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Neutrino Physics – Summary Talk*

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* Invited talk given at the "New Directions in Neutrino Physics at Fermilab Workshop", September 14-16, 1988.

1. INTRODUCTION

This has been a most interesting and timely workshop. I found the talks to be both educational and stimulating. In particular, a number of novel ideas for neutrino experiments and facilities surfaced. Most important, from my perspective, specific suggestions for the next generation of precision measurements and neutrino oscillation searches were discussed. I hope that the proponents of those ideas will follow through with concrete proposals.

We began the workshop with Leon Lederman's introductory remarks and charge to the participants. He asked them to address the question: What new opportunities in neutrino and kaon physics will be made possible by Fermilab's ongoing and proposed upgrades? He was primarily referring to the proposed 120 Gev Main Injector which would have a proton flux about 100 times that of the Tevatron and could produce a copious source of neutrinos and kaons. He also told the audience consisting mainly of experimentalists (in jest, of course) to ignore theorists or "throw theorists out the window." Being a theorist, I can not completely endorse that sentiment. In some areas, theory is presently far ahead of experiment and cannot be ignored. I am thinking of the standard model, its successes and very precise predictions for neutrino cross-sections, CP violation, etc. There a primary role of experiment at this time should be to test the standard model by high precision measurements. Observation of a deviation from the standard model prediction would be a sign of "new physics." In the more speculative areas such as neutrino mass, mixing, electromagnetic moments, etc., theory can provide some guidance, but experimenters really should follow Leon's advice and rely on their own instincts. In my opinion, the adage that any new measurement that improves a bound or constraint by an order of magnitude is worth doing, definitely applies.

With the above theme in mind, I have organized this summary talk as follows: In section 2, I describe the state of neutrino phenomenology. Emphasis is placed on $\sin^2\theta_W$, its present status and future prospects. In addition, some signatures of "new physics" are described. Kaon physics at Fermilab is briefly discussed in section 3. I concentrate on the intertesting rare decay $K_L \to \pi^0 e^+ e^-$ which may be a clean probe of direct CP violation. Neutrino mass, mixing, and electromagnetic moments are surveyed in section 4.

There, I describe the present state and future direction of accelerator based experiments. Finally, I conclude in section 5 with an outlook on the future. Throughout this summary, I have drawn from and incorporated ideas discussed by other speakers at this workshop. However, I have also tried to combine their ideas with my own perspective on neutrino physics and where it is headed.

2. STANDARD MODEL PHENOMENOLOGY – $\sin^2\theta_W$

Neutrino experiments at accelerator facilities generally fall into two categories: 1) Measurements of the Standard Model's properties and parameters, particularly $\sin^2\theta_W$, and 2) Searches for neutrino oscillations. In this section I will discuss the first of these. Oscillation phenomena will be the subject of section 4.

Let me begin by mentioning some topics covered by other speakers. We heard several nice talks on using neutrinos to study QCD. J. Morfin presented an overview of structure function determinations by deep-inelastic scattering experiments. It is interesting that neutrino determined structure functions from different experiments are all in basic agreement while muon determined structure functions (which should be more accurate) are somewhat controversial. Nevertheless, all experiments are pretty much consistent with QCD expectations for $\Lambda_{\overline{MS}} \simeq 100-200$ MeV. Structure functions will of course be the life blood of SSC studies; so, it was nice to learn that a concentrated group effort is underway to determine parton distributions as well as possible. Of course, HERA should contribute a wealth of data to that undertaking when it turns on.

Neutrino production of dileptons has been a source of some controversy in the past. The present situation was nicely summarized by W. Smith. He concluded that in the case of like-sign dileptons, the observed cross-sections are now consistent with perturbative QCD expectations. Opposite sign dilepton production is also consistent with charm production predictions. The latter topic is important for my subsequent discussion of $\sin^2\theta_W$, since charm production threshold uncertainties presently constitute the major error in deep-inelastic $\nu_\mu N$ scattering determinations of that parameter. R. Brock discussed those determinations and also focused on charm threshold effects in his critique of their conclusions. J. Panman described the very successful CERN program in deep-inelastic

$\nu_\mu N$ scattering and its possible future direction. We were told by I. Savin that the UNK 3 TeV fixed target program will have an ambitious neutrino program that should help clarify open issues in deep-inelastic scattering. Similarly, S. Denisov described plans for a tagged neutrino facility at Serpukov. R. Bernstein discussed his ideas for a tagged neutrino facility at Fermilab's Tevatron. The nice feature of such a facility would be direct determination of the individual $\nu_\mu, \overline{\nu}_\mu, \nu_e$ and $\overline{\nu}_e$ fluxes. That would allow precision measurements of individual charged and neutral current cross-sections as well as small mixing oscillation searches.

In the $\nu_\mu - e$ scattering arena, A. Capone updated the ongoing high statistics CHARM II experiment at CERN which hopes to obtain a ± 0.005 statistical uncertainty in its measurement of $\sin^2\theta_W$. Plans for a precise determination of $\sin^2\theta_W$ (within $\pm 1\%$) at LAMPF using a Large Cherenkov Detector (LCD) were outlined by D.H. White. That experiment would also search for oscillation effects and rare decays. The possibility of building a muon storage ring to do neutrino physics was discussed by W. Lee and D. Neuffer. With such a facility, neutrino fluxes would be very well determined and many precision measurements would be possible. I was pleasantly surprised to learn that the Fermilab $\overline{p}$ "Debuncher" could, even in its present configuration, function as an 8.9 GeV muon storage ring. It would be interesting to try out that mode, perhaps for an elastic $\overset{(-)}{\nu}_\mu p$ scattering or oscillation experiment.

To help further motivate some of the new ideas for future neutrino experiments discussed by others, I will describe results obtained from global fits to all existing neutral current data and their implications. Neutrino experiments played a crucial role in that analysis and will be extremely important in future testing of the standard model.

Several years ago, a collaboration[1,2] was formed for the purpose of collecting and analyzing all neutral current data. The idea was to carefully scrutinize experimental and theoretical uncertainties in those results and to consistently include effects of electroweak radiative corrections. That undertaking involved examining a great many diverse measurements (approximately 180) which included: deep-inelastic $\nu_\mu N$ scattering, $W^\pm$ and Z masses, eD scattering asymmetry, atomic parity violation, νe scattering, νp scattering, e^+e^- annihilation, μC scattering, etc. The goals were:

1. To test the standard $SU(2)_L \times U(1)$ model at the tree and quantum loop level.
2. Provide a precise determination of $sin^2\theta_W$ which could be used to rule out or at least constrain various grand unified theories (GUTS).
3. Look for hints of "new physics."

Before stating some of the main results of that global analysis[1], I will briefly describe the basic assumptions that went into it. We assumed 3 generations of fermions and one Higgs doublet with an underlying $SU(3)_C \times SU(2)_L \times U(1)$ gauge symmetry constituted the standard model. Within that framework, radiative corrections to all relevant neutral current processes as well as $W^\pm$ and Z mass formulas were accounted for. Those corrections depend on the couplings and masses in the model. Two of those parameters, m_t (the top quark mass) and m_H (the Higgs scalar mass) are presently undetermined. Whereas, the radiative corrections are not very sensitive to m_H variations, they are sensitive to m_t if it is $\gtrsim$90 GeV.[3,4] Therefore, in some parts of our analysis, we allowed m_H to vary from 10 GeV to 1 TeV and merely required $m_t \lesssim 100$ GeV while in other parts $m_H \simeq 100$ GeV and $m_t = 45$ GeV were assumed for definiteness. Because of the sensitivity to large m_t, we were also able to place bounds on m_t.

Some of the principal results of the global analysis of all existing neutral current data and $W^\pm$ and Z masses were:

1. There is at present no evidence for any deviation from the standard model.
2. For $m_t \lesssim 100$ GeV and $m_H \lesssim 1$ TeV, we found the world average $\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2 = 0.230 \pm 0.0048$. For $m_t \simeq 45$ GeV and $m_H \simeq 100$ GeV, the uncertainties were lowered slightly to $\sin^2\theta_W = 0.230 \pm 0.0044$.
3. Allowing $\rho \equiv m_W^2/m_Z^2 \cos^2\theta_W$ as well as $\sin^2\theta_W$ to vary, we obtained from a two parameter fit to all data $\sin^2\theta_W = 0.229 \pm 0.0064$, $\rho = 0.998 \pm 0.0086$.
4. For $m_t \simeq 45$ GeV and $m_H = 100$ GeV, radiative corrections are confirmed at about the 3σ level, primarily in the comparison of deep-inelastic $\nu_\mu N$ scattering with m_W and m_Z.

5. Consistency of all data at the quantum loop level requires $m_t \lesssim 200$ GeV for $m_H \lesssim 1$ TeV. If $m_H \lesssim 100$ GeV, the tighter constraint $m_t \lesssim 180$ GeV is obtained. Those constraints also apply[5] to a 4th generation mass difference $(m_{t'} - m_{b'})$.
6. Lower limits on extra Z' boson masses were obtained for a variety of popular GUT models. The bounds ranged from 120 GeV to 300 GeV depending on their specific couplings to fermions.

The results described above have many implications. I will elaborate on a few topics in order to show the power of global fits and the need to push for higher precision.

Radiative Corrections and $\sin^2\theta_W$

The weak mixing angle θ_W plays a central role in the $SU(2)_L \times U(1)$ model. Writing out the electroweak neutral current interaction Lagrangian

$$\begin{aligned} \mathcal{L}_{\rm int} = & -eA^\mu(x)\sum_f Q_f \bar{f}\gamma_\mu f - \frac{2e}{\sin 2\theta_W} Z^\mu(x) \\ & \times \sum_f \left(T_{3f}\bar{f}_L\gamma_\mu f_L - \sin^2\theta_W Q_f \bar{f}\gamma_\mu f\right) \\ e = & \, g_2 \sin\theta_W \end{aligned} \tag{2.1}$$

with T_{3f} = weak isospin and Q_f =electric charge, we see that θ_W occurs both at the fermion weak neutral current level and in the normalization of the $SU(2)_L$ coupling g_2 relative to the electric charge e. In addition, for the simplest Higgs doublet scenario, it enters the W-Z mass relationship via

$$m_W = m_Z \cos\theta_W \tag{2.2}$$

So, one tests the standard model and the underlying concept of electroweak unification by measuring $\sin^2\theta_W$ in as many different ways as possible. A deviation in the value obtained from one experiment as compared with another would signal new physics.

Of course, radiative corrections must be accounted for in any precise determination of $\sin^2\theta_W$, so that they will not be confused with new physics. In some cases, electroweak radiative corrections can be quite large. For example, consider the lowest order natural relationship

$$\sin^2\theta_W^0 = \left(e^0/g_2^0\right)^2 = 1 - \left(m_W^0/m_Z^0\right)^2 \tag{2.3}$$

where 0 indicates bare (unrenormalized) parameters. In terms of physical measureable quantities, that relationship is modified by *finite* $O(\alpha)$ loop corrections. The size of those corrections depends on the definitions employed; but for typical definitions they can be quite large. Defining the renormalized weak mixing angle by

$$\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2 \tag{2.4}$$

where m_W and m_Z are physical masses and the renormalized charges e and g_2 via

$$\alpha = e^2/4\pi = 1/137.036 \tag{2.5a}$$

$$G_\mu = \frac{g_2^2}{4\sqrt{2}m_W^2} = 1.16636 \times 10^{-5}\text{GeV}^{-2} \tag{2.5b}$$

leads to[3]

$$\begin{aligned} m_W = m_Z\cos\theta_W &= \left(\frac{\pi\alpha}{\sqrt{2}G_\mu}\right)^{\frac{1}{2}} \frac{1}{\sin\theta_W(1-\Delta r)^{\frac{1}{2}}} \\ &= \frac{37.281\text{GeV}}{\sin\theta_W(1-\Delta r)^{\frac{1}{2}}} \end{aligned} \tag{2.6}$$

where[1]

$$\Delta r = 0.0713 \pm 0.0013 \tag{2.7}$$

for $m_t \simeq 45$ GeV and $m_{\text{Higgs}} \simeq 100$ GeV. The radiative corrections in $\Delta r \simeq O(\alpha)$ are large primarily due to fermion vacuum polarization renormalization of e relative to g_2. It is, however, a mistake to call that effect a QED correction. Infinite fermion loops enter in the photon propagator as well as the $W^\pm$ and Z propagators. The relative correction is calculable (finite) only because of electroweak unification.

When either m_W or m_Z are used to obtain $\sin^2\theta_W$ via Eq. (2.6), Δr causes a sizeable $\simeq 7\%$ shift in the value found. Similarly, neutral current scattering cross sections and interference measurements must be corrected for $O(\alpha)$ quantum loops in extracting $\sin^2\theta_W$. The effects of such electroweak radiative corrections are illustrated in Table 2.1 where several determinations of $\sin^2\theta_W$ are summarized. At present, deep-inelastic ν_μ scattering provides the best determination of $\sin^2\theta_W$. In fact, much of the uncertainty (about $\pm$ 0.004) in that extraction is theoretical in the sense that a model is employed to correct for charm threshold effects. It, therefore, appears that those measurements have been pushed

Table 2.1: Values of $\sin^2\theta_W = 1 - m_W^2/m_Z^2$ before and after electroweak radiative corrections (R.C.) are included. The values $m_t \simeq 45$GeV and $m_H \simeq 100$GeV were employed in the radiative corrections. Atomic parity violation results have been updated to include recent data.

Experiment	$\sin^2\theta_W^{\text{unc.}}$	R. C.	$\sin^2\theta_W$
Atomic P.V.	0.211 ± 0.016	+0.007	0.218 ± 0.016
eD Asymmetry	0.226 ± 0.015 ± 0.013	-0.005	0.221 ± 0.015 ± 0.013
$\overset{(-)}{\nu}_\mu e$	0.221 ± 0.019	+ 0.002	0.223 ± 0.019
$\overset{(-)}{\nu}_\mu P$	0.208 ± 0.033	+ 0.002	0.210 ± 0.033
$\nu_\mu N$ deep-inel.	0.242 ± 0.003 ± 0.005	-0.009	0.233 ± 0.003 ± 0.005
$m_W = 80.9 \pm 1.4$ (UA1 - UA2)	0.212 ± 0.008	+0.017	0.229 ± 0.008
$m_Z = 91.9 \pm 1.8$ (UA1 - UA2)	0.208 ± 0.011	+0.022	0.230 ± 0.011
World Average			0.230 ± 0.0044

about as far as possible unless one supplements the analysis with an experimental study of charm production via unlike sign dimuons,[1] or can find a way to bypass that uncertainty. I will say more about that later.

Comparing the $\sin^2\theta_W^{\text{unc.}}$ and $\sin^2\theta_W$ columns in Table 2.1, it is clear that the uncorrected value obtained from deep-inelastic $\nu_\mu - N$ scattering differs from the m_W and m_Z values, but the corrected $\sin^2\theta_W$ are in good agreement. So, the standard model has been tested at the level of its $O(\alpha)$ radiative corrections (at about the 3σ level) if m_t is actually near 45 GeV.

Experimentalists should strive to measure $\sin^2\theta_W$ as precisely as possible. Fortunately, m_Z measurements to better than ±50 MeV at SLC and LEP will determine $\sin^2\theta_W$ to ± 0.00025 via Eq. (2.6), (but only after m_t and m_H are pinpcinted). Given such high precision, what role can other experiments play? To test the standard model or look for hints of new physics, one *must* compare distinct measurements. Neutrino scattering determinations of $\sin^2\theta_W$ nicely complement the mass measurements, since they are sensitive to different types of new physics. The present uncertainty of about $2 \sim 3\%$ in the average $\sin^2\theta_W$ should provide a benchmark that all new proposals should at least match and strive to better.

Top Quark Mass

Radiative corrections to $\sin^2\theta_W$ can be quite sensitive to the value of m_t, if it is large. In fact, they grow like $\alpha m_t^2/m_W^2$ in most processes.[3,4] Only $\sin^2\theta_W$ determined from deep-inelastic $\nu_\mu N$ scattering (due to a subtle cancellation) is rather insensitive to variations in m_t (see Table 2.2). Therefore, the present good agreement between $\sin^2\theta_W$ obtained from $\nu_\mu N$ data and other experiments such as m_W and m_Z measurements gives us some confidence that m_t is probably in the range $45 \sim 150$ GeV. In fact, as previously stated, by varying m_t in the radiative corrections, we found (for $m_H \leq 1$ TeV)

$$m_t \lesssim 200\text{GeV} \qquad (90\%\ \text{CL}) \tag{2.8}$$

(That bound also applies to a fourth generation quark mass difference $|m_{t'} - m_{b'}|$.)

The effect of changing m_t on our global fits is illustrated in Table 2.3. There I have given the values of $\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2$ and $\sin^2\theta_W\,(m_W)_{\overline{\rm MS}}$ defined by $\overline{\rm MS}$ (modified minimal subtraction). Note that those two distinct definitions of the weak mixing angle differ by terms of $O\left(\alpha m_t^2/m_W^2\right)$; hence, their difference grows as m_t increases. Whereas, $\sin^2\theta_W$ extracted from deep-inelastic $\nu_\mu N$ data is quite insensitive to m_t, that is not so for $\sin^2\theta_W\,(m_W)_{\overline{\rm MS}}$ (see Table 2.2). In contradistinction, $\sin^2\theta_W$ extracted from either m_W or m_Z via Eq. (2.6) is sensitive to the value of m_t through Δr, while $\sin^2\theta_W\,(m_W)_{\overline{\rm MS}}$ extracted from m_W or m_Z (particularly m_W) is much less dependent on m_t. (Of course, a precise determination of both m_W and m_Z would determine $\sin^2\theta_W = 1 - m_W^2/m_Z^2$

Table 2.2: Experimental determinations of $\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2$ for several top quark masses. For uncertainties, see table 2.1.

		$\sin^2\theta_W$	
Experiment	$m_t = 45$ GeV	$m_t = 90$ GeV	$m_t = 180$ GeV
Atomic Parity Violation	0.218	0.217	0.205
eD Asymmetry	0.221	0.218	0.212
$\overset{(-)}{\nu}_\mu e$	0.223	0.220	0.212
$\overset{(-)}{\nu}_\mu p$	0.210	0.207	0.200
$\overset{(-)}{\nu}_\mu N$	0.233	0.232	0.231
m_W	0.229	0.226	0.219
m_Z	0.230	0.226	0.216
Weighted Average	0.230	0.227	0.223

independent of m_t or any radiative corrections.) We, therefore, have a situation in which at present, neither definition's value can be very precisely given without some assumption regarding m_t.

Table 2.3: World average values for the weak mixing angle as a function of m_t (keeping $m_H = 100$ GeV).

m_t (GeV)	$\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2$	$\sin^2\theta_W\,(m_W)_{\overline{\text{MS}}}$
25	0.229 ± 0.0044	0.227 ± 0.0044
45	0.230	0.228
60	0.230	0.228
100	0.227	0.229
200	0.222	0.233
400	0.209	0.248

Gauge Couplings and Guts

The standard $SU(3)_C \times SU(2)_L \times U(1)$ model of strong and electroweak interactions contains 18 independent couplings and masses. Grand Unified Theories[6,7] (GUTS) correlate the three gauge couplings g_3, g_2 and g_1 by embedding the standard model in a compact simple group such as $SU(5)$, $SO(10)$, E_6 etc. Indeed, the high degree of symmetry naturally renders the bare couplings equal, explains charge-color quantization and promotes $\sin^2\theta_W^0$ from an infinite counterterm parameter to a rational number (generally 3/8). Unfortunately, GUTS have so far provided little new insight regarding the 15 mass and quark mixing parameters. Therefore, although GUTS represent a significant theoretical advancement, they cannot be the final word.

Here, I will update the gauge coupling values. Those quantities are extremely important because they provide much of the basis for our belief in GUTS and a severe constraint on model building. In the case of the QCD coupling, the situation has not changed significantly during the last few years. Upsilon decays, high energy jet data, and tau decays are consistent with

$$\Lambda^{(4)}_{\overline{MS}} \simeq 150^{+150}_{-75}\mathrm{MeV} \tag{2.9}$$

(The errors are very conservative.) Assuming $m_t \simeq 45\mathrm{GeV}$ and using $m_W = 80.7\mathrm{GeV}$ (which corresponds to $\sin^2\theta_W = 0.23$), that range leads to[8]

$$\alpha_3(m_W) = 0.107^{+0.013}_{-0.008} \tag{2.10}$$

The analogous electroweak parameters, also defined by $\overline{\mathrm{MS}}$ (modified minimal subtraction) have the short-distance values[1,9]

$$\alpha^{-1}(m_W) = 127.8 \pm 0.3 \tag{2.11}$$

$$\sin^2\theta_W(m_W) = 0.228 \pm 0.0044 \tag{2.12}$$

where the value of $\sin^2\theta_W(m_W)$ in Eq. (2.12) follows from the result $\sin^2\theta_W = 0.230 \pm 0.0044$ (for $m_t = 45$ GeV) found by the global analysis described above. As previously discussed, the $\overline{\mathrm{MS}}$ definition and $\sin^2\theta_W \equiv 1 - m_W^2/m_Z^2$ differ by order α radiative corrections which for $m_t = 45$ GeV imply[1] $\sin^2\theta_W(m_W) = 0.9907\sin^2\theta_W$. For other values of m_t, that relationship is modified as are the central values of both $\sin^2\theta_W(m_W)$ and

$\sin^2\theta_W$ extracted from experiment (i.e., the radiative corrections to each experiment also depend on m_t). (See Tables 2.2 and 2.3.) It should be noted that the world average for the weak mixing angle in Eq. (2.12) has increased from the old $\sin^2\theta_W(m_W) \simeq 0.219$ value of several years ago primarily because of more precise deep-inelastic ν_μ scattering data and refinements in the $W^\pm$ and Z mass determinations. Also, as indicated in Table 2.3, if m_t is > 45 GeV, $\sin^2\theta_W(m_W)$ increases even more. That higher $\sin^2\theta_W(m_W)$ value has very important implications for GUTS, as we shall see.

Employing the relationships

$$\alpha_1(m_W) = 5\alpha(m_W)/3\cos^2\theta_W(m_W) \tag{2.13a}$$

$$\alpha_2(m_W) = \alpha(m_W)/\sin^2\theta_W(m_W) \tag{2.13b}$$

leads to the gauge coupling values

$$\alpha_1(m_W) = 0.0169 \pm 0.0001 \tag{2.14a}$$

$$\alpha_2(m_W) = 0.0344 \pm 0.0007 \tag{2.14b}$$

Several years ago, the central value of $\alpha_2(m_W)$ was 0.036 because $\sin^2\theta_W(m_W)$ was thought to be smaller. Assuming that there are no other new thresholds between the standard model's mass scale of m_W and the grand unification scale of m_X, one can evolve the gauge couplings to higher energies using (for 3 generations)[7]

$$\mu\frac{\partial}{\partial\mu}\alpha_i(\mu) = b_i\alpha_i^2 + \ldots, i = 1,2,3 \tag{2.15a}$$

$$\begin{pmatrix} b_1 \\ b_2 \\ b_3 \end{pmatrix} = -\frac{1}{2\pi}\begin{pmatrix} -41/10 \\ 19/6 \\ 7 \end{pmatrix} \tag{2.15b}$$

and the values of $\alpha_i(m_W)$ given above. If the three couplings meet at a single point, that would be clear evidence for grand unification. When $\alpha_2(m_W)$ was 0.036, they tended to meet near $\mu \sim 2\times10^{14}$GeV. That meeting was taken as strong confirmation of GUTS and perhaps an indication of no new physics thresholds at low or intermediate mass scales. Using the new value for $\alpha_2(m_W)$ in Eq. (2.14b), one finds that is no longer the case (see fig. 2.1).

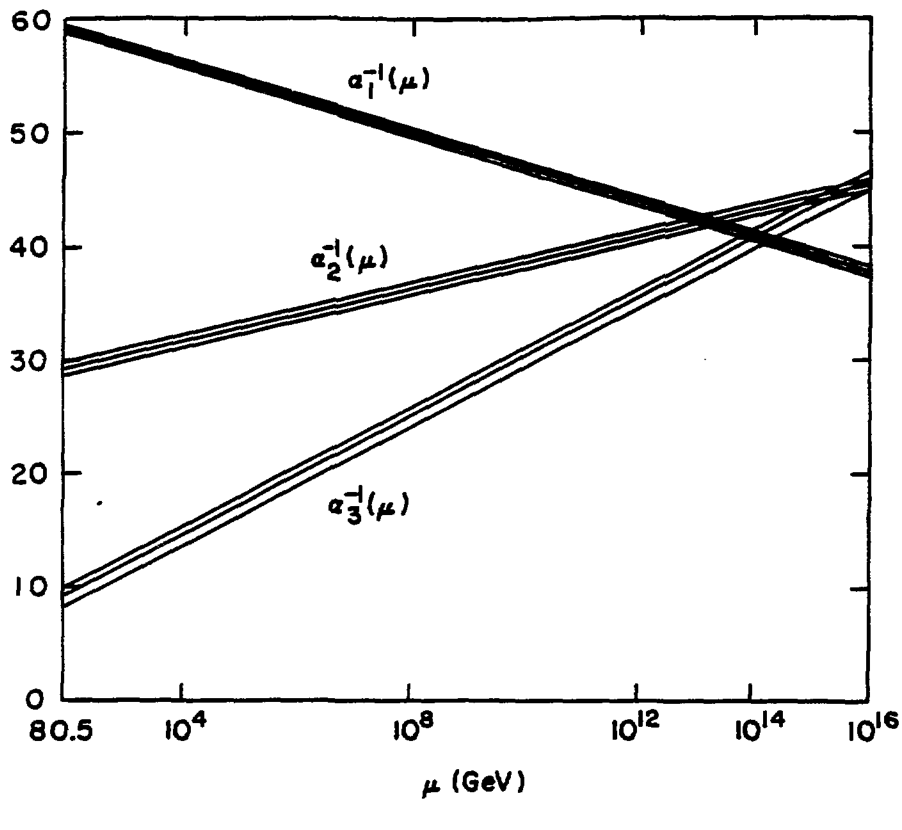

Figure 2.1: Evolution of the $\alpha_i^{-1}(\mu)$, $i = 1, 2, 3$ couplings assuming no new physics beyond the standard model.

The couplings $\alpha_1(\mu)$ and $\alpha_3(\mu)$ continue to meet near 1.5×10^{14}GeV; however, $\alpha_2(\mu)$ now crosses $\alpha_1(\mu)$ at 1.5×10^{13}GeV and meets $\alpha_3(\mu)$ near 1.0×10^{16}GeV. Is grand unification ruled out? No, this development merely implies that new physics thresholds between m_W and m_X must change the evolution of the couplings such that they meet at a single value. In my opinion, the near equality of the couplings at high energies that we find using Eq. (2.15) should still be taken as a strong indication of grand unification. At issue is: What new physics rectifies the evolution and at what energy will it be manifested?

The above remarks are nicely illustrated by the SU(5) Georgi-Glashow model.[6] In the so-called minimal version, one assumes the existence of a great desert between m_W and m_X, the unification mass scale. That simplistic assumption had an appealing consequence, it led to rather definite testable predictions. (The predictions hold in many GUTS with great deserts.) Indeed, using $\alpha^{-1}(m_W) \simeq 127.8 \pm 0.3$ and $\Lambda^{(4)}_{\overline{\mathrm{MS}}} = 150^{+150}_{-75}$MeV, one predicts

$$m_X = \left(2.0^{+2.1}_{-1.0}\right) \times 10^{14}\mathrm{GeV} \tag{2.16}$$

$$\sin^2\theta_W(m_W) = 0.214^{+0.003}_{-0.004} \tag{2.17}$$

Unfortunately, both of these predictions are now ruled out by experiment. The IMB proton decay bound[10]

$$1/\Gamma\left(P \to e^+\pi^-\right) \geq 3.1 \times 10^{32}\mathrm{yr} \tag{2.18}$$

require $m_X \gtrsim 7 \times 10^{14}$GeV, while the $\sin^2\theta_W(m_W)$ prediction conflicts with the world average in Eq. (2.12). (It gets worse if m_t is > 45 GeV.) The latter disagreement is, of course, just another quantitative way of describing the apparent lack of unification of gauge couplings in fig. 2.1 when current $\alpha_i(m_W)$ values are employed. These failures of the minimal SU(5) model do not rule out SU(5) as a viable grand unification group. They do indicate that new physics appendages in the form of additional scalars or fermions must be introduced[11] to render $m_X \geq 10^{15}$GeV and increase the prediction for $\sin^2\theta_W(m_W)$. Another possibility is that a bigger GUT such as SO(10) or E_6 with intermediate stages of symmetry breaking must be employed. I will now describe how low energy supersymmetry[12] may do the trick for SU(5) or any other GUT.

The basic idea of supersymmetry is that each known boson (fermion) has a fermion (boson) partner. In those scenarios, the $\alpha_i(\mu)$ evolution in equations (2.15) charge when we pass the supersymmetry thresholds. In leading order, one finds for three generations of fermions and N_H light Higgs doublets[12,13]

$$\begin{pmatrix} b_1 \\ b_2 \\ b_3 \end{pmatrix} = -\frac{1}{2\pi} \begin{pmatrix} -6 - \frac{3}{10} N_H \\ -\frac{1}{2} N_H \\ 3 \end{pmatrix} \tag{2.19}$$

Taking $N_H = 2$ (the minimal value) and using the $\alpha_i(m_W)$ values Eqs. (2.10) and (2.14) as input, we can solve for m_{SUSY} and m_X. One finds in leading order[9]

$$\ln(m_X/m_W) \simeq \frac{\pi}{2}\left\{\frac{1}{\alpha_2(m_W)} - \frac{1}{\alpha_3(m_W)}\right\} \tag{2.20}$$

independent of m_{SUSY}. Using the values of $\alpha_2(m_W)$ and $\alpha_3(m_W)$ then gives the range of predictions

$$m_X \simeq 2 \times 10^{14} \sim 2 \times 10^{16}\mathrm{GeV} \tag{2.21}$$

The lower mass range corresponds to very large m_{SUSY} while the higher values require m_{SUSY} to be nearer m_W. In SUSY GUTS, one expects the gauge boson mediated decay rate to be

$$\begin{aligned} 1/\Gamma(p \to e^+\pi^0) &\simeq 1.3 \times 10^{29\pm0.7} \\ &\times \left(\frac{m_X}{2 \times 10^{14}\mathrm{GeV}}\right)^4 \mathrm{yr} \qquad \text{(SUSY)} \end{aligned} \tag{2.22}$$

The IMB bound in Eq. (2.18) then rules out the $m_X \lesssim 10^{15}$GeV region in Eq. (2.21) but leaves open the possibility of $m_{SUSY} \lesssim 10^6$GeV as the "new physics" we are looking for. SUSY GUTS also predict[1]

$$\sin^2\theta_W(m_W) = 0.237^{+0.003}_{-0.004} - \frac{4}{15}\frac{\alpha}{\pi}\ln(m_{SUSY}/m_W) \tag{2.23}$$

which is in good accord with experiment (see table 2.1), particularly if m_t, $\Lambda^{(4)}_{\overline{MS}}$ or m_{SUSY} is on the high side. This example illustrates how a new physics threshold (supersymmetry in this case) can bring GUTS into agreement with low energy phenomenology. It also demonstrates the complementarity between proton decay and high energy experiments. If $m_{SUSY} \lesssim 10$TeV, it is likely to be discovered at the SSC, and in this example, the proton decay rate is too slow to observe. On the other hand, if m_{SUSY} is beyond 10 TeV, the value of m_X is lower and the detection of proton decay is more likely. Of course, the use of a single supersymmetry mass scale is rather simplistic. Nevertheless, this example illustrates the importance of pushing the search for proton decay and the precise determination of $\sin^2\theta_W$ as far as possible.

Extra Z' Bosons

Additional neutral gauge bosons (generically called Z' bosons) arise naturally in GUTS larger than SU(5).[14] The SO(10) model has one such additional boson which is often denoted by Z_χ while E_6 has Z_χ as well as a second flavor diagonal neutral boson Z_ψ. Their couplings are specified up to renormalization effects which are calculable if the entire particle spectrum is known.

The interaction Lagrangian for Z_χ and Z_ψ is given by

$$\mathcal{L}_{\rm int} = -\sqrt{3/8}g_2 \tan\theta_W \sum_{i=\chi,\psi} \sqrt{\lambda_i} Z_i^\mu J_\mu^i \tag{2.24a}$$

$$J_\mu^i = \sum_f \left(Q_{f_R}^i \bar{f}_R \gamma_\mu f_R + Q_{f_L}^i \bar{f}_L \gamma_\mu f_L\right) \tag{2.24b}$$

where $\sqrt{\lambda_i}$ is a renormalization parameter that is generally ≈ 1. The charges in Eq. (2.24) are completely specified (for each generation)

$$Q_{e_L}^\chi = Q_{\nu_L}^\chi = -Q_{d_R}^\chi = 3Q_{u_R}^\chi = 3Q_{e_R}^\chi = -3Q_{u_L}^\chi \; -3Q_{d_L}^\chi = 1 \tag{2.25a}$$

$$Q_{\nu_L}^\psi = Q_{e_L}^\psi = Q_{u_L}^\psi = Q_{d_L}^\psi = -Q_{e_R}^\psi = -Q_{u_R}^\psi = -Q_{d_R}^\psi = \sqrt{5/27}. \tag{2.25b}$$

One expects Z_χ and Z_ψ to mix with one another (and probably mix somewhat with the ordinary Z). Ignoring potential small mixing[1] with the usual Z, the mass eigenstates can be denoted by $Z(\beta)$ and $Z'(\beta)$

$$Z(\beta) = Z_\psi \sin\beta + Z_\chi \cos\beta \tag{2.26a}$$

$$Z'(\beta) = Z_\psi \cos\beta - Z_\chi \sin\beta, \qquad -\pi/2 < \beta \leq \pi/2 \tag{2.26b}$$

with $m_{Z(\beta)} \leq m_{Z'(\beta)}$.

Bounds were placed on $m_{Z(\beta)}$ by the neutral current analysis in Ref. 1. They ranged between about 120 GeV for $\cos\beta \simeq -0.6$ to 300 GeV for $\beta \simeq 0$ (i.e. Z_χ). In fig. 2.2, experimental bounds are given for for $Z(\beta)$ mass (called M_2) allowing for possible mixing with the ordinary Z. (θ is the $Z - Z(\beta)$ mixing angle.) Note that the constraint is not

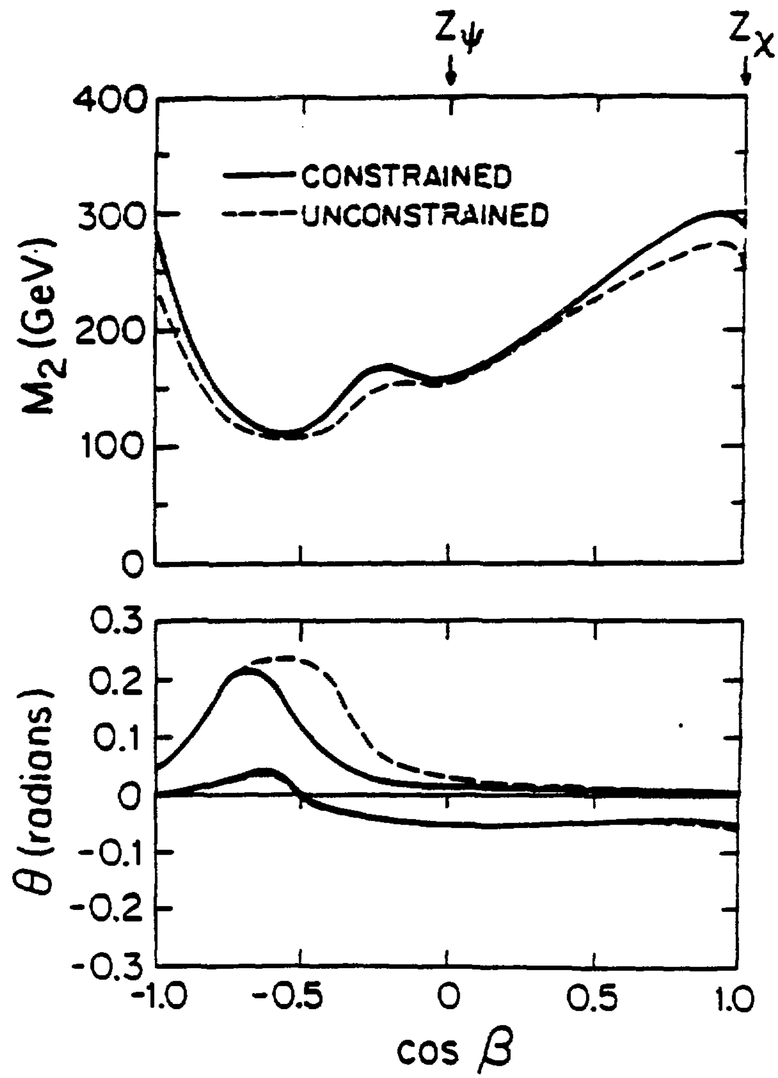

Figure 2.2: Lower limits on the mass (M_2) of an E_6 boson $Z(\beta) = Z_\Psi \sin\beta + Z_\chi \cos\beta$. The dashed line corresponds to an unconstrained Higgs mechanism, i.e. $\rho \neq 1$. Also illustrated is the allowed range of mixing, θ, between $Z(\beta)$ and the standard model Z boson.

very good for $Z(\beta)$ near the superstring inspired E_6 model $\cos\beta \simeq -0.6$. In fact, the data (in particular DESY e^+e^- annihilation results) slightly favors a $Z(\beta)$ near $\cos\beta \simeq -0.6$ which mixes with the ordinary Z. It is very important to push the bounds in fig. 2.2 into the TeV region or better yet find a Z'. To that end, the SSC will have a discovery potential for finding a Z' that should extend to 5-10 TeV.

What if a Z' is discovered? Such a discovery combined with a measurement of its couplings would almost certainly pinpoint the underlying GUT symmetry group. The absolute couplings, which could be obtained by comparing its production and decays with the standard Z, would then give us $\sqrt{\lambda_i}$ in Eq. (2.24) and thus provide further important information about coupling evolutions and new thresholds. I should note that in the E_6 scenario, the mixing angle β should be relatively easy to determine since the branching ratios

$$\frac{\Gamma\left(Z(\beta)\to f\bar{f}\right)}{\Gamma\left(Z(\beta)\to \text{ all}\right)} = \frac{\left(Q^\beta_{f_L}\right)^2 + \left(Q^\beta_{f_R}\right)^2}{\sum_f \left(Q^\beta_{f_L}\right)^2 + \left(Q^\beta_{f_R}\right)^2} \tag{2.27}$$

$$Q^\beta_f \equiv Q^\chi_f \cos\beta + Q^\psi_f \sin\beta$$

depend only on β.

In the time between now and SSC physics, it will be interesting to see if hints of a Z' boson of any kind emerge from low energy phenomenology. In that regard, atomic parity violation and νe scattering experiments may reach high enough precision to probe for such particles up to $\simeq$ 800GeV during the intervening years. If evidence for a Z' is found, an e^+e^- facility capable of sitting on that resonance will be very desirable.

The above examples were meant to illustrate the power of global fits to existing data. In the future, one would hope that significant improvements will occur in many weak neutral current experiments; thereby testing the standard model and its quantum loop corrections at a much more stringent level. Let me be specific as to what goals I see for the next generation of neutrino experiments at accelerator facilities.

In the case of $\sin^2\theta_W$, any new proposal to measure $\sin^2\theta_W$ to within $\pm1\%$ and ρ to within $\pm0.5\%$ would appear worthwhile. These goals seem to be achievable both in $\nu_\mu e$ and deep-inelastic $\nu_\mu N$ scattering. The LCD proposal addresses the $\nu_\mu e$ case. The deep-inelastic $\nu_\mu N$ reaction should either be studied at very high energies where charm threshold effects are not as important or a study of charm production must be undertaken, utilizing opposite-sign dimuons. Alternatively, a means of avoiding charm threshold effects entirely should be found. Deep-inelastic $\nu_\mu N$ played such a key role in the global analysis described above, that I would hate to see it abandoned.

Another interesting reaction to study is elastic $\overset{(-)}{\nu}_\mu p$. The standard model predicts no axial-isoscalar weak neutral currnt coupling in the lowest order SU(2) limit. Strange sea quarks and gluon loop effects, however, induce an axial-isoscalar component. Results from BNL E734[15] indicate it to be about 10-20% of the axial-isovector amplitude. By precisely measuring $d\sigma/dQ^2$, one could measure the form factor and learn something about the spin distribution and gluon content of the proton.[16] In my opinion, the axial-isoscalar neutral current plays a very special role in QCD (it is not conserved because of instanton effects); so its study by whatever means possible is strongly motivated.

3. KAON PHYSICS

During this workshop we heard very nice theoretical and experimental discussions on kaon decays by John Donoghue and Bruce Winstein. From Donoghue's overview it is clear that many interesting aspects of kaon physics remain to be explored, particularly in the realm of rare decays and CP violation. Indeed, he outlined a full research program for testing the standard model and searching for "new physics." There are already a number of ongoing dedicated rare K decay experiments at BNL, CERN, KEK and Fermilab. As an illustration, I have listed in Table 3.1 (updated from a summary talk by B. Winstein) the status of those experiments. Given the theoretical and experimental enthusiasm for kaon physics, not to mention the competitive vieing for new kaon factories by various laboratories, it is likely that we will see more initiatives in the future.

From Winstein's talk we learned that Fermilab because of its higher energy can have advantages over more intense kaon facilities. For example in these decays with a π^0 or many particles in the final state, higher energy kaon beams allow more complete reconstruction with the same angular coverage as well as better photon energy resolution. Those advantages have been very useful in the recent high statistics measurements of $K \to 2\pi$ by E731 and their attempt to determine ϵ'/ϵ.

The proposed new Fermilab Main Injector could also provide a new intense facility for rare K decay studies. With a proton energy of 120 GeV and an intensity about 100 times higher than the Tevatron, it could compete with lower energy, higher intensity kaon factories under discussion.

One particular decay that is likely to play a prominent role in the future Fermilab kaon physics program is $K_L \to \pi^0 e^+ e^-$. That rare (but not forbidden) decay has recently become very popular. It offers a potentially new clean (but still difficult) way to observe direct CP violation. John Donoghue has already described the theoretical uncertainties as well as a strategy for extracting the direct CP violating amplitude. At the risk of being somewhat redundant, I would like to emphasize the importance of searching for that decay by commenting further on the theoretical motivation and uncertainty.

The decay $K_L \to \pi^0 e^+ e^-$ can proceed through CP violating and CP conserving amplitudes. The CP violating amplitudes can be further divided into an indirect part due to a mixing $K_L = K_2 + \epsilon K_1$ (with $K_1 \to \pi^0 e^+ e^-$ mainly through the electromagnetic penguin and weak loops) and a direct CP violating $K_2 \to \pi^0 e^+ e^-$ part due to imaginary contributions of the electromagnetic penguin and weak loop amplitudes. Both of those contributions are predicted to be about the same magnitude.[18] Therefore, the situation is analogous to having $\epsilon'/\epsilon \simeq O(1)$ rather than $\sim 3 \times 10^{-3}$ as in the $K \to 2\pi$ decays. Of course, both amplitudes are also highly suppressed; so, one expects $BR\left(K_L \to \pi^0 e^+ e^-\right)$ to be in the range $10^{-11} \sim 10^{-12}$. Therefore, its measurement presents a difficult experimental challenge. To untangle those distinct sources of CP violation, one would need to measure both $BR\left(K_S \to \pi^0 e^+ e^-\right)$ and $BR\left(K_L \to \pi^0 e^+ e^-\right)$. The former is expected to be about 200 times larger than the latter, but still very hard to determine.

The CP conserving amplitude for $K_L \to \pi^0 e^+ e^-$ is dominated by the two photon intermediate state. There is some controversy[19] about the magnitude of that amplitude with estimates of its contribution to $BR\left(K_L \to \pi^0 e^+ e^-\right)$ ranging from a negligible $O\left(10^{-14}\right)$ to a dominant $O\left(10^{-11}\right)$. Straightening out that controversy is very important. To that end, experimental measurements of $K_L \to \pi^0 \gamma\gamma$, particularly the $\gamma\gamma$ mass spectrum would be very helpful. Even if the CP conserving two photon contribution to $K_L \to \pi^0 e^+ e^-$ turns out to be large, one can still try to unravel the CP violation. For example, Sehgal[19] has pointed out how in that event, interference between CP conserving and violating amplitudes could lead to a large charge asymmetry in the $e^- - e^+$ energy difference which could be used to determine the relative CP violating and conserving amplitudes.

Table 3.1: Status of rare K decay searches. Update of a summary talk table given by B. Winstein at the BNL Workshop on K Decays and CP Violation.

Decay	Experiment	Result	Expected Sensitivity	Comments
$K^+ \to \pi^+ +$ nothing	BNL E787	-	2×10^{-10}	$\sim 10^{-9}$ Next Run
$K^+ \to \pi^+ \mu^+ e^-$	BNL E777	$< 1.1 \times 10^{-9}$	1.5×10^{-10}	Will pursue $K^+ \to \pi^+ e^+ e^-$ $\pi^o \to \mu e < 8 \times 10^{-8}$
$K_L \to \mu e$ $K_L \to e^+ e^-$ $K_L \to \pi^o e^+ e^-$	BNL E780 (Completed)	$< 1.9 \times 10^{-9}$ $< 1.2 \times 10^{-9}$ $< 3.2 \times 10^{-7}$		Will pursue $K_L \to \pi^o e^+ e^-$ E845 at 10^{-10} sensitivity
$K_L \to \mu e$ $K_L \to e^+ e^-$	BNL 791	$< 3.0 \times 10^{-10}$	$\sim 2 \times 10^{-11}$	Will measure $K_L \to \mu^+ \mu^-$ May pursue $K_L \to \pi^o e^+ e^-$
$K_L \to \mu e$ $K_L \to e^+ e^-$ $K_L \to \pi^0 e^+ e^-$	KEK E137 CERN NA31	$\lesssim 3 \times 10^{-9}$ $\lesssim 4 \times 10^{-9}$ $< 4 \times 10^{-8}$	$\sim 2 \times 10^{-11}$	Will pursue $K_L \to \pi^o e^+ e^-$ at 10^{-10} E162
$K_s \to \pi^+ \pi^- \pi^o$	FNAL E621	$\lesssim 1.5 \times 10^{-7}$	$\sim 3 \times 10^{-9}$	Expected Rate 1.2×10^{-9} Upgrade
$K_L \to \pi^o e^+ e^-$	FNAL E731	$\lesssim 4.2 \times 10^{-8}$	$\sim 1 \times 10^{-8}$	May pursue at 10^{-11}

Sorting out direct CP violation in $K_L \to \pi^0 e^+ e^-$ will be quite arduous. It will likely require an ultimate branching ratio sensitivity of $O\left(10^{-12} \sim 10^{-13}\right)$. In addition, complementary studies of the decays $K_S \to \pi^0 e^+ e^-$ and $K_L \to \pi^0 \gamma^+ \gamma^-$ are likely to be necessary. That effort is, however, very worthwhile, since $K_L \to \pi^0 e^+ e^-$ is not subject to the hadronic matrix element uncertainties that plague the interpretation of $K \to 2\pi$ determinations of ϵ'/ϵ. As such, it provides a real quantitative test of direct CP violation predictions in the standard model.

The long road from the present bound $BR\left(K_L \to \pi^0 e^+ e^-\right) < 4 \times 10^{-8}$ (see Table 3.1) to a sensitivity of $O\left(10^{-12} \sim 10^{-13}\right)$ will not be easy. Approved experiments at BNL and KEK will have an anticipated sensitivity of $O\left(10^{-10}\right)$. They can be viewed as exploratory searches for "new physics" or preliminary excursions. I should note that $K_L \to \pi^0 e^+ e^-$ is also an interesting channel to search for a light Higgs or Axion-like particle. (In that regard, $K_L \to \pi^0 \mu^+ \mu^-$ is also an interesting mode to explore.)

The present fixed target Tevatron program may be able to push $BR\left(K_L \to \pi^0 e^+ e^-\right)$ to the 10^{-11} sensitivity level. If lucky, a signal might be observed. Even if that turns out not to be the case, valuable experience will be gained and the stage will be set or future high intensity K facilities where sensitivities at the level of 10^{-13} may be attained.

4. NEUTRINO MASSES, MIXING AND MOMENTS

If neutrinos have small masses, it is likely that they mix and thereby undergo oscillations. In addition, higher order quantum corrections will induce electromagnetic moments which could have interesting consequences. The positive identification of non-zero mass, mixing or moment would be very exciting. It would be a clear signal of physics beyond the standard model.

At this workshop, we heard several nice reviews of neutrino mass and oscillation phenomenology as well as new ideas for future experiments. P. Langacker gave a comprehensive discussion of neutrino mass scenarios and their experimental implications. M. Shaevitz described existing oscillation bounds and the potential for future improvements. S. Parke updated the status of solar and atmospheric neutrino experiments and their constraints on MSW matter enhanced oscillations. Those overviews were complemented by reports from accelerator based oscillation experiments. R. Seto (BNL E776), J. Dumarchez (BNL E816), and S. Freedman (LAMPF E645) discussed their experimental searches for $\nu_\mu \to \nu_e$ oscillations and the resulting bounds. New proposals and ideas for oscillation experiments at BNL, LAMPF and Fermilab were described by S. Aronson, H. White, R. Bernstein, J.D. Bjorken, and N.W. Reay. The BNL proposal envisions a long baseline $\nu_\mu \to \nu_\tau$ disappearance experiment with detectors at 1 km and 10 km. It would be sensitive to Δm^2 as low as $10^{-2} eV^2$. A nice idea to improve the search for $\nu_\mu \to \nu_\tau$ appearance in emulsion was put forward by N.W. Reay. He could conceivably push $\sin^2 2\theta$ down to $\sim 2 \times 10^{-4}$ for $\Delta m^2 \gtrsim 10 eV^2$, a significant improvement. Such a search is motivated in part by the idea that ν_μ or ν_τ with mass $\sim 10 - 50$ eV could be the galactic dark matter. The $\nu_e \to \nu_\tau$ oscillation mode could be sought at LCD or using the tagged Fermilab neutrino facility advocated by R. Bernstein. Possible future long-baseline experiments at Fermilab were discussed by J.D. Bjorken and M. Koshiba while J. Panman described possible future oscillation initiatives by the CHARM II collaboration at CERN.

In my opinion, searches for oscillations at accelerator facilities, particularly $\nu_\mu \to \nu_\tau$, are well motivated and worth doing. I would also like to advocate continued direct searches for neutrino masses and electromagnetic moments. Below, I discuss some of my reasons for that perspective.

Neutrino Masses

In the standard $SU(2)_L \times U(1)$ model, the photon and all three species of neutrinos have zero mass. For the photon, masslessness is a natural consequence of exact electromagnetic gauge invariance; its validity being well verified experimentally by the present bound[20] $m_\gamma < 3 \times 10^{-27}$ eV. However, the masslessness of neutrinos is not on such firm theoretical or experimental footing. Theoretically, $m_\nu = 0$ because only the left-handed component ν_L of each neutrino specie is employed (the right-handed component ν_R is assumed not to exist) and lepton number conservation is required. Relaxing either of these constraints can lead to $m_\nu \neq 0$. Indeed, the present experimental bounds[21,22,23]

$$m_{\nu_e} < 18 \text{ eV}, \tag{4.1a}$$

$$m_{\nu_\mu} < 0.25 \text{ MeV}, \tag{4.1b}$$

$$m_{\nu_\tau} < 35 \text{ MeV}, \tag{4.1c}$$

leave considerable room for speculation that neutrinos actually do possess mass. I might add that the experimental bounds have improved by factors of $2 \sim 5$ since I reviewed[24] this subject in 1981. Progress is slow, but we can expect further improvement by factors of ~ 3 during the next few years. A worthy goal would be a further factor of 5 to 10 reduction.

In the standard model charged leptons and quarks acquire masses by coupling their left- and right-handed components through a Higgs scalar isodoublet which obtains a vacuum expectation value. (Mass terms change chirality, so one must couple left-to-right in order to generate mass.) The fermion masses generated by the Higgs mechanism are totally arbitrary, their values are chosen to agree with experiment. If right-handed components exist, then neutrinos can also be given arbitrary masses by the standard Higgs mechanism. The right-handed components will be sterile under $SU(2)_L \times U(1)$ gauge interactions; so, they do not upset phenomenology. Of course, in such a scenario one must address the question: Why are neutrino masses so much smaller than the masses of other fermions in their generation?

Even without introducing right-handed neutrinos, the neutrino can be given a Majorana mass by coupling ν_L to its (right-handed) charge conjugation ν_L^c (which I subsequently call $\bar{\nu}$, the antineutrino). The mass term would appear as

$$\frac{1}{2}\, m_\nu \; (\bar{\nu}_L^c\, \nu_L + h.c.)\,. \tag{4.2}$$

(It requires anticommuting fields.) A Majorana mass is not possible for other fermions because it would violate electric charge conservation. In the case of the neutrino, lepton number is violated $\Delta L = 2$ by Eq. (4.2).

The simplest way of introducing a Majorana mass is to couple the neutrino-charged lepton $SU(2)_L$ doublet to a Higgs triplet.[25] The implications of such a scenario are the occurrence of neutrinoless double beta decay (so far not observed) and the existence of additional physical scalars associated with the Higgs triplet. In particular, a light scalar and pseudoscalar (Majoran) can exist in such a scenario. This mechanism can accommodate a small neutrino mass, but also fails to naturally explain small masses.

In my opinion, the most compelling rationale for small neutrino masses is found by combining the two mass generating mechanisms mentioned above in the see-saw[26] mechanism.

The basic idea is motivated by grand unified theories (GUTS) which often contain very heavy right-handed Majorana neutrinos, ν_R. Since ν_R is a singlet under $SU(2)_L \times U(1)$, it is expected that large Majorana masses M_R will arise for such fields. On the other hand, the ordinary Higgs mechanism should lead to a Dirac mass, m_f, which couples ν_R to ν_L. Of course, m_f is expected to be a typical fermion mass scale, i.e., of order m_e, m_u or m_d for the first generation. Therefore, one expects a mass matrix of the form

$$\begin{pmatrix} \overline{\nu}_L^c & \overline{\nu}_R \end{pmatrix} \begin{pmatrix} 0 & m_f \\ m_f & M_R \end{pmatrix} \begin{pmatrix} \nu_L \\ \nu_R^c \end{pmatrix} \tag{4.3}$$

Diagonalization leads to a heavy mass $O(M_R)$ Majorana state which is primarily ν_R and a light Majorana state with mass m_f^2/M_R which is the ordinary neutrino. Since $M_R >> m_f$, the induced mass is naturally very small. In such a scenario, one expects the three neutrinos to have a hierarchical mass relationship

$$m_3^2 >> m_2^2 >> m_1^2, \tag{4.4}$$

since

$$m_1^2 : m_2^2 : m_3^2 :: m_u^4 : m_c^4 : m_t^4, \tag{4.5}$$

(or $m_e^4 : m_\mu^4 : m_\tau^4$) in many scenarios. If that is the case, then one might (approximately) anticipate

$$m_1^2 : m_2^2 : m_3^2 :: 1 : 10^{10} : 10^{16}. \tag{4.6}$$

That hierarchy could have important consequences for neutrino oscillations as we shall see. Of course, at issue is what value of M_R sets the neutrino mass scale? Is it the GUT unification scale $\sim 10^{15} \sim 10^{18}$ GeV or something smaller? If the GUT scale enters, then neutrino masses are likely to be very small and extremely hard to determine.

Mixing and Oscillations

If neutrinos have mass, then one expects them to mix. In analogy with quark mixing, the weak interaction states ν_e, ν_μ and ν_τ will be related to the mass eigenstates ν_1, ν_2 and ν_3 by a unitary matrix. It is convenient to parameterize the mixing by[27]

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = \begin{pmatrix} C_1C_3 & S_1C_3 & S_3e^{-i\delta} \\ -S_1C_2 - C_1S_2S_3e^{i\delta} & C_1C_2 - S_1S_2S_3e^{i\delta} & S_2C_3 \\ S_1S_2 - C_1C_2S_3e^{i\delta} & -C_1S_2 - S_1C_2S_3e^{i\delta} & C_2C_3 \end{pmatrix} \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}$$

$$C_i \equiv \cos\theta_i, \quad S_i \equiv \sin\theta_i, \quad i = 1,\ 2,\ 3. \tag{4.7}$$

Although it is probably best to keep an open mind about neutrino mixing, if the quark analogy holds, one expects small θ_i.

Neutrino masses and mixing can lead to many interesting phenomena in laboratory experiments, astrophysics, and cosmology. For example, neutrino decay (e.g., $\nu_\mu \to \nu_e + \gamma$) may occur or neutrinos may have small electromagnetic moments. A particularly popular possibility is neutrino oscillation from one specie to another. I will briefly survey that broad subject.

If neutrinos have masses and mix, then neutrino oscillations will occur. A neutrino ν_ℓ produced at time 0 may at a later time be observed as a $\nu_{\ell'}$. To illustrate the basic features, first consider the simplified case of only two species ν_ℓ and $\nu_{\ell'}$ such that

$$\begin{pmatrix} \nu_\ell \\ \nu_{\ell'} \end{pmatrix} = \begin{pmatrix} \cos\theta & \sin\theta \\ -\sin\theta & \cos\theta \end{pmatrix} \begin{pmatrix} \nu_1 \\ \nu_2 \end{pmatrix} \tag{4.8}$$

with ν_1 and ν_2 mass eigenstates with masses m_1 and m_2 respectively. If at time $t = 0$, a $\mid \nu_\ell >=\mid \nu_1 > \cos\theta + \mid \nu_2 > \sin\theta$ is produced with momentum $\vec{p}$, then at a latter time t, it will evolve to

$$i\frac{d}{dt} \mid \nu_i >= E_i \mid \nu_i >, i = 1,\ 2$$

$$\mid \nu(t) >= e^{-iE_1 t} \mid \nu_1 > \cos\theta + e^{-iE_2 t} \mid \nu_2 > \sin\theta \tag{4.9}$$

The probability of detecting a ν_ℓ at distance $R \simeq t$ (for $\beta \simeq 1$) is therefore given by

$$P(R)_{\nu_\ell \to \nu_\ell} = \mid < \nu_\ell \mid \nu(t) > \mid^2 = 1 - \sin^2 2\theta \sin^2 \left[\frac{\Delta m_{21}^2 R}{4 \mid \vec{p} \mid} \right] \tag{4.10a}$$

while the probability of observing a ν_ℓ' is

$$P(R)_{\nu_\ell \to \nu_\ell'} = \mid < \nu_\ell' \mid \nu(t) > \mid^2 = \sin^2 2\theta \sin^2 \left[\frac{\Delta m_{21}^2 R}{4 \mid \vec{p} \mid} \right] \tag{4.10b}$$

where $\Delta m_{21}^2 \equiv m_2^2 - m_1^2$. (I have used $E_i \simeq \mid \vec{p} \mid + \frac{1}{2} m_i^2 / \mid \vec{p} \mid$.) It oscillates between ν_ℓ and ν_ℓ' with a characteristic oscillation length L_{21}

$$L_{21} = \frac{4\pi \mid \vec{p} \mid}{\Delta m_{21}^2} \tag{4.11}$$

So, low energy neutrinos are more likely to oscillate. If one believes that Δm_{ij}^2 are all small as in the see-saw mechanism, then one must indeed employ neutrinos with an extremely

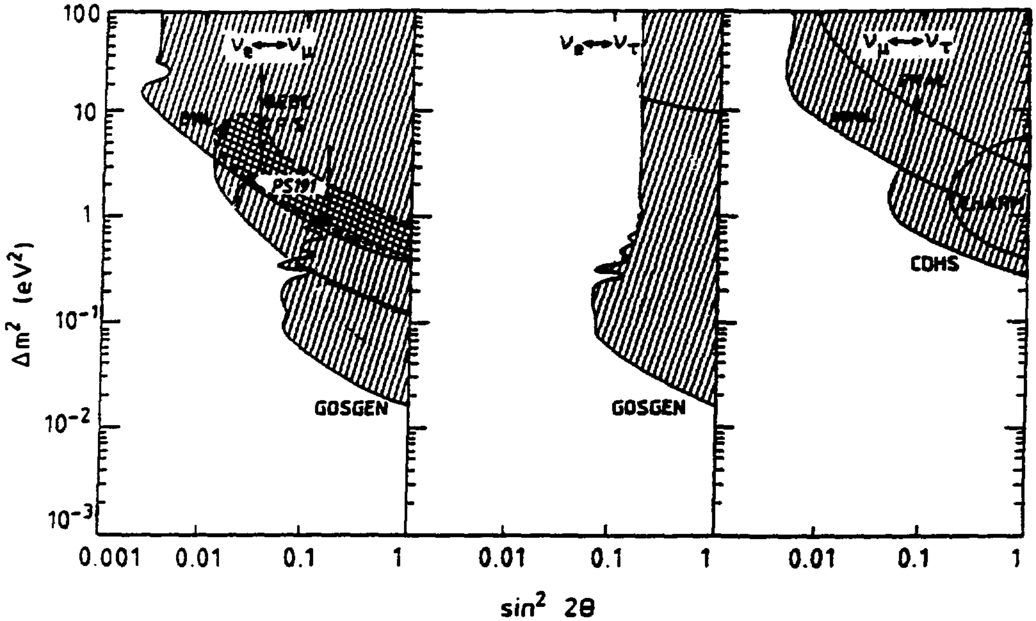

Figure 4.1: Bounds on neutrino oscillations.[28] These constraints were summarized by F. Vannucci in 1987. The only update that I have included is deletion of an oscillation signal at the Bugey reactor. Recent bounds from BNL and LAMPF on $\nu_\mu \leftrightarrow \nu_e$ oscillations are similar to the BEBC PS results.

large R/E_ν to have any hope of observing oscillations. At present, laboratory reactor and accelerator experiments looking for $\nu_e \leftrightarrow \nu_\mu$ oscillations have ruled out[29] $\Delta m^2_{21} \gtrsim 10^{-2}$ eV^2 for $\sin^2 2\theta \gtrsim 0.1$ and $\Delta m^2_{21} \gtrsim 1$ eV^2 for $\sin^2 2\theta \gtrsim 0.01$. Bounds on $\nu_\mu \to \nu_\tau$ are not as good (see M. Shaevitz's talk). To explore much smaller Δm^2_{21} requires atmospheric or astrophysical neutrino experiments in which R can be very large.

In the real world, there are (at least) three neutrinos, and their mixing can be quite complicated. A full description of neutrino oscillations in the 3 generation scenario will depend on 6 parameters, θ_1, θ_2, θ_3, δ, Δm^2_{31} and Δm^2_{21}). However, things are considerably simplified in the hierarchy case $m_3^2 >> m_2^2,\ m_1^2$. Then oscillations effectively become a 2x2 problem with "effective" mixing parameters. The relevant Δm^2 will depend on R/E_ν. To illustrate how decoupling occurs, consider the case $R/E_\nu \sim 0\left(1/\Delta m^2_{31}\right) << 1/\Delta m^2_{21}$ which may be most relevant for terrestrial experiments. In that case, the oscillation probabilities are simply governed by[8]

$$P\left(R\right)_{\nu_\mu \leftrightarrow \nu_e} \simeq \sin^2\theta_2 \sin^2 2\theta_3 \sin^2\left[\frac{\Delta m^2_{31} R}{4E_\nu}\right] \tag{4.12a}$$

$$P(R)_{\nu_\mu \leftrightarrow \nu_\tau} \simeq \sin^2 2\theta_2 \cos^4 \theta_3 \sin^2 \left[\frac{\Delta m_{31}^2 R}{4E_\nu}\right] \tag{4.12b}$$

$$P(R)_{\nu_e \leftrightarrow \nu_e} \simeq 1 - \sin^2 2\theta_3 \sin^2 \left[\frac{\Delta m_{31}^2 R}{4E_\nu}\right] \tag{4.12c}$$

The oscillation length L_{21} is too long to effect the evolution. For small mixing angles, one finds

$$P(R)_{\nu_\mu \leftrightarrow \nu_\tau} \simeq P(R)_{\nu_\mu \leftrightarrow \nu_e} / \sin^2 \theta_3 \tag{4.13}$$

This suggests that accelerators which produce ν_μ beams may be much better suited to search for $\nu_\mu \rightarrow \nu_\tau$ oscillations rather than $\nu_\mu \rightarrow \nu_e$. So far, most experiments have concentrated on $\nu_\mu \rightarrow \nu_e$.

For the other extreme $R/E_\nu \sim 0\left(1/\Delta m_{21}^2\right) >> 1/\Delta m_{31}^2$, (very long distances or small energies) oscillations are governed by Δm_{21}^2 in an effective 2 neutrino system. So, for example one finds

$$P(R)_{\nu_e \rightarrow \nu_e} = \sin^4 \theta_3 + \cos^4 \theta_3 \left(1 - \sin^2 2\theta_1 \sin^2 \frac{\Delta m_{21}^2 R}{4E_\nu}\right) \tag{4.14}$$

It is interesting to compare this formula with the observed flux of solar neutrinos where a long standing problem exists. One expects neutrinos produced in the solar core to give rise to $\nu_e +^{37} Cl \rightarrow e^- +^{37} Ar$ at the rate of[29]

$$7.9 \pm 2.5 SNU \qquad \left(1SNU = 10^{-36} \text{captures/atom} - \text{s}\right) \tag{4.15}$$

However, Davis and collaborators have observed for the last 17 years an average flux of[30]

$$2.2 \pm 0.3\ SNU \tag{4.16}$$

The discrepancy constitutes the solar neutrino problem. What happened to the missing flux? A simple solution is that it was never present, i.e., the theoretical prediction in (4.15) is wrong. More exotic is the possibility that the ν_e flux was depleted by oscillations. That could occur if Δm_{21}^2 is fine-tuned to the earth-sun distance. More naturally for $\Delta m_{21}^2 >> R/E_\nu$, one finds from (4.15) a survival probability average

$$P(R)_{\nu_e \rightarrow \nu_e} = \sin^4 \theta_3 + \cos^4 \theta_3 \left(\cos^4 \theta_1 + \sin^4 \theta_1\right) \tag{4.17}$$

which can be quite small for particular mixing angles. The smallest value of (4.17) is $1/3$ for $\sin^2 \theta_1 = 1/2$, $\sin^2 \theta_3 = 1/3$. Such a depletion would solve the solar neutrino problem;

but it still appears somewhat contrived and requires relatively large mixing which runs counter to our experience with quark mixing and theoretical prejudices.

A truly ingenious solution to the solar neutrino puzzle has been proposed by Mikheyev and Smirnov.[31] Employing an analysis of matter effects on neutrino oscillations by Wolfenstein,[13] they showed that for a large range of neutrino masses and mixing parameters, neutrino oscillations between ν_e and ν_μ or ν_τ in the sun's interior could be significantly enhanced. In fact, for some neutrino energies one can get a nearly complete transformation of ν_e into another neutrino flavor. That scenario (referred to as the MSW effect) provides an elegant natural solution to the solar neutrino puzzle and hence has become very popular. It enters into the planning for future solar neutrino experiments as well as their interpretation. I describe in section 4 some of its salient features and potential implications. I might note that for more than a year Ray Davis[30] has been averaging $4 \sim 5$ SNU. One does not know whether that represents a statistical fluctuation or an important hint regarding resolution of the solar neutrino puzzle.

Electromagnetic Moments

Although neutrinos are electrically neutral, they can have electromagnetic form factors. In the case of 4 component Dirac neutrino[33]

$$\begin{aligned} < \nu \mid J_\mu^{em} \mid \nu > = \overline{u}_\nu [(F(q^2) + \gamma_5\ G(q^2))\gamma^\alpha (g_{\alpha\mu} - \frac{q_\alpha q_\mu}{q^2}) \\ + (M(q^2) + \gamma_5 E(q^2)) i\sigma_{\mu\alpha} q^\alpha] u_\nu \end{aligned} \tag{4.18}$$

The F and G form factors are charge radii while $M\,(0)$ is the magnetic dipole moment and $E\,(0)$ is the electric dipole moment of the neutrino. In the case of Majorana neutrinos (which are self-conjugate), $F = M = E = 0$, i.e., only $G\,(q^2)$ can be non-vanishing.

Both Dirac and Majorana neutrinos can also have transition moments that are off diagonal, i.e., connect mass eigenstates ν_2 and ν_1. Magnetic and electric transition moments can give rise to neutrino decay or flavor precession as we shall see.

I will say a few words about the neutrino electromagnetic and transition moments. First consider a transition moment $\kappa_{21}\ e/2m_e$ between ν_2 and ν_1 such that the decay amplitude for $\nu_2 \to \nu_1 + \gamma$ is given by[15]

$$\mathcal{M} = \kappa_{21} \frac{e}{2m_e} \mathcal{E}^\mu q^\nu \overline{\nu}_1 \sigma_{\mu\nu} (1 - \gamma_5)\, \nu_2 \tag{4.19}$$

In that case, the decay rate for $\nu_{2_L} \to \nu_{1_R} + \gamma$ is

$$\Gamma\left(\nu_{2_L} \to \nu_{1_R} + \gamma\right) = \frac{\alpha}{2m_e^2}\left(\frac{m_2^2 - m_1^2}{m_2}\right)^3 \kappa_{21}^2 \tag{4.20}$$

Of course, for κ_{21} and $m_2 - m_1$, very small, as expected, the predicted decay rate will be tiny. Indeed, it can easily exceed the age of the universe.

If the neutrino is 4 component (Dirac), it can have a (diagonal) magnetic and/or electric dipole moment. Of course, the electric dipole moment would violate T and would most likely be the smaller. In the standard $SU(2)_L \times U(1)$ with massive Dirac neutrinos, one finds that one loop induced magnetic dipole moment $\mu = \kappa e/2m_e$ is[34,35]

$$|\mu| = \frac{3eG_\mu m_\nu}{8\sqrt{2}\pi^2} \leq 3 \times 10^{-19}\left(m_\nu/1\text{eV}\right) e/2m_e \tag{4.21}$$

which is very small. (Transition moments are generally even smaller since they involve flavor mixing.) One, therefore, expects $\kappa \lesssim 10^{-19}$ for very light neutrinos, unless some new chiral changing physics enters.

Magnetic, electric, and transition moments can be bounded by laboratory experiments and/or astrophysics experiments. In the case of $\nu - e$ scattering, the existence of any such moment would increase the cross- section by (neglecting neutrino mass)[36]

$$\Delta\frac{d\sigma(\nu e)}{dy} = \frac{\pi\alpha^2\kappa^2}{m_e^2}\left(\frac{1}{y} - 1\right) \tag{4.22}$$

where $y = \left(E_e' - m_e\right)/E_\nu$. Note that it is rather insensitive to the initial ν energy. Therefore, low energy experiments where the usual cross-section $\sim E_\nu$ are most sensitive to κ. From existing experimental data, one finds[37,38]

$$|\kappa_{\nu_\mu}| < 10^{-9} \tag{4.23}$$

$$|\kappa_{\nu_e}| < 4 \times 10^{-10} \tag{4.24}$$

(Those bounds also apply to electric and transition dipole moments.) It will be difficult to push those direct experimental bounds much farther, since the effect in (4.22) goes like κ^2. Better bounds,[39]

$$|\kappa_\nu| < 8.5 \times 10^{-11} \tag{4.25}$$

are obtainable from stellar evolution arguments. Recently, Fukugita and Yazaki[40] have argued that such bounds can be extended to $\kappa \lesssim 10^{-11}$ using observed data on Red Giants.

A value of $\kappa \gtrsim 10^{-11}$ would have important implications for solar neutrinos, as we shall see in the next section. Therefore, laboratory experiments should also aim to achieve that level of sensitivity. I might note that rather conservative analyses[41] of supernova dynamics combined with the observation of expected neutrino flux from 1987a give the bound $\kappa < 10^{-12}$. That bound does not apply to Majorana transition moments.

Matter Effects

To understand the effect of matter on neutrino propagation, consider the amplitude for low energy neutrino, ν_ℓ, scattering off a fermion f.

$$\mathcal{M}\left(\nu_\ell f \to \nu_\ell f\right) = -i\frac{G_\mu}{\sqrt{2}}\,\overline{\nu}_\ell\,\gamma^\alpha\left(1-\gamma_5\right)\nu_\ell\overline{f}\gamma_\alpha\left(C^V_{\nu_\ell f} + C^A_{\nu_\ell f}\gamma_5\right)f \tag{4.26}$$

where C^V and C^A are constants and $G_\mu = 1.16636 \times 10^{-5}\ GeV^{-2}$. In the case of an unpolarized medium, the coherent forward scattering of neutrinos with momentum p_ν can be described by an index of refraction n_{ν_ℓ} given by[42]

$$p_\nu\left(n_{\nu_\ell} - 1\right) = -\sqrt{2}\,G_\mu \sum_{f=e,u,d} C^V_{\nu_\ell f}N_f \tag{4.27}$$

where the N_f is the fermion number density and

$$C^V_{\nu_\ell f} = T_{3f} - 2Q_f\sin^2\theta_W \qquad f \neq \ell \tag{4.28a}$$

$$C^V_{\nu_\ell \ell} = 1 + T_{3_\ell} - 2Q_\ell\sin^2\theta_W \tag{4.28b}$$

with $T_{3f} = \pm\frac{1}{2}$ the weak isospin. (Quantum loop corrections to (4.28) are given in ref. 43). Note that $\nu_e e$ scattering differs from $\nu_\mu e$ and $\nu_\tau e$ scattering because of an additional charged current component in the amplitude. That effect leads to an important difference in indices of refraction

$$-p_\nu\left(n_{\nu_e} - n_{\nu_\mu}\right) = -p_\nu\left(n_{\nu_e} - n_{\nu_\tau}\right) = \sqrt{2}\,G_\mu N_e \tag{4.29}$$

which can significantly modify oscillations. An alternative description of the difference in (4.29) can be given by writing the amplitude in (26) as an effective Lagrangian

$$\mathcal{L} = -\sqrt{2}\,G_\mu\,\overline{\nu}_\ell\,\gamma^\alpha\,\frac{1-\gamma_5}{2}\,\nu_\ell\,\overline{f}\,\gamma_\alpha\left(C^V_{\nu_\ell f} + C^A_{\nu_\ell f}\gamma_5\right)f \tag{4.30}$$

We can average the Lagrangian over the background matter medium using

$$< C^V_{\nu_\ell f}\overline{f}\gamma_0 f >= C^V_{\nu_\ell f}N_f \tag{4.31}$$

whereas all other terms average to zero. One can, therefore, interpret the medium as providing an external potential

$$V = \sqrt{2}\, G_\mu \sum_f C^V_{\nu_e f}\, N_f \tag{4.32}$$

felt by neutrinos. With such a potential, neutrino propagation is described by replacing

$$i\frac{d}{dt} \to i\frac{d}{dt} - V \tag{4.33}$$

in the equation of motion. (For antineutrinos $V_{\bar{\nu}} = -V_\nu$.)

Althought the magnitude of the potential is small

$$\sim 4 \times 10^{-14}\ \mathrm{eV}\left(N_f/6 \times 10^{23} cm^{-3}\right),$$

it can have truly remarkable consequences when it interferes with an equally small vacuum energy difference $\Delta m^2/2E_\nu$. Several examples will be illustrated below.

MSW Oscillations

Only the difference of potentials affects neutrino oscillations in matter. In ordinary matter that difference is due to an additional repulsive interaction for the ν_e with electrons in the medium coming from charged current interactions. In that way, the coupled equations of motion in matter become (neglecting quantum corrections)[43] up to common diagonal terms

$$\begin{aligned} i\frac{d}{dt}\begin{pmatrix}\nu_e\\ \nu_\mu\\ \nu_\tau\end{pmatrix} = &\; U\begin{pmatrix}0 & 0 & 0\\ 0 & \Delta m^2_{21}/2p_\nu & 0\\ 0 & 0 & \Delta m^2_{31}/2p_\nu\end{pmatrix}U^{-1}\\ &+\begin{pmatrix}\sqrt{2}G_\mu N_e & 0 & 0\\ 0 & 0 & 0\\ 0 & 0 & 0\end{pmatrix}\begin{pmatrix}\nu_e\\ \nu_\mu\\ \nu_\tau\end{pmatrix}\end{aligned} \tag{4.34}$$

where U is the unitary matrix in (4.7).

To illustrate the MSW effect, consider the two generation limit, i.e., $\nu_e - \nu_\mu$ oscillations with $S_2 = S_3 = 0$ in (4.7). For that case (34) becomes

$$i\frac{d}{dt}\begin{pmatrix}\nu_e\\ \nu_\mu\end{pmatrix} = \begin{pmatrix}\sqrt{2}\, G_\mu\, N_e & \frac{-\Delta m^2_{21}}{4p_\nu}\sin 2\theta\\ -\frac{\Delta m^2_{21}}{4p_\nu}\sin 2\theta & \frac{\Delta m^2_{21}}{2p_\nu}\cos 2\theta\end{pmatrix}\begin{pmatrix}\nu_e\\ \nu_\mu\end{pmatrix} \tag{4.35}$$

As we saw before, in a vacuum with $N_e = 0$, $\nu_e \to \nu_\mu$ oscillations are governed by (see (4.10))

$$P(R)_{\nu_e \to \nu_\mu} = \sin^2 2\theta \sin^2 [\pi R / L_{12}] \tag{4.36a}$$

with

$$L_{21} = 4\pi p_\nu / \Delta m_{21}^2 \tag{4.36b}$$

the vacuum oscillation length. Solving (4.35) for constant N_e, one finds the modified matter oscillation result

$$P(R)_{\nu_e \to \nu_\mu} = \sin^2 2\theta_m \ \sin^2 \left(\frac{\pi R}{L_m} \right) \tag{4.37a}$$

where

$$\sin^2 2\theta_m = \frac{\sin^2 2\theta}{1 - 2\frac{L_{21}}{L_0} \cos 2\theta + \left(\frac{L_{21}}{L_0}\right)^2} \tag{4.37b}$$

$$L_m = \frac{L_{21}}{\left[1 - 2\frac{L_{21}}{L_o} \cos 2\theta + \left(\frac{L_{21}}{L_o}\right)^2\right]^{1/2}} \tag{4.37c}$$

$$L_o = \frac{2\pi}{\sqrt{2}\, G_\mu\, N_e} \tag{4.37d}$$

so, we see that the matter mixing angle θ_m can be large even if the vacuum angle θ is small. In fact, if the condition

$$L_{21}/L_o = \cos 2\theta \quad \text{(Resonance Condition)} \tag{4.38}$$

which corresponds to an electron density

$$N_e = \frac{\Delta m_{21}^2 \cos 2\theta}{2\sqrt{2}\, G_\mu\, p_\nu} \tag{4.39}$$

is satisfied, one finds $\theta_m = 45^o$, maximal mixing. Of course, with the full 3 generation mixing matrix, the situation is more complicated but the physics is basically the same. The potential in matter can cause energy level crossings and thereby enhance oscillations. (It can also, of course, quench oscillation if the splitting is increased by the matter potential.)

The MSW effect actually comes about from a changing electron density profile in the sun. Neutrinos produced in the solar core start out at a high electron density, but experience a decreasing density as they propagate outward (i.e., the potential changes).

If the resonance condition in (4.39) is satisfied during transit, ν_e can resonate completely into ν_μ (or ν_τ). Assuming that the vacuum angle θ is small, they will not convert back to ν_e as they propagate to the earth.

The MSW solution to the solar neutrino problem is very elegant and naturally attractive. It can resolve the longstanding low neutrino flux observations for Δm^2 in the range $10^{-7} \sim 10^{-4}$ eV^2 and for a large range of mixing angles $0.001 < \sin^2\theta \leq 0.4$. It has also inspired many other applications of the effect of matter on neutrinos Two applications, neutrino decay and spin-flavor precession will be briefly surveyed.

Neutrino Decay

If neutrinos have mass and mix, then the heavier neutrino species should decay into lighter ones. Some vacuum decays such as $\nu_{2_L} \rightarrow \nu_{1_R} + \gamma$ (see (4.20)) will, however, be significantly altered in dense matter because ν_{2_L} and ν_{1_R} experience different potentials. Even more important, in a hot dense plasma, photons acquire an effective mass, the plasma frequency. That mass will inhibit and can quite easily eliminate the possibility of radiative neutrino decay in matter.

A more interesting scenario considered by Berezhiani and Vsotsky[45] involves the decays $\nu_e \rightarrow \overline{\nu}_e + M$ (or $\overline{\nu}_e \rightarrow \nu_e + M$ in neutron-rich matter) and $\overline{\nu}_\mu \rightarrow \nu_\mu + M$ in matter with M the hypothetical spin-0 Majoran previously mentioned. Such decays are forbidden in the vacuum because CPT requires ν and $\overline{\nu}$ to have the same mass. However, ordinary neutral matter with $N_e > N_n$ provides a repulsive potential for ν_e and $\overline{\nu}_\mu$ and attractive potential for $\overline{\nu}_e$ and ν_μ (see (4.28)). Those potentials give rise to effective mass differences that govern the decay rate. The rates for such decays in the earth, sun and supernova were estimated by Berezhiani and Vysotsky using experimental bounds on the Majoran couplings to neutrinos. The most interesting possibility seems to be in a supernova. There the neutron-rich core region can lead to $\overline{\nu}_e \rightarrow \nu_e + M$, $\overline{\nu}_\mu \rightarrow \nu_\mu + M$, $\overline{\nu}_\tau \rightarrow \nu_\tau + M$ and thus skew the flux toward a higher neutrino and depleted antineutrino flux rate. Such an effect would be energy dependent and could be affected by flavor mixing.

Neutrino Spin-Flavor Precession

One speculative solution to the solar neutrino flux problem was suggested long ago (1971) by Cisneros.[46] He argued that if neutrinos had a magnetic moment $\kappa \sim 10^{-14}$ and large transverse magnetic fields $\sim 10^6 G$ existed in the sun, then neutrinos could undergo spin precession $\nu_{e_L} \rightarrow \nu_{e_R}$ as they traversed the solar radius. Since ν_{e_R} is sterile, one would observe an effective depletion of neutrino flux by $\sim 1/2$ in accord (approximately) with the Davis experiment. That scenario has been revived and improved by Okun, Voloshin and Vysotsky (OVV).[47] Their motivation came from the observation that there appears to be an anticorrelation between neutrino flux and variations in sunspot activity, i.e., large magnetic field disturbances.[47] Indeed, there is mounting evidence that solar neutrino flux may follow the 11-year solar cycles.

The precession scenario for solar neutrinos has been extensively studied by OVV. They noted that magnetic or electric dipole moments could give $\nu_{e_L} \rightarrow \nu_{e_R}$ precession, while flavor transition moments could result in the combined spin-flavor precession $\nu_{e_L} \rightarrow \nu_{\mu_R}$ or ν_{τ_R} (for Dirac neutrinos) or $\nu_e \rightarrow \overline{\nu}_\mu$ or $\overline{\nu}_\tau$ (for Majorana neutrinos). In all cases, however, they concluded that a relatively large moment $\kappa \sim 10^{-10} \sim 10^{-11}$ was necessary for realistic solar magnetic fields $10^3 \sim 10^4 G$. Such a scenario renders $\kappa \lesssim 10^{-11}$ an interesting domain for direct experimental searches.

I would like to conclude by describing work done in collaboration with C-S. Lim[48] on the effect of matter on neutrino spin-flavor precision (similar observations have been made by E. Akhmedov).[49] We noted that in a vacuum, neutrino mass differences can quench spin flavor precession. The different potentials felt by distinct flavor neutrinos in matter is also capable of quenching precession. However, when both a mass difference and potential difference exist, they can cancel one another. In that case, a resonance situation can arise and $\nu_L \rightarrow \nu'_R$ can proceed unimpeded (provided magnetic fields are present). For a varying density, one can even have an MSW effect in which essentially complete conversion can occur. Let me illustrate how that could come about for Majorana neutrinos (the more realistic case) by considering $\nu_e \rightarrow \overline{\nu}_\mu$ spin-flavor precession. In that case, (ignoring other species and components), the precession evolution in a strong magnetic field B is governed by[29]

$$i\frac{d}{dt}\begin{pmatrix}\nu_e\\ \overline{\nu}_\mu\end{pmatrix}=\begin{pmatrix}V_{\nu_e} & \mu^* B\\ \mu B & \frac{\Delta m_{21}^2}{2E_\nu}\cos 2\theta - V_{\nu_\mu}\end{pmatrix}\begin{pmatrix}\nu_e\\ \overline{\nu}_\mu\end{pmatrix} \tag{4.40}$$

where θ is the mixing angle, and V_{ν_e}, V_{ν_μ} are matter potentials and μ is a transition moment.

For a neutral medium $N_e = N_p$, one finds (N_n = neutron density)

$$V_{\nu_e} = \frac{G_\mu}{\sqrt{2}}\,(2N_e - N_n)$$

$$V_{\nu_\mu} = \frac{G_\mu}{\sqrt{2}}(-N_n) \tag{4.41}$$

Assuming that B and the number densities are static, one finds that the transition probability for spin-flavor precession is given by

$$P(t)_{\nu_e\to\overline{\nu}_\mu} = \frac{(2\mu B)^2}{\Delta^2 + (2\mu B)^2}\sin^2\left(\sqrt{\Delta^2 + 4\mu^2 B^2}\,t/2\right) \tag{4.42}$$

$$\Delta = \sqrt{2}G_\mu(N_e - N_n) - \frac{\Delta m_{21}^2}{2E_\nu}\cos 2\theta \tag{4.43}$$

For the special condition $\Delta = 0$, spin-flavor precession is unimpeded and (4.42) becomes

$$P(t)_{\nu_e\to\overline{\nu}_\mu} = \sin^2(\mu B t) \tag{4.44}$$

i.e., maximal spin-flavor precession.

That condition is achieved when the density $N_e - N_n$ just allows a cancellation in (4.43). Of course, for an adiabatically changing density profile, it is possible to start with $\Delta > 0$ and encounter a resonance region, $\Delta \simeq 0$, where essentially all ν_e of a given energy are converted to $\overline{\nu}_\mu$ and then reach a region where either $\Delta < 0$ or B is no longer large enough to give a significant precession. In that case essentially total conversion of $\nu_e \to \overline{\nu}_\mu$ can occur.

For the above resonant spin-flavor precession scenario to proceed, one needs non-zero flavor transition moments, large magnetic fields and dense matter. In the sun, where $< B > \simeq 10^3 \sim 10^4 G$ is expected, the transition moment κ must be relatively large $\sim 10^{-10} - 10^{-11}$ for large spin-flavor precession of neutrinos with $E_\nu \sim 10$ MeV to occur. Nevertheless, it will be interesting to see whether the hints of a time variation in the neutrino flux correlated with magnetic field fluctuations is confirmed by subsequent experiments. A more likely candidate to look for spin-flavor precession is a supernova. In that case, very large $B \sim 10^{12} \sim 10^{15} G$ can occur. Hence, one is sensitive to $\kappa \sim 10^{-19} \sim 10^{-23}$ a realistic range of transition moments. The signature of such a phenomenon will be an interchange of ν_e and $\bar{\nu}_\mu$ (or $\bar{\nu}_\tau$) supernova spectra. In addition, for $N_n \gtrsim N_e$ one may have $\bar{\nu}_e \to \nu_\mu$ or ν_τ near the core. Observing such an effect would be extraordinary. Will we be so lucky to have one occur close enough to observe and will our detectors be ready if one occurs? I hope so.

5. OUTLOOK

Neutrino physics has entered an interesting era. After many years of struggle, high precision experiments at accelerator facilities are possible, but require major efforts. There are several ideas for measuring $\sin^2\theta_W$ in $\nu_\mu e$ and deep-inelastic $\nu_\mu N$ scattering with 1% accuracy. They range from the LCD detector at LAMPF and contemplated muon storage ring to a new generation of deep-inelastic scattering experiments. In the latter case, one must overcome charm threshold uncertainties either by going to very high energies or by studying charm production in more detail. In my opinion, any experiment capable of measuring $\sin^2\theta_W$ to $\pm 1\%$ is worthy of very serious consideration.

Another issue that is starting to become popular involves elastic $\overset{(-)}{\nu}_\mu p$ scattering. By measuring the induced axial-isoscalar neural current contribution, one is observing the gluon and strange content of the proton. A dedicated experiment to map out the q^2 dependence of those effects is very worthwhile and should be seriously examined.

At the more exotic level, searches for neutrino oscillations are at an exciting stage. The Kamiokande collaboration has started to confirm the solar neutrino flux measurement of R. Davis. Several new solar neutrino experiments, which should clarify things, are starting up. If the elegant MSW effect is responsible for the depletion of solar flux, it suggests a Δm^2_{21} of $10^{-4} \sim 10^{-7} eV^2$. Although that domain is beyond accelerator capabilities, it could easily correspond to $\Delta m^2_{32} \simeq 10^{-2} \sim 10^2$ eV^2 in $\nu_\mu \to \nu_\tau$ oscillation phenomena. One should, therefore, search for $\nu_\mu - \nu_\tau$ oscillations as far as possible. Disappearance experiments at lower energies can probe the small Δm^2_{32} region, but only at relatively large mixing. Higher energy facilities, such as at Fermilab, are capable of $\nu_\mu - \nu_\tau$ appearance experiments. Once above τ threshold, the ν_τ could produce τ leptons. That allows one to explore very small mixing but not so small Δm_{32}^2. Detecting a τ would not only be evidence for oscillations, but would also confirm the ν_τ's existence.

Neutrino physics is very hard, but the discovery potential is high and well worth the effort.

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