TRI-PP 88-84
9

TRI-PP-88-84
Sep 1988

Light Neutrino Masses in the Fritzsch Model with Horizontal Peccei-Quinn Symmetry

C.Q. Geng and John N. Ng
TRIUMF Theory Group, 4004 Wesbrook Mall, Vancouver, B.C., Canada V6T 2A3

Abstract

We study the pattern of light neutrino masses in a model of Fritzsch-type quarks and lepton mass matrices endowed with a global horizontal Peccei-Quinn symmetry. The model places the masses of $(\nu_e,\ \nu_\mu,\ \nu_\tau)$ to be in the range $(10^{-7},\ 10^{-4},\ 10)$ eV and $(10^{-2},\ 10,\ 10^{6})$ eV for stable and unstable ν_τ cases respectively.

(sumitted to Physical Review D)

1 Introduction

One of the unsatisfactory features of the standard model is the lack of understanding of the strong CP-problem.[1] The parameter θ which multiplies a $P,\ T$ violating term in the QCD Lagrangian has to be fine tuned to $< 10^{-9}$ in order to agree with non-observation of the neutron dipole moment. An elegant solution of this problem is to add a global Peccei-Quinn (PQ) $U(1)_{PQ}$ symmetry[2] which breaks at a scale f_a in the order of $10^8 - 10^{12}$ GeV according to the experimental, astrophysical, and cosmological considerations.[3] This high-scale breakdown of $U(1)_{PQ}$ gives rise to an invisible pseudo-Goldstone boson, the invisible axion.[4-6] There are two well-known minimal types of invisible axion models: (i) Dine-Fischler-Srednicky-Zhitnitskii (DFSZ)[5] and (ii) Kim-Shifman-Vainshtein-Zakharov (KSVZ)[6] types, in which the breaking scale f_a is obtained by the large vacuum expectation value (VEV) of a complex Higgs-singlet χ field.

If neutrinos are not massless, one can achieve an understanding of the smallness of their masses within the framework of invisible axion models. This is done by introducing one singlet right-handed neutrino for each family. Then by identifying f_a with the scale Λ of the see-saw mechanism[7] the following naive neutrino mass hierarchy

$$M_{\nu_e} : M_{\nu_\mu} : M_{\nu_\tau} \sim m^2_{u,e} : m^2_{c,\mu} : m^2_{t,\tau}\ , \qquad (1.1)$$

can be derived.[8] It is common to assume a generation independent diagonal Dirac neutrino mass matrix and a unique right-handed Majorana mass for each generation.

In general, the neutrino mass hierarchy (NMH) of (1.1) need not hold true and it can be different for different structures of fermion mass matrices. However, in practice, it is hard to get exact expression of NMH because of the difficulty in diagonalizing the complicated 6×6 neutrino mass matrix in general. Naturally, one hopes to obtain a simple matrix coming from the model itself. The first goal in constructing such models, to achieve realistic charged fermion mass matrices which are constrained by the experimental informations on their mass patterns and the Kobayashi-Maskawa (KM) matrix of quark mixings.[9] It is well-known that the most promising quark mass matrices are the Fritzsch-type matrices[10] which lead to the KM matrix. They are based on the observed quark masses. A way to derive the Fritzsch quark matrices is to impose a global horizontal family symmetry on the model. It should be very interesting if this symmetry can be identified as a PQ symmetry and thus relating the fermion mass matrices and the strong CP-problem simultaneously. Obviously, the PQ charges of the symmetry, denoted by $U(1)^H_{PQ}$, are no longer generation-independent. To incorporate this we have to go beyond the DFSZ and KSVZ models. It has been pointed out by Davidson and Wali (DW)[11] that the two Higgs doublets and one Higgs singlet minimal model with Fritzsch-type quark mass matrices and a global horizontal-PQ symmetry exists and these charges are uniquely determined. In their model, flavor changing neutral current induced by the pseudo-Goldstone boson called familon-axion are suppressed if $U(1)^H_{PQ}$ breaking scale is large. In this paper, we will study in detail the DW model. Especially, we will explore the light neutrino masses in various cases.

The paper is organized as follows. In Sect. 2, we present details of the model and study the fermion mass matrices. In Sect. 3, we examine the light neutrino masses with the experimental and cosmological bounds imposed on them. The concluding remarks are given in Sect. 4.

2 The Model

The model contains two doublets $\phi_i (i = 1, 2)$ and one singlet χ Higgs fields with $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)^H_{PQ}$ symmetry where $U(1)^H_{PQ}$ is a global horizontal PQ symmetry. The Higgs and fermion fields have the following $U(1)^H_{PQ}$ transformations:

$$\phi_1 \to e^{-i\alpha}\,\phi_1\,,\quad \phi_2 \to e^{i\alpha}\,\phi_2\,,\quad \chi \to e^{-i\alpha}\,\chi$$
$$F^i_L = (q, u^c, d^c, l, \nu^c, e^c)^i_L \to e^{iZ_i\alpha}\,F^i_L\,, \tag{2.1}$$

where $q^i_L = \binom{u}{d}^i_L$ and $l^i_L = \binom{\nu}{e}^i_L$ are the generic left-handed quark and lepton doublets respectively and $i = 1, 2, 3$ are the family indices. The fermionic PQ charges are denoted by Z_i.

We now determine the horizontal PQ charge Z_i of each fermion family by requiring that the up- and down-quark mass matrices be the Fritzsch-type matrix. The PQ charge for the quarks can be written in a matrix form

$$Z = (Z_{ij} \equiv Z_i + Z_j)\,, \tag{2.2}$$

explicitly

$$Z = \begin{pmatrix} 2Z_1 & Z_1 + Z_2 & Z_1 + Z_3 \\ Z_1 + Z_2 & 2Z_2 & Z_2 + Z_3 \\ Z_1 + Z_3 & Z_2 + Z_3 & 2Z_3 \end{pmatrix}\,. \tag{2.3}$$

In order to have a Fritzsch-type matrix:[10]

$$F = \begin{pmatrix} 0 & F_{12} & 0 \\ F_{21} & 0 & F_{23} \\ 0 & F_{32} & F_{33} \end{pmatrix}\,, \tag{2.4}$$

for quark mass matrices one must satisfy the following conditions:

$$|Z_1 + Z_2| = |Z_2 + Z_3| = 2|Z_3| = 1\,, \tag{2.5}$$

and

$$2|Z_1|,\ 2|Z_2|,\ |Z_1+Z_3| \neq 1 \ , \tag{2.6}$$

because the $U(1)^H_{PQ}$charges for $\phi_{1,2}$ are ± 1. From (2.5) and (2.6), one finds that the only solutions are

$$(Z_1, Z_2, Z_3) = \pm\left(\frac{5}{2}, -\frac{3}{2}, \frac{1}{2}\right) \ . \tag{2.7}$$

The two solutions become equivalent when the signs of the ϕ_i's charges are re-defined and we shall take the plus sign in (2.7) for the remaining paper. The PQ charges in (2.7) are precisely the one given by DW.[11] It should be noted that the PQ charges for ϕ_1 and ϕ_2 must be opposite in order to have both up- and down-quark mass matrices be Fritzsch-type. It is easily seen that the assignment of PQ charges in (2.7) results in the same amount of color anomaly on $[SU(3)_c]^2 \cdot U(1)^H_{PQ}$ as in the DFSZ model.[1,5]

We thus have the following Yukawa couplings:

$$\begin{aligned}\mathcal{L}_Y \ = \ & \left(h^u_{12}\bar{q}^1_L\tilde{\phi}_1 u^2_R + h^u_{21}\bar{q}^2_L\tilde{\phi}_1 u^1_R + h^u_{23}\bar{q}^2_L\tilde{\phi}_2 u^3_R + h^u_{32}\bar{q}^3_L\tilde{\phi}_2 u^2_R + \right.\\ & \left. + h^u_{33}\bar{q}^3_L\tilde{\phi}_1 u^3_R\right) + \left(u \to d,\ \tilde{\phi}_{1(2)} \to \phi_{2(1)}\right) + (q \to l,\ u \to \nu,\ d \to e) + \\ & + \left(h^N_{12}\nu^{1T}_R C\nu^2_R\chi + h^N_{21}\nu^{2T}_R C\nu^1_R\chi + h^N_{23}\nu^{2T}_R C\nu^3_R\chi^+ + \right.\\ & \left. + h^N_{32}\nu^{3T}_R C\nu^2_R\chi^+ + h^N_{33}\nu^{3T}_R C\nu^3_R\chi\right) + \text{h.c.} \ ,\end{aligned} \tag{2.8}$$

and the most general Higgs potential:

$$\begin{aligned} V \ = \ & \sum_i m^2_i\phi^+_i\phi_i + \sum_{i,j} a_{ij}(\phi^+_i\phi_i)(\phi^+_j\phi_j) + b_{12}(\phi^+_1\phi_2)(\phi^+_2\phi_1) + \\ & [c\phi^+_1\phi_2\chi^2 + \text{h.c.}] \ . \end{aligned} \tag{2.9}$$

We write the Higgs fields and their VeV's as follows:

$$\phi_i = \begin{bmatrix} \phi^+_i \\ \frac{1}{\sqrt{2}} e^{i\theta_i}(v_i + R_i + iI_i) \end{bmatrix} \ , \quad \chi = \frac{1}{\sqrt{2}} e^{i\theta_3}(\Lambda + R_3 + iI_3) \ , \tag{2.10}$$

and

$$\langle\phi_i\rangle = \frac{1}{\sqrt{2}} e^{i\theta_i} v_i \ , \quad \langle\chi\rangle = \frac{1}{\sqrt{2}} e^{i\theta_3}\Lambda \ , \tag{2.11}$$

where $v = \sqrt{v_1^2 + v_2^2} = \left(\sqrt{2}\,G_F\right)^{-1/2} \approx 250$ GeV. Spontaneous symmetry breaking takes place at two different scales based on the VEV's of the doublets and singlet Higgs field. The Higgs potential V shown in (2.9) is exactly the same as the DFSZ model.[5] From (2.9)–(2.11), one finds that the axion field is

$$a = \frac{1}{f_a}\left[\frac{2v_1v_2}{v^2}(v_2I_1 - v_1I_2) + \Lambda I_3\right] \ , \tag{2.12}$$

with the axion decay constant

$$f_a = \frac{1}{2v}[4v_1^2v_2^2 + v^2\Lambda^2]^{1/2} \approx \frac{\Lambda}{2} \ , \tag{2.13}$$

for $\Lambda \gg v$. In order to obtain an invisible familon-axion, f_a would have to be in the range of $10^8 - 10^{12}$ GeV. The axion acquires a small mass via the color anomaly, and is given by

$$m_a = m_\pi \frac{f_\pi}{f_a}\frac{3\sqrt{m_u m_d}}{(m_u + m_d)} \ . \tag{2.14}$$

From the Yukawa couplings in (2.8), one obtains the following fermion mass matrices:

$$M_u \ = \ \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & h^u_{12}v_1e^{-i\theta_1} & 0 \\ h^u_{21}v_1e^{-i\theta_1} & 0 & h^u_{23}v_2e^{-i\theta_2} \\ 0 & h^u_{32}v_2e^{-i\theta_2} & h_{33}v_1e^{i\theta_1} \end{pmatrix} \ ,$$

$$M_d \ = \ \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & h^d_{12}v_2e^{i\theta_2} & 0 \\ h^d_{21}v_2e^{i\theta_2} & 0 & h^d_{23}v_1e^{i\theta_1} \\ 0 & h^d_{32}v_1e^{i\theta_1} & h^d_{33}v_2e^{i\theta_2} \end{pmatrix} \ ,$$

$$M^D_\nu \ = \ M_u\left(h^u_{ij} \to h^\nu_{ij}\right) \ , \qquad M_e = M_d\left(h^d_{ij} \to h^e_{ij}\right) \ ,$$

and

$$M^N \ = \ \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & h^N_{12}\Lambda e^{i\theta_3} & 0 \\ h^N_{21}\Lambda e^{i\theta_3} & 0 & h^N_{23}\Lambda e^{-i\theta_3} \\ 0 & h^N_{32}\Lambda e^{-i\theta_3} & h^N_{33}\Lambda e^{i\theta_3} \end{pmatrix} \ , \tag{2.15}$$

where $M_{u,d,e}$, M_ν^D and M^N are the u-, d-, and e-type, Dirac neutrino, and right-handed Majorana neutrino mass matrices respectively. It is interesting that all the mass matrices in (2.15) are Fritzsch-type. Using the framework on the Fritzsch fermion mass matrices,[12] one assumes that

$$|h^f_{ij}\, v_k| = |h^f_{ji}\, v_k| \approx \sqrt{m_{f_i} m_{f_j}} \ , \tag{2.16}$$

where $f = u,\ d$ and e, $k = 1,2$ and $i,j = 1,2,3$ are the family indices.

3 Light Neutrino Masses

The general form of the neutrino mass matrix is given by

$$M = \begin{pmatrix} 0 & M_\nu^D \\ M_\nu^{DT} & M^N \end{pmatrix} \ , \tag{3.1}$$

where M_ν^D and M^N are 3×3 matrices as given in Eq. (2.15). We will assume that

$$M_\nu^D \approx M_u \ , \tag{3.2}$$

based on the similarity of their structures and a common Yukawa coupling h^N for all families of right-handed neutrino, i.e.,

$$h^N = h^N_{ij} \ . \tag{3.3}$$

Therefore we have for the light neutrino mass matrix

$$M_\nu = -M_\nu^{DT}\,(M^N)^{-1}\,M_\nu^D = -M_u^T\,(M^N)^{-1}\,M_u \tag{3.4}$$

by virtue of the "see-saw" mechanism. Ignoring the phases in (3.4), since we are not interested in CP violating phenomena in this paper, we find

$$M_\nu \approx -\frac{m_t^2}{\frac{1}{\sqrt{2}}h^N\,\Lambda} \begin{pmatrix} 0 & \frac{m_u m_c}{m_t^2} & 0 \\ \frac{m_u m_c}{m_t^2} & \frac{m_c}{m_t} - 2\cdot\frac{m_c}{m_t}\sqrt{\frac{m_c}{m_t}} & -\sqrt{\frac{m_u m_c}{m_t^2}} + \sqrt{\frac{m_c}{m_t}} \\ 0 & -\sqrt{\frac{m_u m_c}{m_t^2}} + \sqrt{\frac{m_c}{m_t}} & 1 \end{pmatrix} . \tag{3.5}$$

Following the discussions in Ref. 13, one estimates the three light neutrino masses to be

$$M_{\nu_e} \approx \frac{m_u\,\sqrt{m_u m_c}}{\sqrt{2}\,h^N\Lambda} \ , \tag{3.6a}$$

$$M_{\nu_\mu} \approx \frac{2\sqrt{2}m_c\,\sqrt{m_u m_c}}{h^N\,\Lambda} \ , \tag{3.6b}$$

and

$$M_{\nu_\tau} \approx \frac{\sqrt{2}\,m_t^2}{h^N\,\Lambda} \ . \tag{3.6c}$$

The neutrino mass hierarchy is

$$M_{\nu_e} \ : \ M_{\nu_\mu} \ : \ M_{\nu_\tau} \approx \frac{m_u\,\sqrt{m_u m_c}}{2} \ : \ 2m_c\,\sqrt{m_u m_c} \ : \ m_t^2 \ . \tag{3.7}$$

To calculate the neutrino masses in (3.5), one has to know the value of m_t. Recently, the ARGUS collaboration has observed a large $B_d^0 - \overline{B}_d^0$ mixing.[14] It has been suggested that the large mixing results from a large m_t or a large charged Higgs effect.[16,17] Albright, Jarlskog and Lindholm[18] have shown that in the Fritzsch-type of models, the ARGUS result requires t-quark mass in the range of $95 \le m_t^{\rm phys} \le 105$ GeV without the charged Higgs effect and $70 \sim m_t^{\rm phys} \le 105$ GeV with the charged Higgs contribution respectively. The relation between the physical mass and running mass for the first order QCD correction is[19]

$$m_t^{\rm phys} = m_t(\mu = m_t)\left[1 + \frac{4}{3\pi}\,\alpha_s(m_t)\right] \ , \tag{3.8}$$

and[20]

$$m_t^{\rm phys} \approx 0.6\ m_t(\mu = 1\ {\rm GeV}) \ , \tag{3.9}$$

for the mass range between 40 GeV and 180 GeV. Therefore the ARGUS experiment would imply

$$70 \sim m_t(\mu = 1\ {\rm GeV}) \le 175\ {\rm GeV} \ . \tag{3.10}$$

Taking m_t (1 GeV) $\sim$ 150 GeV, m_c (1 GeV) $\sim$ 1.3 GeV, m_u (1 GeV) $\sim$ 5 MeV, and $\Lambda \approx 2f_a$ we get

$$M_{\nu_e} \approx 5.7 \times 10^{-2}\ \mathrm{eV} \cdot \left(\frac{10^7\ \mathrm{GeV}}{h^N f_a}\right) = 5.7 \times 10^{-7} \left(\frac{10^{12}\ \mathrm{GeV}}{h^N f_a}\right)\ \mathrm{eV}\ , \quad (3.11a)$$

$$M_{\nu_\mu} \approx 60\ \mathrm{eV} \left(\frac{10^7\ \mathrm{GeV}}{h^N f_a}\right) = 6.0 \times 10^{-4} \left(\frac{10^{12}\ \mathrm{eV}}{h^N f_a}\right)\ \mathrm{eV}\ , \quad (3.11b)$$

and

$$M_{\nu_\tau} \approx 6.4\ \mathrm{eV} \left(\frac{10^7\ \mathrm{GeV}}{h^N f_a}\right) = 64 \left(\frac{10^{12}\ \mathrm{GeV}}{h^N f_a}\right)\ \mathrm{eV}\ . \quad (3.11c)$$

The current experimental and cosmological constraints on the neutrino masses are given by[21]

$$M_{\nu_e} < 18\ \mathrm{eV}\ , \quad M_{\nu_\mu} < 250\ \mathrm{keV}\ , \quad M_{\nu_\tau} < 35\ \mathrm{MeV}\ , \quad (3.12a)$$

and[22]

$$\sum M_{\nu_i} < 65\ \mathrm{eV}\ , \quad (3.12b)$$

for stable neutrinos and[23]

$$M_{\nu_i} > 1\ \mathrm{MeV}\ , \quad (3.12c)$$

for unstable ν_i, respectively. From Eqs. (3.12a,c) we see that the only possible unstable neutrino is ν_τ with

$$1\ \mathrm{MeV} < M_{\nu_\tau} < 35\ \mathrm{MeV}\ . \quad (3.13)$$

Now we consider the following two cases,

(a) Stable ν_τ:

From (3.11c) and (3.11b), one finds that

$$h^N f_a > 10^{12}\ \mathrm{GeV}\ , \quad (3.14)$$

which requires $h^N > 10^4$ for $f_a \sim 10^8$ GeV and $h^N > 1$ for $f_a \sim 10^{12}$ GeV. With such unnatural large Yukawa coupling of Majorana right-handed neutrinos, we obtain

$$M_{\nu_e} < 2.9 \times 10^{-7}\ \mathrm{eV}\ , \quad M_{\nu_\mu} < 6.0 \times 10^{-4}\ \mathrm{eV} \quad \text{and } M_{\nu_\tau} < 64\ \mathrm{eV}. \quad (3.15)$$

With the above mass ranges we can understand the solar-neutrino puzzle via Mikheyev-Smirnov-Wolfenstein (MSW) mechanism[24] due mainly to the $\nu_e - \nu_\mu$ oscillations.

(b) Unstable ν_τ:

In this case, the ν_τ mass range of (3.12) imposes the following constraint:

$$6.4 \times 10^7\ \mathrm{GeV} > h^N f_a > 1.8 \times 10^6\ \mathrm{GeV}\ . \quad (3.16)$$

On the other hand, the value of M_{ν_μ} which must be less than 65 eV implies

$$h^N f_a > 9.2 \times 10^6\ \mathrm{GeV}\ . \quad (3.17)$$

Combining (3.16) and (3.17), we find that

$$6.4 \times 10^7\ \mathrm{GeV} > h^N f_a > 9.2 \times 10^6\ \mathrm{GeV}\ , \quad (3.18)$$

and the Yukawa coupling: $0.64 > h^N > 0.092$ for $f_a \sim 10^8$ GeV and $6.4 \times 10^{-5} > h^N > 9.2 \times 10^{-6}$ for $f_a \sim 10^{12}$ GeV, respectively, since the range of $h^N f_a$ in (3.17) is now small, we have a better prediction for neutrino masses which are given by

$$8.9 \times 10^{-3}\ \mathrm{eV} < M_{\nu_e} < 6.2 \times 10^{-2}\ \mathrm{eV}\ ,$$

$$9.3\ \mathrm{eV} < M_{\nu_\mu} < 65\ \mathrm{eV}\ , \quad (3.19)$$

and

$$1\ \mathrm{MeV} < M_{\nu_\tau} < 7\ \mathrm{MeV}\ .$$

It should be noted that to have a natural value of Yukawa couplings h, cases (a) and (b) correspond to the two extreme PQ breaking scales 10^{12} GeV and 10^8 GeV respectively.

4 Concluding Remarks

In this paper, we have studied the neutrino masses in the horizontal type of PQ symmetric model. The model has two doublets and one singlet Higgs fields which are the same as the DFSZ invisible axion model. The requirement of a Fritzsch matrix for quark masses uniquely determines the horizontal PQ charges for all families. The charged lepton, Dirac neutrino, and right-handed Majorana neutrino mass matrices are also all Fritzsch-type. Three light neutrino masses are obtained by the "see-saw" mechanism. The neutrino mass hierarchy is different from the "naive" mass hierarchy. For the case of all three stable neutrinos, the solar neutrino problem can be solved with the MSW mechanism due to the $\nu_e - \nu_\mu$ oscillations. For existing unstable ν_τ case, the masses of ν_e, ν_μ, and ν_τ are predicted to be order of 10^{-2} eV, 10 eV and 1 MeV, respectively.

The model also contains the following interesting feature: (i) it can easily be extended to the GUT's scheme by replacing the standard group;[11,25] (ii) the lepton mixings can be calculated as in Ref. 13; (iii) the model has the cosmological axion domain wall problem[26] just like the DSFZ model. This problem may be solved in the new inflationary model[27] or in the theory with $U(1)^H_{PQ}$ being embedded in a larger symmetry[28] so that the Lazavides-Shafi mechanism[29] can be applied.

Acknowledgment

This work was supported in part by the Natural Sciences and Engineering Research Council of Canada.

References

1. For reviews, see J.E. Kim, Phys. Rep. **150**, 1 (1987); H.Y. Cheng, Phys. Rep. **158**, 1 (1988).

2. R.D. Peccei and H.R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977); Phys. Rev. D **16**, 1791 (1977).

3. D.A. Dicus *et al.*, Phys. Rev. D **18**, 1829 (1978); D **22**, 839 (1980); M. Fukugita, S. Watamura and M. Yoshimura, Phys. Rev. D **26**, 1840 (1982); N. Iwamoto, Phys. Rev. Lett. **53**, 1198 (1984); J. Preskill, M.B. Wise and F. Wilczek, Phys. Lett. **120B**, 127 (1983); L.F. Abbott and P. Sikivie, Phys. Lett. **120B**, 133 (1983); M. Dine and W. Fischler, Phys. Lett. **120B**, 137 (1983).

4. S. Weinberg, Phys. Rev. Lett. **40**, 223 (1978); F. Wilczek, Phys. Rev. Lett. **40**, 279 (1978).

5. M. Dine, W. Fischler and M. Srednicki, Phys. Lett. **B104**, 199 (1981); A.P. Zhitnitskii, Sov. J. Nucl. Phys. **31**, 260 (1980).

6. J.E. Kim, Phys. Rev. Lett. **43**, 103 (1979); M.A. Shifman, A.I. Vainshtein, and V.I. Zakharov, Nucl. Phys. **B166**, 493 (1980).

7. T. Yanagida, in Proc. Workshop on Unified theory and baryon number in the universe, eds. O. Sawada and A. Sugamoto, (KEK, 1979); M. Gell-Mann, P. Ramond and R. Slansky, in Supergravity, eds. P. van Nieuwenhuizen and D. Freedman (North-Holland, 1980).

8. R.N. Mohapatra and G. Senjanović, Z. Phys. **C 17**, 53 (1983); P. Langacker, R.D. Peccei and T. Yanagida, Mod. Phys. Lett. **A1**, 541 (1986); K. Kang and

M. Shin, Mod. Phys. Lett. **A1**, 585 (1986); M. Shin, Phys. Rev. Lett. **59**, 2515 (1987).

9. M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1972).
10. H. Fritzsch, Phys. Lett. **70B**, 436 (1977); **73B**, 317 (1978).
11. A. Davidson and K.C. Wali, Phys. Rev. Lett. **48**, 11 (1982).
12. H. Georgi and D. Nanopoulos, Nucl. Phys. **B155**, 52 (1979).
13. K. Kang and M. Shin, Mod. Phys. Lett. **A1**, 585 (1986); Phys. Lett. **185B**, 163 (1987); and Nucl. Phys. **B287**, 687 (1987).
14. ARGUS Collab., H. Albrecht *et al.*, Phys. Lett. **B192** (1987) 245.
15. I.I.Y. Bigi and A.I. Sanda, Phys. Lett. **194B**, 307 (1987); L.L. Chau and W.Y. Keung, UC Davis preprint UCD-87-02; V. Barger, T. Han, D.V. Nanopoulos, and R.T.N. Phillips, Phys. Lett. **194B**, 312 (1987); D. Du and Z. Zhao, Phys. Rev. Lett. **59**, 1072 (1987); J.F. Donoghue *et al.*, Phys. Lett. **195B**, 285 (1987); H. Harari, and Y. Nir, Phys. Lett. **195B**, 586 (1987); W. Hou and A. Soni, Phys. Lett. **196B**, 92 (1987); J.R. Cudell, F. Halzen, X.G. He and S. Pakvasa, Phys. Lett. **196B**, 227 (1987).
16. S.L. Glashow and E.E. Jenkins, Phys. Lett. **196**, 233 (1987).
17. C.Q. Geng and J.N. Ng, TRIUMF preprint, to be published.
18. C.H. Albright, C. Jarlskog and B. Lindholm, Phys. Rev. D **38**, 872 (1988).
19. J. Gasser and H. Leutwyler, Phys. Rep. **87**, 77 (1982).

20. Cf.Y. Nir, Nucl. Phys. **B306**, 14 (1988).
21. Particle Data Group, Phys. Lett. **B170** (1986); The ARGUS collaboration, DESY preprint, DESY 87-148.
22. R. Cowsik and J. McClelland, Phys. Rev. Lett. **29**, 669 (1972).
23. H. Harari and Y. Nir, Nucl. Phys. **B292**, 251 (1987).
24. S.P. Mikheyev and A.Y. Smirnov, Sov. J. Nucl. Phys. **42**, 913 (1985); L. Wolfenstein, Phys. Rev. D **17**, 2369 (1978).
25. A. Davidson, V.P. Nair and K.C. Wali, Phys. Rev. D **29**, 1504; **15**, (1984).
26. P. Sikivie, Phys. Rev. Lett. **48**, 1156 (1982).
27. S.Y. Pi, Phys. Rev. Lett. **52**, 1725 (1984).
28. F. Wilczek, Phys. Rev. Lett. **49**, 1549 (1982).
29. G. Lazarides and Q. Shafi, Phys. Lett. **B115**, 21 (1982).