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# THEORETICAL SURVEY OF ELECTRON-POSITRON PHYSICS*

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## ABSTRACT

The success as well as the incompleteness of the standard model is recalled. We survey some of what has been checked and review the status of some of the parameters which are undetermined within the standard model. Possible areas where there may be problems are indicated and searches for phenomena which come from beyond the standard model are emphasized.

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# 1. INTRODUCTION

Everyone is aware of the success of the standard model. But we are also not satisfied with it as a final theory of Nature. It is incomplete. Many parameters, such as quark masses and weak mixing angles are arbitrary, and simply put in from outside the standard model. Nor is the reason for quark and lepton families or generations explained, let alone a connection between quarks and leptons. And then we have the hierarchy problem—how is a scale as "small" as the weak scale determined if we start at a grand-unified scale or at the Planck scale?

So we are constantly probing; checking and then rechecking with more accuracy the predictions of the standard model. We are looking for a discrepancy, a crack which will provide us with an insight as to what lies beyond the standard model.

With this in mind this survey will be somewhat larger in scope than just electron-positron physics. We will review first some of what has been checked of the standard model's assumptions or predictions. Then we will review some of the parameters which are inputs to the standard model and in particular the latest values for the magnitudes of the Kobayashi-Maskawa matrix elements. Then we proceed to two areas where there may be problems: CP violation and tau lepton decay. Finally we examine a couple of places in which new phenomena could show up: heavy neutral leptons and extra vector bosons associated with a larger gauge group.

## 2. STANDARD MODEL PARAMETERS

The standard model has been subjected to extensive reviews.[1] We will not attempt another comprehensive review here, but instead will concentrate on a few salient points and new analyses.

There are a number of very impressive measurements over the past few years that make us quite confident about the relevance of the standard model as a description of Nature. None is so direct as the discovery[2] of the $W$ and then[3] of the $Z$ by the UA1 and UA2 collaborations at CERN. Their masses are in accord with theoretical predictions. Moreover, if the masses are converted to a value of $\sin^2\theta_W$ and compared with other determinations from neutral current experiments in both purely leptonic (such as $e^+e^- \to \mu^+\mu^-$ and neutrino-electron elastic scattering) and semi-leptonic (such as deep inelastic neutrino-nucleon and polarized electron-deuteron inelastic scattering) processes, there is very impressive agreement among the various experiments.[4] These results are not yet quite at the level of testing the electroweak radiative corrections at the one loop level, but the time for doing so is coming soon.

The axial-vector couplings of the $\mu$ and $\tau$ leptons and $c$ and $b$ quarks to the $Z$ have been determined from the front-back asymmetry in electron-positron annihilation. The lepton couplings are now determined to roughly 10% from the combined PEP and PETRA measurements[5] shown in Figure 1 and are in agreement with the standard model within two standard deviations. Similar results hold for the $c$ and $b$ quarks, with larger errors. Said in a different form, if our choice of the difference of weak isospins, $I_{3R} - I_{3L}$ (which is proportional to $g_A$), is restricted to half-integral values, then the data unambiguously favor the values which result if one makes the standard model assignments of weak isospin for these quarks and leptons.

Essentially the same measurements put limits on an effective four-Fermi interaction arising from compositeness.[6] Writing the coefficient of this interaction as $g^2/\Lambda^2$, and taking $g^2/4\pi = 1$, the recent measurements[7] of $e^+e^- \to e^+e^-$ and $e^+e^- \to \mu^+\mu^-$ put a lower limit on $\Lambda$ of about 2 TeV.

One of the sets of parameters which we are unable to predict from inside the standard model is the relation between the mass eigenstates and weak eigenstates for quarks. The relation between them is expressed by a matrix known as the Kobayashi-Maskawa[8] matrix, as they exhibited an explicit parametrization in the six quark case in 1973. We define this unitary matrix as

$$\begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix} \tag{1}$$

where we have followed the usual convention of leaving the charge $2e/3$ quarks unmixed, and expressed the relation by a rotation acting on the mass eigenstates to obtain the weak eigenstates of the charge $-e/3$ quarks.

Since the matrix elements come from outside the standard model, as do the quark masses, they are a potential window on the physics beyond: best of all, if we knew what to do with this information, it is something which we have in our hands today. At a somewhat less ambitious level, there may be a connection between the quark masses and the matrix elements,[9] at a minimum reducing the number of free parameters, and perhaps giving us insight into a symmetry or dynamics which relates one generation to another.

In addition, knowledge of these matrix elements is useful for more mundane engineering purposes in calculating the effects of loops involving virtual heavy quarks, as in the next section. Last, but not necessarily so in importance, checking on unitarity of the matrix assuming there are three generations could, by its failure, provide evidence for a fourth generation. (We immediately note that there is no evidence for this at the present time. See below.)

The values of each of the matrix elements can in principle be determined from weak decays of the relevant quarks, or, in some cases, from deep inelastic neutrino scattering. Our knowledge of the various matrix elements at this time comes from the following sources[10]

(1) $|V_{ud}| = 0.9729 \pm 0.0012$ from nucleon beta decay when compared to muon decay, after incorporating newly refined corrections from Marciano and Sirlin[11] for the order $\alpha$ structure dependent terms and leading log radiative corrections, which are summed using the renormalization group.

(2) $|V_{us}| = 0.221 \pm 0.002$ from $K_{e3}$ and hyperon decays, with a new analysis by Leutwyler and Roos[12] which takes account of isospin violation between charged and neutral $K$ decays and brings the values extracted for $|V_{us}|$ from the two decays into agreement at the 1% level.

(3) $|V_{cd}| = 0.24 \pm 0.03$ from neutrino production of charm[13] and subsequent semileptonic decays.

(4) $|V_{cs}| > 0.66$ from the width for $D \to Ke\nu$ and the (very conservative) assumption that the relevant form factor occurring in $D_{e3}$ decay, $f_+^D(0) < 1$.

(5) $|V_{ub}/V_{cb}| < 0.19$ from the newer (and less stringent) limit on $b \to u/b \to c$ of $< 0.08$.[14]

(6) $0.037 < |V_{cb}| < 0.053$ from the present world average $b$ lifetime.[14]

Putting these constraints together with unitarity and the assumption of three generations gives the following limits on the Kobayashi-Maskawa matrix elements:

$$\begin{pmatrix} 0.9742 \text{ to } 0.9756 & 0.219 \text{ to } 0.225 & 0 \text{ to } 0.008 \\ 0.219 \text{ to } 0.225 & 0.973 \text{ to } 0.975 & 0.037 \text{ to } 0.053 \\ 0.002 \text{ to } 0.018 & 0.036 \text{ to } 0.052 & 0.9986 \text{ to } 0.9993 \end{pmatrix} . \qquad (2)$$

The data are consistent with there being just three generations: neither do they preclude there being more than three generations. If we assume there are four or more generations, then the constraints from unitarity are of course much less effective:

$$\begin{pmatrix} 0.9710 \text{ to } 0.9748 & 0.218 \text{ to } 0.224 & 0 \text{ to } 0.01 & \cdots \\ 0.192 \text{ to } 0.288 & 0.66 \text{ to } 0.98 & 0.037 \text{ to } 0.053 & \cdots \\ 0 \text{ to } 0.14 & 0 \text{ to } 0.72 & 0 \text{ to } 0.999 & \cdots \\ \vdots & \vdots & \vdots & \end{pmatrix} . \qquad (3)$$

The known matrix elements may be used together with unitarity to restrict potential matrix elements to additional quarks; for example, we must have $|V_{ub'}| < 0.088$ from the known elements in the first row. The magnitudes of the matrix elements given above may be used to obtain the angles involved in any particular parametrization, such as that of Kobayashi-Maskawa[8] or that of Maiani.[15]

## 3. CP VIOLATION

More than thirty years after its discovery, the only place where CP violation has been definitively observed is still the neutral $K$ system, and in particular, in the decay $K \to \pi\pi$. This amplitude may be split into parts, $A_0$ and $A_2$, that involve final $\pi\pi$ isospin 0 and 2, respectively. Experimentally it is known that $|A_0|$ is about twenty times $|A_2|$. This is one manifestation of the $\Delta I = 1/2$ rule for strangeness-changing nonleptonic decays. This rule states that transitions which result in an isospin change of 1/2 (such as from a $K$ meson to two pions with total isospin 0) have amplitudes that are much bigger than transitions which result in an isospin change of 3/2 (such as from a $K$ meson to two pions with total isospin 2).

The usual convention is to choose phases so that the larger amplitude, $A_0$ is real; then the parameter $\epsilon$ that involves CP violation in the mass matrix is given by

$$\epsilon \approx \frac{1}{2\sqrt{2}}\, e^{i\pi/4}\, \frac{ImM_{12}}{ReM_{12}}, \tag{4}$$

where $ReM_{12}$, the real part of the off-diagonal element of the $K^0 - \overline{K}^0$ mass matrix is equal to one-half the mass difference between $K_S$ and $K_L$, $\Delta M_K$.

The other parameter for the neutral $K$ system, $\epsilon'$, which is connected to CP violation in the decay amplitude, is in this convention,

$$\epsilon' = \frac{i}{\sqrt{2}}\, e^{i(\delta_2 - \delta_0)}\, \frac{ImA_2}{A_0}, \tag{5}$$

with $\delta_0$ ($\delta_2$) the $\pi\pi$ phase shift in the isospin 0 (2) channel.

We calculate $ImM_{12}$ from the box diagram. This should be a reliable short distance calculation since it is momentum scales between $m_c$ and $m_t$ which contribute an imaginary part to $M_{12}$ in the standard model. However, if we do the calculation in the usual quark basis, we must remember that $A_0$ generally has a phase, $e^{i\xi}$, due to virtual heavy quarks in loops contributing to the decay amplitude. To get back to the usual convention where $A_0$ is real, we redefine the $K^0$ and $\overline{K}^0$ states so as to eliminate this phase. This induces phase changes elsewhere, so that to lowest order in the small quantity $\xi$:

$$ImM_{12}^{sd} \rightarrow ImM_{12}^{sd} + 2\xi ReM_{12}^{sd}, \tag{6}$$

and therefore in the basis where $A_0$ is real,

$$\epsilon = \frac{1}{\sqrt{2}}\, e^{i\pi/4} \left( \frac{ImM_{12}^{sd}}{\Delta M_K} + 2\xi \frac{ReM_{12}^{sd}}{\Delta M_K} \right). \tag{7}$$

Here the superscript $sd$ indicates the short distance contribution to the particular mass matrix element. Since the amplitude $A_2$ is real in the quark basis, it now picks up a phase

$$A_2 \rightarrow A_2 e^{-i\xi} \tag{8}$$

and correspondingly we have:

$$\epsilon' \approx \frac{1}{\sqrt{2}}\, e^{i\pi/4}\, \frac{A_2}{A_0}\, (-\xi), \tag{9}$$

where we have used the experimental $\pi\pi$ phase shifts to write (approximately) the phase of $\epsilon'$.

In the expression for $\epsilon$, we drop the term involving $\xi$. This is good to 20% or better in light of the recent measurements of $\epsilon'$ (see below). Similarly, because of the small measured limit on $\epsilon'$, long distance contributions to $\epsilon$ (which should be of order $20\epsilon'$) are also very likely neglectable.[16] Taking the short distance contribution corresponding to the box diagram and $m_q^2 \ll M_W^2$, we then arrive at the expression:

$$\epsilon \approx \frac{e^{i\pi/4}}{\sqrt{2}} \frac{BG_F^2 f_K^2 m_K}{6\pi^2 \Delta M_K} \times s_1^2 s_2 s_3 s_\delta [-\eta_1 m_c^2 + \eta_2 s_2 (s_2 + s_3 c_\delta) m_t^2 + \eta_3 m_c^2 \ell n(m_t^2/m_c^2)], \tag{10}$$

where the $s_i$ are the sines of the Kobayashi-Maskawa angles $\theta_i$, $i = 1, 2, 3$, which are known to be small so that the approximation $c_i = \cos\theta_i = 1$ has been used in Eq. (10). A non-zero value of the angle $\delta$ is indicative of CP violation in the Kobayashi-Maskawa parametrization.[8] The factors $\eta_1$, $\eta_2$, and $\eta_3$ are due to strong interaction (QCD) corrections and have the values 0.7, 0.6, and 0.4, respectively, with usual quark and W boson masses.[17] The infamous parameter $B$ is the ratio of the actual value of the matrix element between $K^0$ and $\overline{K}^0$ states of the operator composed of the product of two $V$-$A$ neutral, strangeness changing currents divided by the value of the same matrix element obtained by inserting the vacuum between the two currents.

One can see that there can be a potential problem in getting the right-hand-side of Eq.(10) to reproduce the experimental value of $|\epsilon| = 2.27 \times 10^{-3}$ if the combination of Kobayashi-Maskawa angles, $m_t$, and $B$ are not large enough. To get some idea of where the experimental situation places us, take $B = 1/3$, a $b$ quark lifetime of 1 picosecond (near the present world average[14]), and (as it was a year ago), $b \to u/b \to c < .04$. Then for $m_t \lesssim 60\ GeV$ we would have trouble satisfying Eq.(10). For $B \approx 1$, there is no problem satisfying Eq.(10) for any so far unexcluded value of $m_t$, as long as $b \to u/b \to c \gtrsim 0.02$. Smaller values of $b \to u/b \to c$ get us into trouble by lowering the upper bound on $s_3$.

The story from the past year or so in this regard is one of retreat from a confrontation with the standard model. On the experimental side, the upper limit on $b \to u/b \to c$ has become less stringent as it was realized that the theoretically motivated electron spectra used to fit $B$ meson decay are not a good fit to the much improved measured spectra from CESR.[18] As a result the upper limit has risen to a more conservative 0.08 rather than the previous 0.04. Thus $s_3$ can be larger, relieving some of the pressure on the right hand side of Eq.(10).

On the theoretical side, an argument[19] that $B \approx 1/3$ based on using SU(3) to relate $B$ to the amplitude $A_2$, was found to have potential 100% corrections from the next order in chiral SU(3) breaking and is therefore unreliable.[20] Since that time there has been a flurry of papers on the subject,[21] with values of $B$ ranging from 0.3 to 1.5 or so. In my opinion the subject is not conclusively settled, and as long as B can be near unity, the standard model is far from the danger of being excluded as having a credible origin for CP violation through the phase in the Kobayashi-Maskawa matrix.

The situation for $\epsilon'$ has also undergone a similar relaxation. From Eq.(9) we derive, by inserting experimentally measured quantities, that:[22]

$$\epsilon'/\epsilon = -15.6\xi = 6.0\ s_2 s_3 s_\delta \left(\frac{Im\widetilde{C}_6}{-0.1}\right) \left(\frac{<\pi\pi|Q_6|K^0>}{1.0\ GeV^3}\right) \tag{11}$$

where $Q_6$ is the "penguin" operator in the short distance expansion of the strangeness-changing weak Hamiltonian responsible for $K$ decay,[23] and $Im\widetilde{C}_6$ is the imaginary part of the corresponding Wilson coefficient with the Kobayashi-Maskawa factor taken out.

The value of $-0.1$ for this last quantity is relatively stable from calculation to calculation, as the imaginary part depends on momentum scales from $m_c$ to $m_t$ where the short distance expansion is well justified. The value of the matrix element of $Q_6$ is much less certain. If it is large enough to explain the magnitude of $A_0$ (which is a long standing puzzle) then combined with the value of $s_2 s_3 s_\delta$ needed to fit $\epsilon$ (see above), it yields the prediction $\epsilon'/\epsilon \approx +10^{-2}$. This was basically the original observation in Ref. 23: if the "penguin" operator is to be an explanation of the $\Delta I = 1/2$ rule and the magnitude of $A_0$, then $\epsilon'/\epsilon$ should be at the 1% level.

The most recent experiments on the other hand obtain:[24]

$$\epsilon'/\epsilon = (-0.46 \pm 0.53 \pm 0.24) \times 10^{-2} \tag{12a}$$

and,[25]

$$\epsilon'/\epsilon = (+0.17 \pm 0.82) \times 10^{-2} \tag{12b}$$

At the same time theoretical calculations of the matrix element of $Q_6$ have changed dramatically. Early calculations gave numbers of order 1 $\text{GeV}^3$, giving hope that one could explain the magnitude of $A_0$ on the basis of "penguins" and that $\epsilon'/\epsilon$ is of order 1%. In the last couple of years, however, calculations incorporating current algebra constaints in a correct manner give much smaller numbers.[26] These predict values for $\epsilon'/\epsilon$ of order $2 \times 10^{-3}$. These calculations, while not unassailable, have as impeccable theoretical credentials as any others; but they do not allow one to understand the magnitude of the overall $K \to \pi\pi$

amplitude. It is certainly possible to take the attitude that the newer calculations of the matrix element are the correct ones, and the origin of the $\Delta I = 1/2$ rule lies elsewhere. It is also possible to still claim that "penguins" are the primary explanation of the $\Delta I = 1/2$ rule, and either that this is not decisively excluded by the present round of experiments or that the origin of CP violation lies elsewhere than the standard model.[27]

The future in the calculational realm probably belongs to the lattice gauge calculations of matrix elements.[28] While still in their infancy, there are some hopeful signs, but the calculations are done on small lattices and involve approximations or extrapolations which do not allow definitive conclusions. Still, they hold the promise of eventually providing a reliable calculation of these quantities. Experimentally, the next round should push the error bars down to the $10^{-3}$ level. That should allow a real conclusion to be drawn as to whether the "penguin" operator plays a significant role in $K$ decay if it is assumed that CP violation has its origin in the standard model.

# 4. TAU DECAY

All the properties of the tau are consistent with its being a third generation lepton, i.e., just another copy of the electron and muon, albeit much heavier. The front-back asymmetry measurements referred to earlier[5] and the momentum spectrum of the final charged lepton in its purely leptonic decay assure us that its assignment to a left-handed weak doublet is correct. The other member of this doublet, the tau neutrino, must be distinct from the electron and muon neutrinos. The upper bound on the tau neutrino mass has recently been lowered to 70 MeV, below the mass of the charged lepton of the previous generation.[29]

The one problem of consequence comes in comparing the sum of the exclusive decay mode measurements with the inclusive charged-prong multiplicity measurements. In particular, it is difficult to account for the origin of all the decays of the tau that result in one charged prong.[30] This arises as follows.

Let us normalize all the theoretical calculations of decay branching ratios to that for $\tau \rightarrow \nu_\tau e \bar{\nu}_e$, which we take to be 17.9% (the world average value[29] is $17.9 \pm 0.3\%$). Then the present experimental situation is summarized in Table I for decays of the tau involving three charged prongs. First of all we see that the sum of the exclusive modes (almost entirely $\tau^- \rightarrow \nu_\tau \pi^- \pi^- \pi^+$ and $\tau^- \rightarrow \nu_\tau \pi^- \pi^- \pi^+ \pi^0$) is in agreement with the three charged-prong inclusive branching ratio. Second, where there is a theoretical prediction, it is in good agreement with experiment.[29,30]

The agreement between theory (where there is a prediction of some accuracy) and experiment (where there is a definite measurement) is also very good in the case of tau decays involving one charged prong, shown in Table II. Note in particular the decays $\tau \rightarrow \nu_\tau \pi$ and $\tau \rightarrow \nu_\tau K$, whose rates follow from those for

$\pi \to \mu\nu$ and $K \to \mu\nu$, respectively. There is also the major decay $\tau^- \to \nu_\tau \pi^- \pi^0$, whose rate follows from using CVC to relate it to an integral over $e^+e^- \to \pi^+\pi^-$ cross sections, with a result that is in excellent agreement with experiment. While there is not a precise experimental number with which to compare it, the theoretical prediction for $\tau^- \to \nu_\tau \pi^- 3\pi^0$ follows in a similar way from an integral over the well measured cross section for $e^+e^- \to 2\pi^+ 2\pi^-$ , and one has every reason to have confidence in the prediction.[30]

The theoretical upper bounds on the modes in Table II follow from using isotopic spin invariance to bound one charge combination of the final hadrons in terms of another combination which has been measured, e.g. $\tau^- \to \nu_\tau \pi^- 4\pi^0$ is bounded from the measured decay $\tau^- \to \nu_\tau 3\pi^- 2\pi^+$. In particular, the major decay mode $\tau^- \to \nu_\tau \pi^- 2\pi^0$ has a branching ratio which must be less than that of the well measured mode $\tau^- \to \nu_\tau 2\pi^- \pi^+$. From the properties of this latter mode,[29] which is dominated by the $A_1$ resonance decaying into $\rho\pi$, the inequality should be an equality. The 2% limit on the last three modes is a fairly generous upper limit.

The sum of the branching ratios from theory is $\leq 81.7\%$ in Table II, while the measured one charged prong inclusive branching ratio is $86.6 \pm 0.3\%$.[29] We have accounted for all the purely leptonic modes and all the modes of the form $\tau^- \to \nu_\tau (n\pi)^-$ of any consequence as well as Cabibbo suppressed modes. Where are the remaining 5% of one prong decays?

One possibility is that the branching ratio for $\tau \to \nu_\tau e \bar{\nu}_e$, which we took to be 17.9% (and to which we normalized all our theoretical predictions), should be $\approx 19\%$. This would scale up all the predicted branching ratios by $\approx 6\%$, making the agreement of experiment and theory worse in Table II. It would also put the $\tau$

lifetime measurements[14] about one standard deviation away from their predicted standard model value. While the data over the past year has tended to make the discrepancy more significant, it is still possible that most of the problem could go away by the branching ratios moving upward.[31]

A second possibility is that there are other conventional decay modes which we have neglected as very likely small, but which in fact have branching ratios amounting to several percent. A candidate for such a role is $\tau \rightarrow \nu_\tau \eta \pi \pi$. There is no particular argument that would demand a sizable branching ratio for this channel, but no direct evidence that limits it to the few tenths of a percent level at which one might have naively estimated it.

Finally, there is the possibility of new physics. There can not be a new elementary charged particle into which the $\tau$ decays, for it would have been pair produced in electron-positron annihilation and made the value of $R$ disagree with experiment. Thus we need a new neutral particle produced along with conventional ones, and/or a distortion of the expected branching ratios due to the effects of a heavy virtual particle. It is hard to find a scenario for this situation which is not very contrived (especially if it is not to be in conflict with other existing experiments). Perhaps the whole problem will just creep away a few percent at a time.

# 5. HEAVY NEUTRAL LEPTONS

The most straightforward extension of the presently known neutral leptons, i.e. the electron, muon, and tau neutrinos, is to add a fourth, so-called "sequential" neutrino as a part of a fourth generation consisting of a charge $2e/3$ quark, a charge $-e/3$ quark, a charge $-e$ lepton, and a neutrino. The left-handed quarks and leptons reside in weak isospin doublets, while all right-handed fermions are singlets, as for the first three generations. Moreover, in the canonical version of the standard model there is no right-handed neutrino field at all and simply no way to obtain a massive neutrino: the neutrino is left-handed and massless. More generally, going slightly beyond the standard model, it is entirely possible to supply a right-handed singlet field and make such a neutrino a massive fermion. Then, just as for the quark sector, the weak and mass eigenstates will not coincide. This is conveniently expressed in terms of a unitary matrix $U$ which "rotates" the neutrino mass eigenstates (the column index, labelled by the generation number) to the weak eigenstates (the row index, labelled by the corresponding charged lepton). For example, the neutrino which is in a left-handed doublet with the electron, $\nu_e$, is the superposition of mass eigenstates, $\nu_j$ , given by

$$\nu_e = \sum_{j=1}^{4} U_{e\,j} \nu_j \tag{13}$$

in the case of four generations. Here, since all the neutral leptons have the same value of weak isospin, there are no lepton flavor-changing neutral currents.[32] This generalizes to adding an arbitrary number of generations. The unitary property of the matrix $U$ implies that the $Z$ couples to each pair of neutrino mass eigenstates with the same universal strength, but there are no couplings of the $Z$ (at tree level) to different mass eigenstates.

Neutral leptons are also predicted which behave as singlets under the weak isospin of the standard model. Such is the case, for example, in some grand unified models such as $O(10)$, where the 16 dimensional representation includes all the usual quarks and leptons of one generation plus a right-handed singlet, neutral lepton.[33] In left-right symmetric theories as well, there are often heavy neutral leptons.[34] In particular, they arise as the partners of the usual charged leptons under the additional $SU(2)$ gauge group with right-handed couplings, but are singlets under the usual $SU(2)$ gauge group with left-handed couplings of the standard model.

A particularly attractive reason for having further singlet neutral leptons is the so-called "see-saw" mechanism[35] for generating neutrino masses. With each left-handed neutrino one associates a right-handed partner, $N$, as in the sixteen dimensional representation of $O(10)$, so that the mass matrix involving $\nu$ and $N$ looks like

$$M = \begin{pmatrix} 0 & m_D \\ m_D & M \end{pmatrix} . \tag{14}$$

Here $m_D$ is a Dirac mass, presumably comparable to a quark or charged lepton mass of that generation, and $M$ is a large Majorana mass. The eigenstates of this matrix will have masses of $m_D^2/M$ and $M$ approximately. Thus we have light neutrinos as well as heavy ones. The mixing matrix elements, $U_{\ell N}$, between the light neutrinos (which are members of the usual left-handed isodoublets with corresponding charged leptons, $\ell$) and the heavy neutral singlets are of order $m_D/M$. We can also arrange matters so that the light neutrino remains massless, but still mixes with a singlet Dirac heavy neutral lepton.[36,37]

Another possibility is the existence of mirror neutrinos, i.e. neutral leptons in right-handed doublets under the usual $SU(2)$, together with corresponding mirror charged leptons.[38] Such neutral leptons are left-handed singlets; hence the name mirror fermions as they are just the mirror image of the usual fermions in the standard model. A recent case in point is provided by the $O(18)$ grand unified model[39] which predicts a fourth generation with the usual left-handed weak interaction couplings as well as four(!) generations of quarks and leptons with right-handed couplings. Each of these generations contains a neutral lepton with a mass below about 40 GeV.

At the same time that theoretical interest has been increasing,[40–42] the scope and sensitivity of experimental techniques has also been undergoing a dramatic change.[43] High flux hadron beams at both low and high energy now permit very low limits being placed on the emission of neutral leptons in leptonic or semileptonic hadron decays. At electron-positron colliders, center-of-mass energies and integrated luminosities have reached the level where weak interaction cross sections are capable of being responsible for production of detectable numbers of new neutral leptons. This, along with the concurrent development of precision vertex detectors at such machines,[44] has made it possible to extend the mass range under investigation by over an order of magnitude. While previously one was limited by the mass of the decaying hadron, in many cases the kinematic limit now is simply the center-of-mass energy of the electron-positron collision. The near future should see this line of analysis take another large jump in sensitivity through the study of $Z$ decays into neutral lepton pairs. This will permit a high sensitivity sweep of the mass range from zero to half the $Z$ mass.

A variety of techniques have been used over the years in high energy physics experiments to search for neutral heavy leptons. Before the searches at electron-positron colliders, most of the previous limits on mixing of a heavy neutrino with the electron or muon neutrino came from hadron leptonic or semileptonic decays. A direct and powerful technique[45] is to study pion or kaon decays at rest, searching for additional monochromatic peaks in the electron or muon momentum spectra in the leptonic decays $\pi \to e\nu$, $\pi \to \mu\nu$, $K \to e\nu$, or $K \to \mu\nu$, as appropriate. Upper limits on the mixing with a fourth generation neutrino extracted using this technique are shown[46–49] as the limiting curves for $|U_{e4}|^2$ labelled (1), (2), and (3) in Figure 2 and the curves for $|U_{\mu 4}|^2$ labelled (3) and (12) in Figure 3. At the present time, limits on the square of these mixing matrix elements are in the region of $10^{-6}$ from such experimental searches.

A different, but related technique in that it relies on the decay in flight of the same hadrons, namely $\pi$ and $K$ mesons, arose when it was realized[50] that the heavy neutral leptons so produced would themselves decay downstream. It was therefore possible to use existing neutrino detectors or modifications thereof to conduct searches for the subsequent decay of the heavy neutral leptons produced in high intensity $\pi$ and $K$ beams. In the potentially observed experimental rate, the square of a mixing matrix element enters twice: first in the production where the neutral lepton is born in $\pi$ or $K$ decay through mixing with the electron or muon neutrino, and then again when the heavy lepton decays weakly into ordinary leptons or leptons and quarks. Depending on the beam and the particular sensitivities of the detector these two mixing matrix elements could be the same or different. Limits obtained from $\pi$ and $K$ beams using this technique[51,52] are given by curves (4) and (10) in Figure 2, where it is seen that for neutral lepton masses below the $K$ mass they now provide the best upper limits on $|U_{e4}|^2$, going down to below $10^{-8}$.

Once we are utilizing hadron decays in flight, we need not restrict our attention to $\pi$ and $K$ decays, as we are limited by their masses. Charmed meson decays extend the range of investigation considerably. To enhance their proportion of the neutrino flux one employs a beam dump to absorb longer lived hadrons before they can decay weakly, leaving short lived mesons (and baryons) which have "prompt" decays, typically a small fraction of an absorption length. With the increased mass range opened up by the charmed meson masses, both purely leptonic and semileptonic decays are potential sources of new neutral leptons in an interesting mass range. Some of the recent limits[53–55] obtained in this manner are given by curves (5), (9), and (11) in Figures 2 and 3 for $|U_{e4}|^2$ and $|U_{\mu 4}|^2$, respectively. It is also possible to get limits on $|U_{e4}U_{\mu 4}|$ using hadron decays in flight (either with beams or in beam dumps) by, for example, producing the heavy neutrino in association with a muon and detecting its charged-current decay involving an electron.[51,52]

There is no reason to stop at charm with this technique. Hadrons containing bottom and even top quarks can be employed; one is only limited by their production cross sections in hadron-hadron collisions. Indeed it has been pointed out[56] that beam dump experiments utilizing $B$ meson decays could extend the limits to masses of roughly 2.5 GeV$/c^2$ at the $10^{-6}$ level. One should also note that all the limits we have discussed up to now involve charged-current weak interactions to produce the heavy neutral lepton in association with ordinary charged leptons. Hence the relevant mixing matrix elements and amplitudes at the production vertex are independent of whether a doublet, a singlet, or a mirror neutral lepton is under scrutiny.

To extend the search to the domain of yet higher masses we turn to electron-positron colliders where we are limited kinematically only by the center-of-mass energy for $e^+e^- \to \overline{\nu}_4\nu_e$ (and half of it for pair production of heavy neutrinos). If the neutrino is in a left-handed doublet, as for a fourth generation neutrino, then its production proceeds in lowest order through W exchange with a cross section[57] for $s \ll M_W^2$

$$\sigma(e^+e^- \to \overline{\nu}_4\nu_e) = |U_{e4}|^2 \, \frac{G_F^2 s}{6\pi} \, (1 - m^2/s)^2(1 + m^2/2s) \tag{15}$$

where $m$ is the neutrino mass and $s$ the square of the center-of-mass energy.

Such a process will be observed as an electron-positron collision which results in an event with missing energy and momentum (due to the $\nu_e$) on one side and, if $m$ is not too large, a jet of decay products of $\nu_4$ on the other side, i.e. a monojet event. Such events have been searched for recently at PEP and PETRA in another context.[58–61] By combining the PEP data one obtains[62] the upper limit shown as curve (7) in Fig. 2. This is already a considerable improvement over the limit following from universality[63] shown as curve (6), and can be improved still further, as we discuss below.

In contrast with Eq. (15), the cross section for production of a pair of "sequential" Dirac neutrinos through a "virtual" $Z$:

$$\sigma(e^+e^- \to \overline{\nu}_4\nu_4) =$$

$$\frac{G_F^2 s}{24\pi} \, \frac{(1 - 4m^2/s)^{1/2}(1 - m^2/s)}{1 - s/M_Z^2} \, (1 - 4\sin^2\theta_W + 8\sin^4\theta_W) \tag{16}$$

does not contain a factor of $|U_{e4}|^2$. With the presently accumulated luminosity at PEP, dozens of such heavy neutrino pairs should have be produced in each experimental region as long as their mass is at least slightly below the beam energy. Since the decay width is still proportional to the square of a mixing matrix element, a search for secondary decay vertices sweeps out a region in the $|U|^2$ versus $m$ plane. The results of such a search between 0.2 and 10 cm from the interaction point are shown[64] as the diagonal region going up to masses of almost 14 GeV/$c^2$, bounded by curve (8) in Figures 2 and 3. The analogous result for a "sequential" neutrino mixed with the tau neutrino is shown in Fig. 4.

What happens to all these bounds if we had been considering a singlet neutrino rather than a "sequential" neutrino in a doublet? All the bounds that come from hadron leptonic or semileptonic decay remain the same (aside from slight shifts in limits for those experiments which rely on detecting particular decays whose predicted branching ratios change somewhat due to additional neutral-current interactions).[65] The cross section for $e^+e^- \to Z_{virtual} \to \overline{N}N$ picks up a factor of $|U|^4$ over Eq. (16), making it negligibly small given present experimental sensitivities and interesting values of $|U|^2$, and we unfortunately lose the restrictions obtained from the search for secondary vertices in electron-positron annihilation. In the case of $e^+e^- \to N\overline{\nu}_e$, we now have both $W$ exchange and direct-channel $Z$ contributions and the expression for the cross section becomes

$$\sigma(e^+e^- \to N\overline{\nu}_e) = $$
$$|U_{eN}|^2 \, \frac{G_F^2 s}{24\pi} \left(1 - \frac{m^2}{s}\right)^2 \left(1 + \frac{m^2}{2s}\right) \, (1 + 4\sin^2\theta_W + 8\sin^4\theta_W) \, . \qquad (17)$$

If $\sin^2\theta_W = 0.22$, then the cross section in Eq. (17) is 0.57 times that in Eq. (16) and the monojet searches have somewhat reduced sensitivity.

In contrast to the reduced sensitivity to mixing of $N$ with the electron neutrino in monojet searches, we gain from the neutrino flavor changing coupling of the $Z$ a direct-channel $Z$ contribution to $e^+e^- \to N\overline{\nu}_\mu$ which is proportional to $|U_{\mu N}|^2$:

$$\sigma(e^+e^- \to N\overline{\nu}_\mu) = \\ |U_{\mu N}|^2 \frac{G_F^2 s}{24\pi} \left(1 - \frac{m^2}{s}\right)^2 \left(1 + \frac{m^2}{2s}\right) \left(1 - 4\sin^2\theta_W + 8\sin^2\theta_W\right) . \tag{18}$$

An identical expression with $\tau$ in place of $\mu$ follows from potential mixing with the tau neutrino. The limit that is obtained for $|U_{\mu N}|^2$ from Eq. (18) plus existing experiments is not better than that from universality.

However, the presence of neutrino flavor changing couplings of the $Z$ does lead to a new limit on $|U_{\mu N}|^2$ from the lack of observation of the process $\nu_\mu$+nucleus $\to$ $N + \ldots$, followed by decay of the $N$. An additional region has been excluded by the CHARM collaboration[53] in this way.

The experimental situation for mirror neutrinos is somewhat of a cross between those for sequential and right-handed singlet heavy neutral leptons. Since mirror neutrinos lie in right-handed doublets, they have full strength couplings to the $Z$ (like sequential neutrinos) and the limits from electron-positron annihilation experiments pertain to them (curve (8) in Figures 2, 3, and 4). On the other hand, being left-handed singlets, they also have flavor changing couplings to the $Z$ and limits following from production through such couplings also are applicable.

In the near future it should be possible to considerably extend many of the limits by further analysis of existing data. Combining all the experiments it should be possible, for example, to exclude values of $|U_{e4}|^2$ and $|U_{\mu4}|^2$ above $10^{-6}$ up to masses of about 20 GeV/$c^2$. Upper limits on $|U_{e4}|^2$ of $10^{-2}$ to $10^{-3}$ should be obtainable from PETRA data for masses up to 30 GeV/$c^2$ from monojet searches.

A dramatic increase in sensitivity will occur when the SLC and LEP come into full operation. With six percent of $Z$ decays going into neutrino-antineutrino for each light neutrino in a weak doublet, additional sequential or mirror neutrinos will be very amenable to detection. The presence of such particles with low masses can be ascertained by "neutrino counting", i.e. tagging $Z$ decays into unseen neutrals and seeing if the resulting decay width is accounted for by the known neutrinos. Those with higher masses will generally decay through mixing, but here the technique of using vertex detectors will enable us to exclude masses up to about 40 GeV/$c^2$ and mixing matrix elements squared down to around $10^{-10}$. Monojet searches looking for $Z \rightarrow N\overline{\nu}$ could be sensitive to weak singlet heavy neutral leptons with masses up to a large fraction of the $Z$ mass and squares of mixing matrix elements to the known neutrinos down to about $10^{-5}$.

## 6. EXTRA $Z$'s

Over the years there have been numerous suggestions in the direction of extending the electroweak gauge group beyond $SU(2) \times U(1)$. In particular, many grand unified theories have extra $Z$'s as do left-right symmetrical theories such as $SU(2)_L \times SU(2)_R \times U(1)$.

The investigation of this area has received a big boost from the advent of superstrings, where the combined low energy gauge group is generally larger than $SU(3)_c \times SU(2)_L \times U(1)_Y$.[66–69] An early favorite in this regard is to have $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)$ at "low" energies, but this is by no means the only possibility.[67,68,69] The patterns of bosons and fermions following under various assumptions has been analyzed by many people.[70] A number of these possibilities can already be ruled out, depending in part on what other theoretical assumptions are made (such as the number of scales of symmetry breaking). Still, a range of possibilities with at least an extra U(1) and associated neutral gauge boson still remains, and it is interesting to examine how they are constrained by existing experiments.

The constraints on an extra $Z$ which we have at our disposal come from neutral current experiments, the measured $W$ and $Z$ masses, and the lack of observation of a second $Z$ in collider experiments up to now. One needs to apply these constraints to each proposed low energy model in turn to see what limits can be placed on the mass and mixing of the extra $Z'$ with the Z of the standard model. The results[71] are somewhat surprising in that a $Z'$ with a mass as low as about 130 GeV is allowed, as long as the mixing with the standard model $Z$ is fairly small (so as not to displace it too far from its nominal mass). This comes about in part because the couplings of the $Z'$ given by such models tend to be small to ordinary fermions. This also makes them harder to produce in ordinary lepton or hadron collisions than the usual $Z$, and it makes their decay branching ratios to ordinary leptons smaller if the "exotic" channels for decay are open. In spite of this, up to masses of several hundred GeV they have production cross sections in $\bar{p}p$ collisions which are adequate enough that we will be watching the Tevatron collider to see if it turns up direct evidence of their existence in the next few years.

# 7. CONCLUSION

As we noted at the beginning, the agreement of the standard model and experiment is extremely impressive. The pickings are pretty slim if one is looking for areas of potential discrepancy.

Nevertheless, it is worth keeping an eye on the situation with respect to CP violation and, to a lesser extent, the decays of the tau. While there is not yet any smoking gun, another round of experiments is in progress. We are guaranteed continued improvement in the accuracy of our knowledge of various quantities and the possibility for surprises along the way certainly exists.

We have also seen that there are very sensitive limits on certain possibilities that go beyond the standard model. In particular, the limits on heavy neutral leptons are very impressive and promise to improve by several orders of magnitude in sensitivity to small mixing angles, as well as to extend considerably further in mass range during the next few years.

However, the limits on other possibilities are rather poor. In particular, I have in mind the constraints which present experiments put on extra $Z$'s. New physics may be just around the corner.

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## TABLE I

THREE CHARGED PRONG DECAYS OF THE $\tau$

| Decay Mode | Branching Ratio (%) Theory[a] | Experiment[b] |
|---|---|---|
| $\tau^- \rightarrow \nu_\tau 2\pi^-\pi^+$ | | $8.1 \pm 0.6$ |
| $\tau^- \rightarrow \nu_\tau 2\pi^-\pi^+\pi^0$ | 4.9 | $5.0 \pm 0.6$ |
| $\tau^- \rightarrow \nu_\tau (K_s\pi)^-$ | 0.3 | $0.3 \pm 0.1 \pm 0.1$ |
| $\tau^- \rightarrow \nu_\tau K^-\pi^-\pi^+$ | | $0.22 \pm 0.14$ |
| $\tau^- \rightarrow \nu_\tau K^-K^+\pi^-$ | | $0.22 \pm 0.14$ |
| $\tau^- \rightarrow \nu_\tau 2\pi^-\pi^+2\pi^0$ | | |
| $\tau^- \rightarrow \nu_\tau 2\pi^-\pi^+2\pi^0$ | $< 0.4$ | |
| *TOTAL* | | $13.2 \pm 0.3$ |

($a$) All theoretical branching rates are normalized to that for $\tau \rightarrow \nu_\nu e\overline{\nu}_e$ taken as 17.9%. Calculations from Ref. 30.

($b$) Experimental values are taken from Ref. 29.

## TABLE II

ONE CHARGED PRONG DECAYS OF THE $\tau$

| Decay Mode | Branching Ratio (%) Theory[a] | Experiment[b] |
|---|---|---|
| $\tau^- \rightarrow \nu_\tau e^- \overline{\nu}_e$ | 17.9 (Input) | $17.5 \pm 0.7$ |
| $\tau^- \rightarrow \nu_\tau \mu^- \overline{\nu}_\mu$ | 17.4 | $18.1 \pm 0.6$ |
| $\tau^- \rightarrow \nu_\tau \pi^-$ | 10.9 | $10.3 \pm 1.2$ |
| $\tau^- \rightarrow \nu_\tau \pi^- \pi^0$ | 22.0 | $22.0 \pm 2.1$ |
| $\tau^- \rightarrow \nu_\tau \pi^- 2\pi^0$ | $\leq 8.1 \pm 0.6$ | |
| $\tau^- \rightarrow \nu_\tau \pi^- 3\pi^0$ | 1.0 | |
| $\tau^- \rightarrow \nu_\tau \pi^- 4\pi^0$ | $< 0.1$ | |
| $\tau^- \rightarrow \nu_\tau \pi^- 5\pi^0$ | $< 0.1$ | |
| $\tau^- \rightarrow \nu_\tau K^-$ | 0.7 | $0.6 \pm 0.2$ |
| $\tau^- \rightarrow \nu_\tau (K\pi)^-$ | 0.9 | $0.9 \pm 0.3 \pm 0.3$ |
| $\tau^- \rightarrow \nu_\tau (K\pi\pi)^-$, $\tau^- \rightarrow \nu_\tau (K\overline{K}\pi)^-$, $\tau^- \rightarrow \nu_\tau (K\overline{K})^-$ | $< 2$ | |
| *TOTAL* | $< 81.7$ | $86.6 \pm 0.3$ |

(*a*) All theoretical branching rates are normalized to that for $\tau \rightarrow \nu_\nu e \overline{\nu}_e$ taken as 17.9%. Calculations from Ref. 30.

(*b*) Experimental values are taken from Ref. 29.

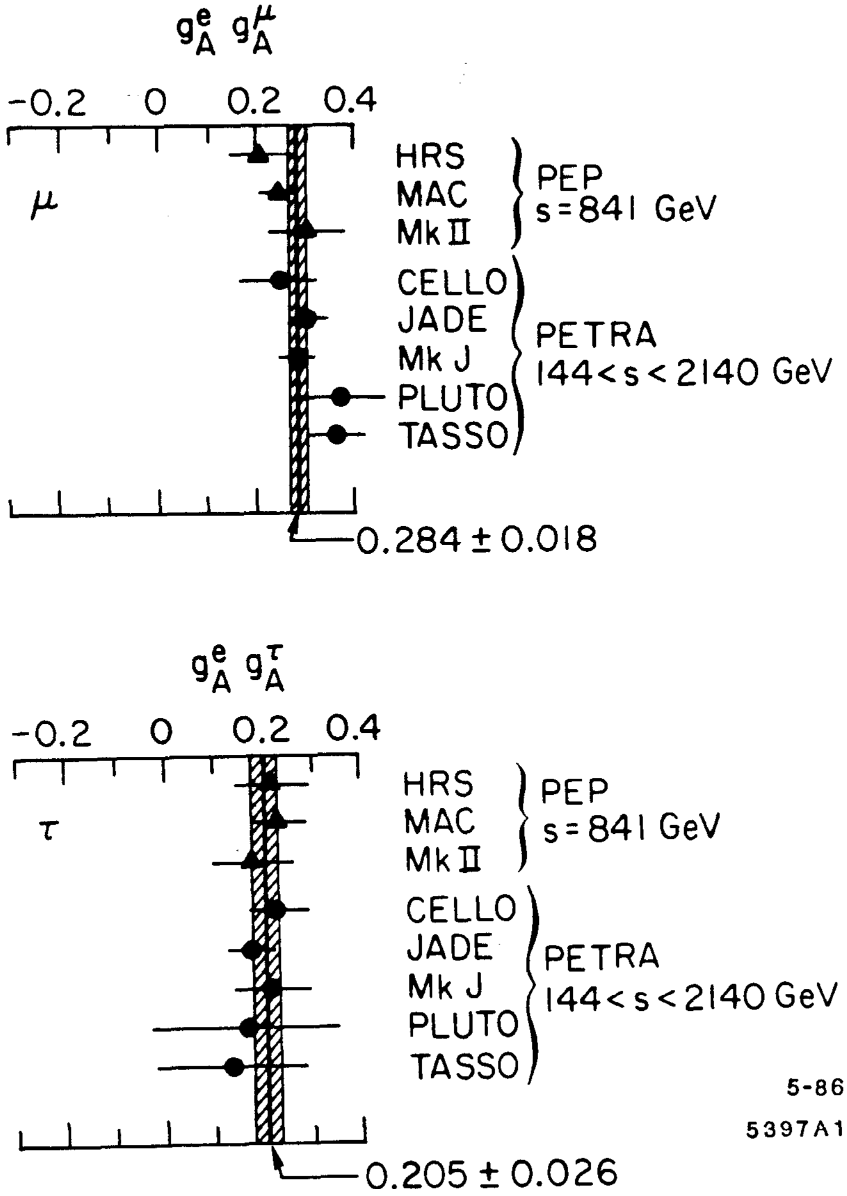
## FIGURE CAPTIONS

1. The $g_A^e g_A^\mu$ and $g_A^e g_A^\tau$ values derived from experiments measuring the respective $\mu + \mu^-$ and $\tau^+\tau^-$ front-back asymmetries at PEP and PETRA. The world average values are indicated, to be compared to the standard model expectation for both quantities of +0.25. Figure from Ref. 5.

2. Limits on $|U_{e4}|^2$ as a function of the mass $M_4$ of a sequential, fourth generation neutrino as obtained from (1) TRIUMF $\pi \to e\nu$, Ref. 46; (2) SIN $\pi \to e\nu$, Ref. 47; (3) KEK $K \to e\nu$, Ref. 48; (4) CHARM experiment at CERN with a wide band beam, Ref. 51; (5) CHARM experiment at CERN using a beam dump, Ref. 53; (6) Universality, Ref. 63; (7) Monojet searches at PEP, Ref. 62; (8) Mark II secondary vertex search at PEP, Ref. 64; (9) Beam dump experiment at Fermilab, Ref. 54; (10) Wide band beam experiment at the PS at CERN, Ref. 52; (11) BEBC experiment at CERN using a beam dump, Ref. 55.

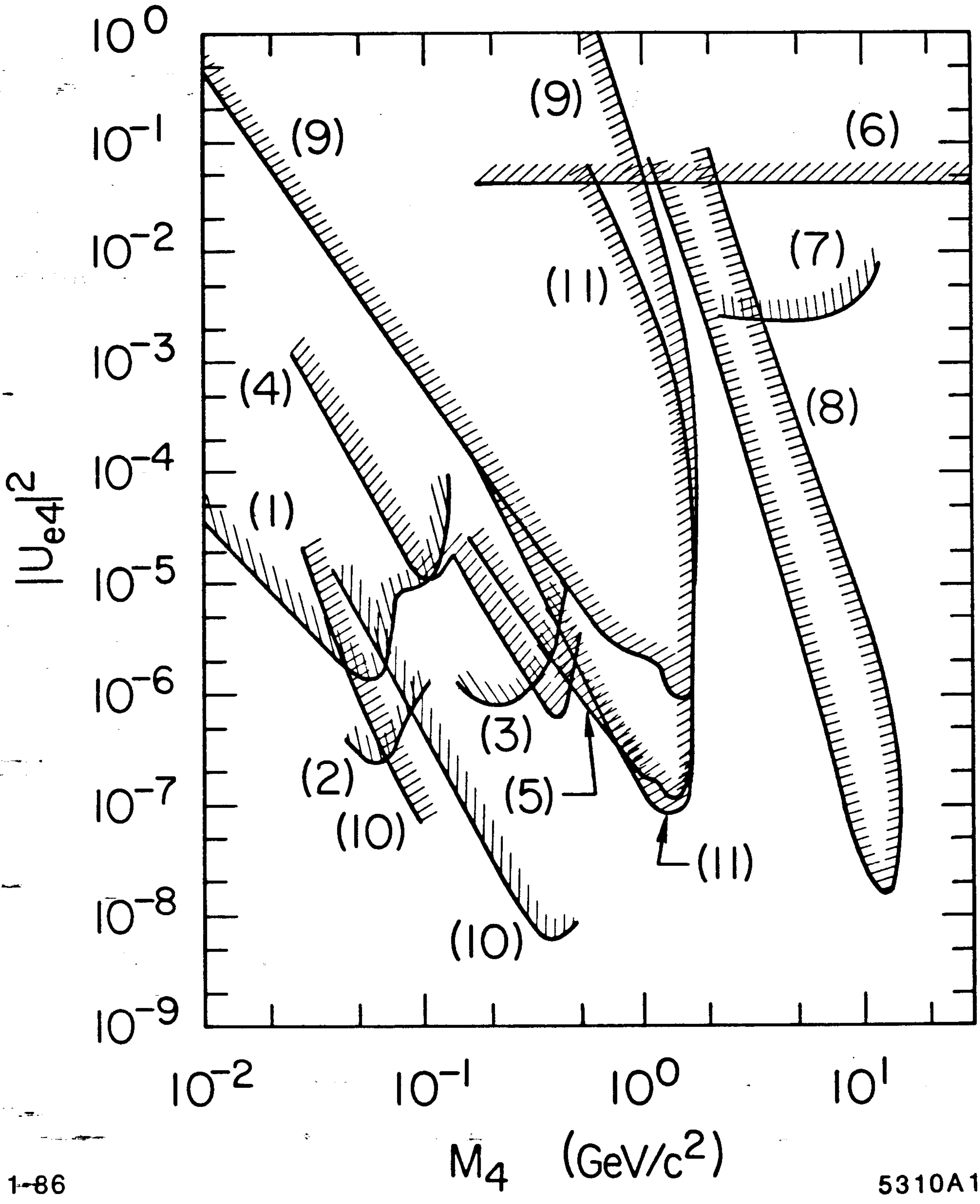
3. Limits on $|U_{\mu 4}|^2$ as a function of the mass $M_4$ of a sequential, fourth generation neutrino as obtained from (3) KEK $K \to \mu\nu$, Ref. 48; (5) CHARM experiment at CERN using a beam dump, Ref. 53; (6) Universality, Ref. 63; (8) Mark II secondary vertex search at PEP, Ref. 64; (9) Beam dump experiment at Fermilab, Ref. 54; (11) BEBC experiment a CERN using a beam dump, Ref. 55; (12) SIN $\pi \to \mu\nu$, Ref. 49.

4. Limits on $|U_{\tau 4}|^2$ as a function of the mass $M_4$ of a sequential, fourth generation neutrino as obtained from (8) Mark II secondary vertex search at PEP, Ref. 64.



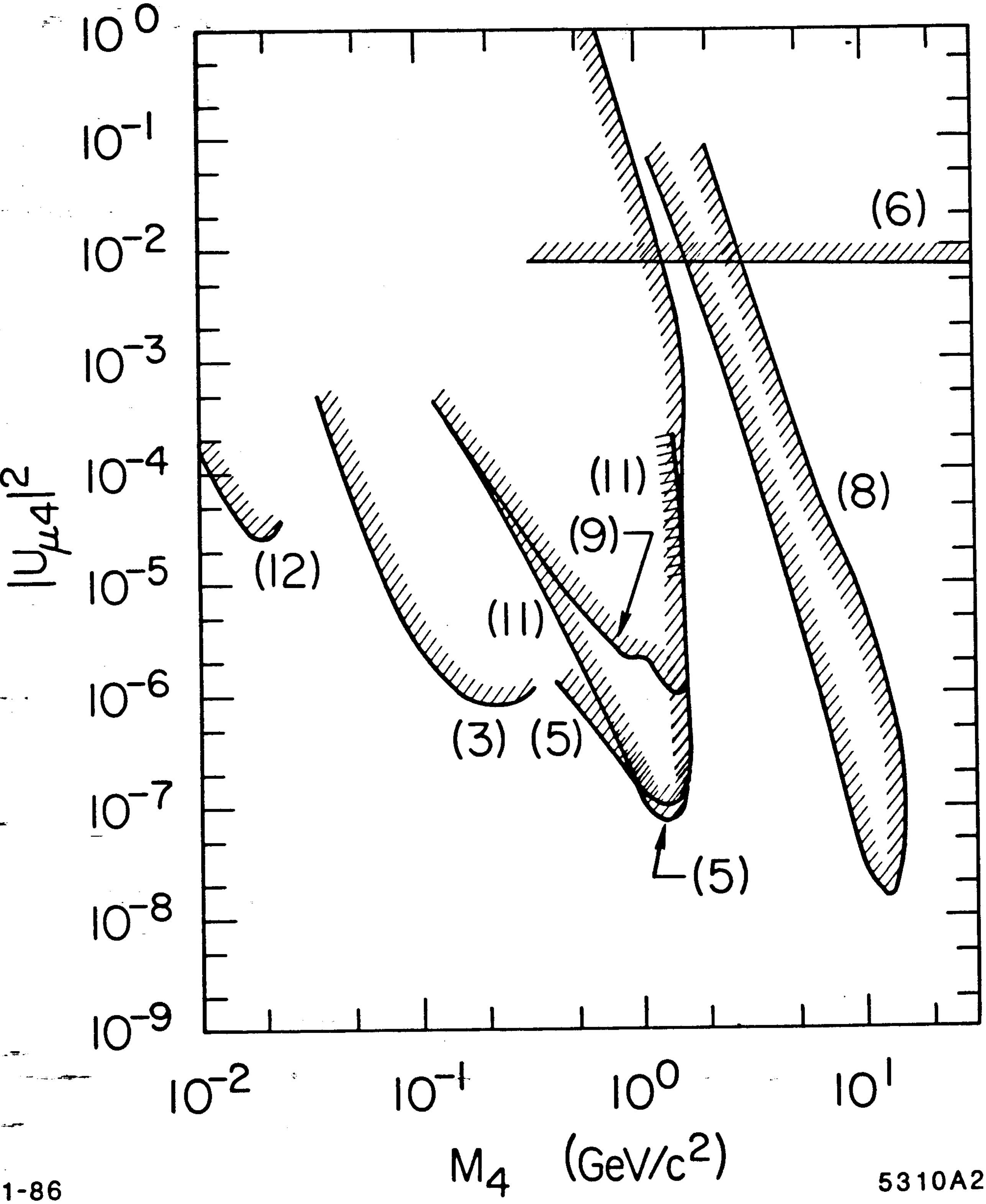
$g_A^e\, g_A^\mu$
-0.2 0 0.2 0.4
μ
HRS
MAC
Mk II
PEP
s=841 GeV
CELLO
JADE
Mk J
PLUTO
TASSO
PETRA
144<s<2140 GeV
0.284 ± 0.018
$g_A^e\, g_A^\tau$
-0.2 0 0.2 0.4
τ
HRS
MAC
Mk II
PEP
s=841 GeV
CELLO
JADE
Mk J
PLUTO
TASSO
PETRA
144<s<2140 GeV
5-86
5397A1
0.205 ± 0.026


Fig. 1



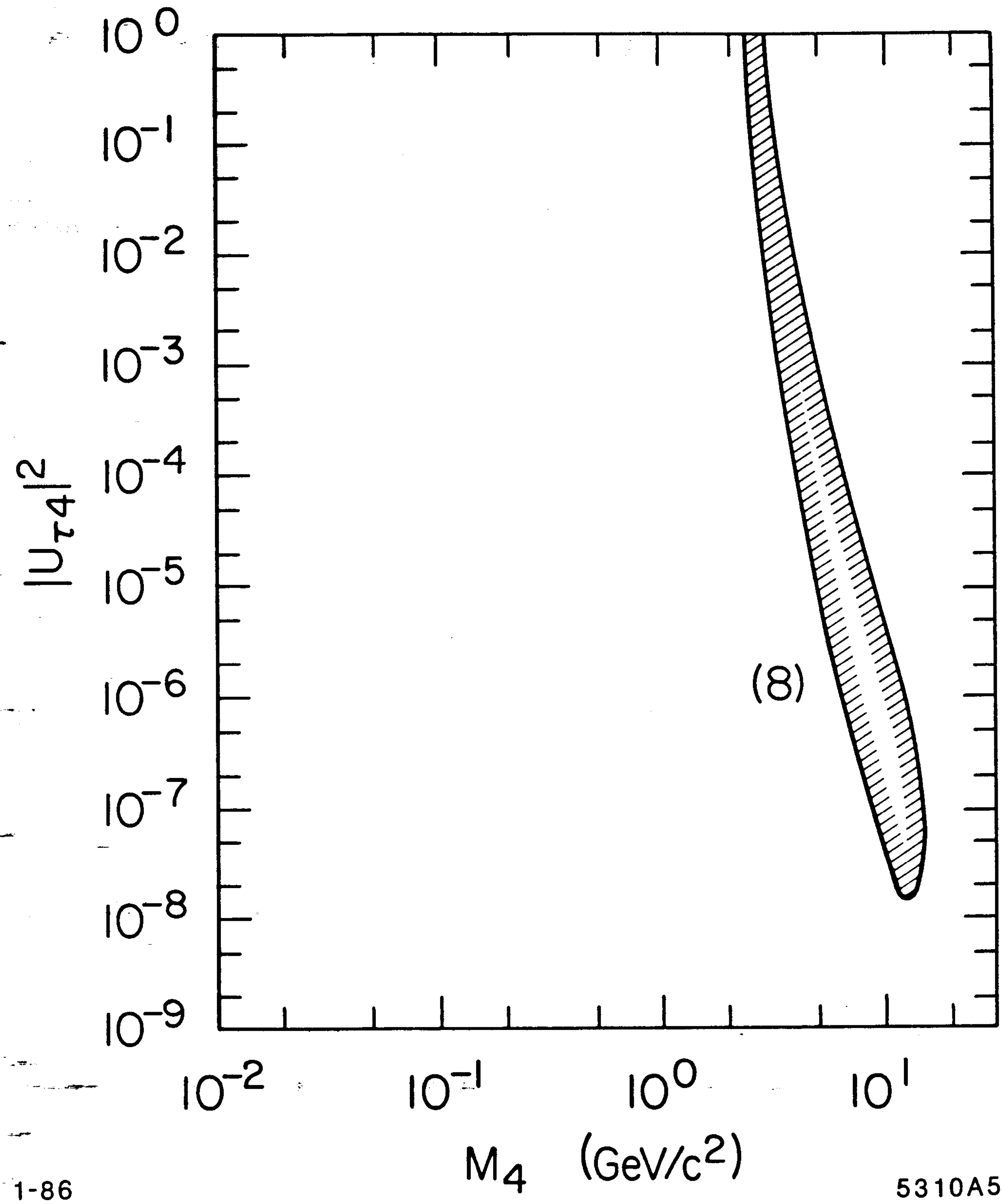
$10^0$
$10^{-1}$
$10^{-2}$
$10^{-3}$
$10^{-4}$
$10^{-5}$
$10^{-6}$
$10^{-7}$
$10^{-8}$
$10^{-9}$
$|U_{e4}|^2$
$10^{-2}$
$10^{-1}$
$10^0$
$10^1$
$M_4$ (GeV/c$^2$)
(1)
(2)
(3)
(4)
(5)
(6)
(7)
(8)
(9)
(9)
(10)
(10)
(11)
(11)
1-86
5310A1


Fig. 2



$10^0$
$10^{-1}$
$10^{-2}$
$10^{-3}$
$10^{-4}$
$10^{-5}$
$10^{-6}$
$10^{-7}$
$10^{-8}$
$10^{-9}$
$|U_{\mu 4}|^2$
(6)
(11)
(8)
(9)
(12)
(11)
(3)
(5)
(5)
$10^{-2}$
$10^{-1}$
$10^0$
$10^1$
$M_4$ (GeV/$c^2$)
1-86
5310A2


Fig. 3



$10^0$
$10^{-1}$
$10^{-2}$
$10^{-3}$
$10^{-4}$
$10^{-5}$
$10^{-6}$
$10^{-7}$
$10^{-8}$
$10^{-9}$
$|U_{\tau 4}|^2$
(8)
$10^{-2}$
$10^{-1}$
$10^0$
$10^1$
$M_4$ (GeV/c$^2$)
1-86
5310A5


Fig. 4