta_u$ and $\Theta_d$ are, to the precision given, the same as in Case 1.

The leptonic charged weak current written in terms of weak eigenstates $\nu'_\mu$ and $N'_\mu$ reads as follows:

$$\mathcal{J}_\ell^\alpha = \cos\Theta_\mu \bar{\mu}\gamma^\alpha(1-\gamma_5)\nu'_\mu - \sin\Theta_\mu \bar{\mu}\gamma^\alpha(1+\gamma_5)N'_\mu. \quad (15)$$

The mass eigenstate neutrinos are obtained from the weak eigenstates by a rotation defined by:

$$\begin{pmatrix} \nu_\mu \\ N_\mu \end{pmatrix} = \begin{pmatrix} \cos\Phi_\mu & -\sin\Phi_\mu \\ \sin\Phi_\mu & \cos\Phi_\mu \end{pmatrix} \begin{pmatrix} \nu'_\mu \\ N'_\mu \end{pmatrix}. \quad (16)$$

The standard formulae for neutrino oscillation (see, for example, [10]) give us the following expressions for the different $\nu'_\mu$, $N'_\mu$ oscillation probabilities

$$\begin{aligned} P_{\nu'_\mu\nu'_\mu}(x) &= P_{N'_\mu N'_\mu}(x) = 1 - \tfrac{1}{2}\sin^2 2\Phi_\mu\left(1-\cos\frac{2\pi x}{L}\right), \\ P_{\nu'_\mu N'_\mu}(x) &= P_{N'_\mu\nu'_\mu}(x) = \tfrac{1}{2}\sin^2 2\Phi_\mu\left(1-\cos\frac{2\pi x}{L}\right), \end{aligned} \quad (17)$$

where $x$ is the distance between production and interaction points and the oscillation length $L$ is given by

$$L = \frac{4\pi E_\nu}{\Delta m^2}. \quad (18)$$

One can see that if $L$ is large compared to $x$, i.e. if the mass difference $\Delta m^2$ is small compared to the average beam energy $E_\nu$, the oscillation between $\nu'_\mu$ and $N'_\mu$ is negligible and $\Phi_\mu$ dependence will disappear. In the opposite extreme case of a fairly large $\Delta m^2$ (although with $m_{\nu_\mu}$, $m_{N_\mu} \ll m_e$) the oscillating term averages to zero and the probabilities depend only on $\Phi_\mu$ but not on $\Delta m^2$.

The $\Delta m^2$ is, however, unknown and we must hence consider the general case. Weighting the beam, according to (15), with $\cos^2\Theta_\mu$ and $\sin^2\Theta_\mu$ for $\nu'_\mu$ and $N'_\mu$, respectively, and taking oscillations into account, we find

$$\begin{aligned} &C_{2(1)} + P_{\nu_\mu} C_{3(4)} \\ &= (V_q^2 + A_q^2)(P_{\nu'_\mu\nu'_\mu}(x)\cos^4\Theta_\mu + P_{N'_\mu N'_\mu}(x)\sin^4\Theta_\mu) \\ &\overset{+}{(-)} 2V_q A_q (P_{\nu'_\mu\nu'_\mu}(x)\cos^4\Theta_\mu - P_{N'_\mu N'_\mu}(x)\sin^4\Theta_\mu). \end{aligned} \quad (19)$$

[Note that the non-diagonal oscillation probabilities $P_{\nu'_\mu N'_\mu}(x)$ and $P_{N'_\mu\nu'_\mu}(x)$ do not appear in this formula since they correspond to situations when the states $\nu'_{\mu L}$ and $N'_{\mu R}$ oscillate to the sterile states $N'_{\mu L}$ and $\nu'_{\mu R}$, respectively.] Using then (17) one can see that dependence on the oscillation drops out of the $\rho_{\text{th}}^2$:

$$\rho_{\text{th}}^2 = \frac{1 - \dfrac{2\lambda_q}{1+\lambda_q^2}\cdot\dfrac{1-\tan^4\Theta_\mu}{1+\tan^4\Theta_\mu}}{1 + \dfrac{2\lambda_q}{1+\lambda_q^2}\cdot\dfrac{1-\tan^4\Theta_\mu}{1+\tan^4\Theta_\mu}}. \quad (20)$$

The quantity measured in the CDHS experiment hence does not depend either on the angle $\Phi_\mu$ or on the squared mass difference $\Delta m^2$.* Furthermore, it is very insensitive to changes in values of the angle $\Theta_\mu$ close to zero. The angle $\Theta_\mu$ is strictly bounded by other data [2], $|\Theta_\mu| \leqq 3.2°$, and, in fact, the value of

* Information about these quantities could be obtained, for example, in the neutrino oscillation experiments

$\rho^2_{\rm th}$ to the precision we use is independent of its variations within this range. We therefore finish up with the same numerical limits on $\lambda_q$ and the same $(\Theta_u, \Theta_d)$ boundary as in Case 1.

## 4. The Left-Right Symmetric Model

There are two parameters in the $SU(2)_L \times SU(2)_R \times U(1)$ left-right symmetric model that are relevant to the charged-current processes [2, 4]: the squared ratio of charged vector boson masses, $r=(m_{W_1}/m_{W_2})^2$, and the angle $\omega$ that measures the possible mixing between the weak eigenstate gauge bosons $W_L$ and $W_R$. In the case of no mixing $W_1 = W_L$ and $W_2 = W_R$.

The differential cross-sections for $\overset{(-)}{\nu}_\mu q$ scatterings are the following

$$\frac{d\sigma^{\nu_\mu q}}{dx\,dy} = Q(x)\{(1+\eta_{AA})^2 + 4\eta_{AV}^2 + 4P_{\nu_\mu}\eta_{AV}(1+\eta_{AA}) + (1-y)^2(1-\eta_{AA})^2\},$$

$$\frac{d\sigma^{\bar\nu_\mu q}}{dx\,dy} = Q(x)\{(1-\eta_{AA})^2 + (1-y)^2[(1+\eta_{AA})^2 + 4\eta_{AV}^2 - 4P_{\bar\nu_\mu}\eta_{AV}(1+\eta_{AA})^2]\}, \tag{21}$$

where the longitudinal polarizations $P_{\overset{(-)}{\nu}_\mu}$ are given by

$$P_{\nu_\mu} = -P_{\bar\nu_\mu} = \frac{2\eta_{AV}\eta_{AA}}{\eta_{AA}^2 + \eta_{AV}^2}, \tag{22}$$

and in the zero-momentum transfer limit we have (notations used here are introduced in [4])

$$\eta_{AV} = \frac{(r-1)\varepsilon}{1+r\varepsilon^2}, \quad \eta_{AA} = \frac{r+\varepsilon^2}{1+r\varepsilon^2}, \tag{23}$$

where

$$\varepsilon = \frac{1+\tan\omega}{1-\tan\omega}.$$

It is assumed here that neutrinos are Dirac particles.

The expression for $\rho^2_{\rm th}$ is now given by

$$\rho^2_{\rm th} = \frac{(1-\eta_{AA})^2}{(1+\eta_{AA})^2 + 4\eta_{AV}^2 + 4P_{\nu_\mu}\eta_{AV}(1+\eta_{AA})}. \tag{24}$$

The experimental constraint in (6) determines an area in the $(r, \omega)$ plane. This is presented in Fig. 2.

If neutrinos are Majorana particles with heavy right-handed components, the leptonic part of the charged right-handed currents is completely suppressed at low energies since the right-handed neutrino could not be produced. In this case, as noted in [11], one cannot obtain any information about the parameters of the model, for example from muon or pion decay. The data on $\rho^2$, however, do restrict $r$ and $\omega$ even in this case. The expression for $\rho^2_{\rm th}$ in the Majorana case is obtained from (24) simply by setting $P_{\nu_\mu} = -1$, since the beam now contains only left-handed neutrinos. The boundary in this case, depicted in Fig. 2, is very close to the boundary obtained for the Dirac case. The reason is that the polarization $P_{\nu_\mu}$, as given in (24), very weakly depends on small deviations of the parameters from their $V-A$ values:

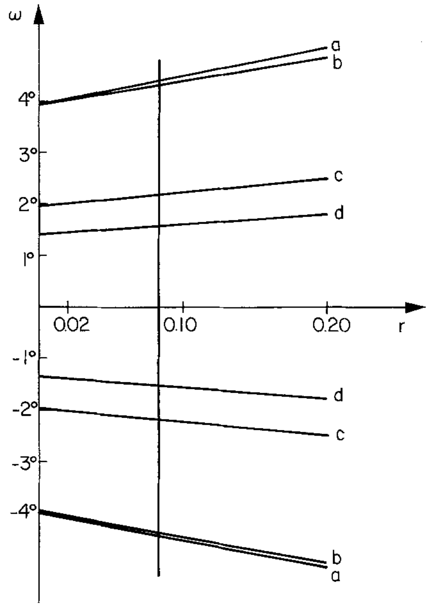


**Fig. 2.** Bounds obtained for the parameters of the left-right symmetric model. a) $\overset{(-)}{\nu}_\mu N$ data, the muon neutrino is a Majorana particle and $\nu_{\mu R}$ is heavy; b) $\overset{(-)}{\nu}_\mu N$ data, the muon neutrino is a Dirac particle; c) muon decay data, $\nu_{eR}$ light and $\nu_{\mu R}$ heavy (or vice-versa); d) muon decay data, $\nu_e$ and $\nu_\mu$ are Dirac particles. Allowed region (68 % C.L.) in each case is between the respective pair of lines. The vertical line is the upper limit for the parameter $r$, $r \leq 0.084$ (68 % C.L.), as given in [2]

$$P_{\nu_\mu} \simeq -1 + 2(r+\omega)^2. \tag{25}$$

As can be seen from Fig. 2, $\rho^2$ restricts essentially only the mixing angle $\omega$. If electron neutrino and muon neutrino are not both Majorana particles with heavy right-handed component, more precise data on this angle come, however, from the Michel parameter $\rho_\mu$ of muon decay. We will present, for comparison, also the boundaries which correspond to the experimental limits on this quantity [12]:

$$\rho_\mu = 0.7517 \pm 0.0026. \tag{26}$$

If neutrinos $\nu_e$ and $\nu_\mu$ are Dirac particles, or, what is effectively equivalent, Majorana particles with light

($\ll m_e$) right-handed components, the Michel parameter is given by [4]

$$\rho_\mu = \frac{3}{8} \frac{(1+\eta_{AA})^2 + 4\eta_{AV}^2}{1+\eta_{AA}^2 + 2\eta_{AV}^2}. \tag{27}$$

If both of the neutrinos are "true" Majorana particles, i.e. their right-handed components are very heavy, $\rho_\mu$ is trivially equal to 3/4 and was pointed out above, no information about the parameters is obtained. In the odd case that one of the neutrinos is a Dirac particle (or a Majorana particle with a light right-handed component) and the other one is a "true" Majorana particle we find [$\rho_\mu$ is symmetric with respect to the exchange of the rôles of the two neutrinos*]

$$\rho_\mu = \frac{3}{4} \frac{1}{1+\eta_M}, \tag{28}$$

where

$$\eta_M \equiv \frac{(1-\varepsilon^2)(1-r)}{(1+\varepsilon)^2 + r(1-\varepsilon)^2}.$$

Both Equations (27) and (28) give limits on the $(r, \omega)$ plane which are much tighter than those obtained from the quantity $\rho^2$.

## 5. Summary

To summarize:

i) The recent data [6, 7] on the high-energy $\overset{(-)}{\nu}_\mu N$ scattering imply a stringent upper limit on the deviation of the charged weak current of light quarks from the exact $V-A$ form. Expressed in terms of a general one-boson-exchange parametrization:

$$0.868 \leqq \lambda_q \equiv \frac{A_q}{V_q} \leqq 1.152 \qquad (68\%\,\text{C.L.}).$$

This puts the status of quark currents about on the same level of accuracy as that of leptons [2].

ii) Applied to the fermion/mirror-fermion mixing model parametrization the experimental limits determine a domain in the $(\Theta_u, \Theta_d)$ plane depicted in Fig. 1. The result is not dependent on the unknown mass $m_{N_\mu}$ of the muonic mirror neutrino, if $m_{N_\mu} > m_K$ or $m_{N_\mu} \ll m_e$.

iii) In the left-right symmetric model the experimental bounds limit the mixing angle to the values $|\omega| \lesssim 4°$ (0.07 rad) (68 % C.L.) (Fig. 2). This result is independent of whether the muon neutrino is a Dirac or Majorana particle. The more stringent limits on $\omega$ from muon decay are also presented.

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* The other muon decay parameters, for example $\delta$ and $\xi$, are not symmetric and will therefore distinguish the two possibilities