Why – and Where – to Look for Heavy Majorana Neutrinos

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Much of the motivation for the present Summer Study comes from the wish to push precision studies of the Standard Model to the limits of experimental finesse and interpretative acumen; and, having thus cleared the roster for the first steps beyond its obvious shortcomings, to find traces of *New Phenomena* in all the expected spots. The agenda of this very fruitful study period is replete with ingenious ways to explore the weak-coupling as well as the strong-coupling mechanisms of electroweak symmetry breaking (EWSB). In this context, it is always amazing to some of us with what levity our community accepts the most glaring break in an orderly lepton/quark symmetry on the SU_2 level, declaring the three observed neutrino species to be massless, their right-handed incarnations to be non-existent.

In the process, it is equally astounding that we tend to define a global symmetry, lepton number, on a logical par with the jealously guarded baryon number, which is not even defined in the case of massless neutral leptons: lepton number conservation is therefore a somewhat meaningless concept where neutrinos are involved; even the much more strongly supported concept of lepton *flavor*, often loosely confused with lepton *number*, is easily violated by the concept of Majorana neutrinos.

It is quite obvious that our considerations of where we are likely to find experimental signals for *Beyond-the Standard-Model* phenomena such as massive neutrinos, and for their appropriate group-theoretical assignment, should be seen as one of the most important tasks of the new generation(s) of experimentation that our future facilities will have to face up to.

In the following, I will argue that the Linear Electron Colliders that are being discussed here will have a chance of making decisive contributions to the neutrino mass and lepton number puzzles, if we run them in the electron–electron modes; more specifically, we see this set of experimentation as uniquely promising, much more so than the classical methods of ferreting out neutrino properties from the well-hidden nuclear physics foxholes that may permit the massiveness and Majorana character of light neutrinos to be observed via the phenomenon known as neutrinoless double beta decay, hitherto unobserved.

I. MAJORANA NEUTRINOS – HOW MASSIVE?

Let us recall that mass terms in the Lagrangean always connect right-handed and left-handed fields. Take the only known neutral leptons of the 1st generation (or family)

$$\nu_{eL}, \nu_{eR}^c = \hat{C}\nu_{eL} \ . \qquad (1)$$

Then introduce a new field $N_R \neq \nu_R^c$ with its charge- (and parity-) conjugate field $N_L^c = \hat{C}N_R$. We now can write a mass term by coupling this new field to the left-handed electron neutrino:

$$-L_{\rm Dirac} = m_D \bar{\nu}_L N_R + {\rm h.c.} \qquad (2)$$

which obviously implies the existence of a 4-component Dirac particle. This term in the Lagrangean conserves lepton number and flavor. Next, however, take the left-handed electron–neutrino field and couple it to its charge-conjugate ν_R^c, and also find a mass term:

$$-L_{\rm Majorana} = \tfrac{1}{2} m_M \bar{\nu}_L \nu_R^c + {\rm h.c.} \qquad (3)$$

This term "creates" or "annihilates " two neutrinos; it defines a 2-component Majorana particle, and violates lepton number and flavor by two units.

Let us revert now to the Standard Model: in the three known generations, neutrino masses are small, and within experimental error compatible with zero. The SM therefore accepts $m_\nu = 0$, thus decreeing there cannot be a helicity flip. This implies that all known evidence for lepton number conservation (notably the fact that reactor-emitted antineutrinos from neutron decay generate only positron tracks in downbeam detectors, interpreted as the conservation of $L = -1$, shared by antineutrinos and positrons) loses its meaning: helicity conservation alone implies the same selection rule. How then are we going to find out whether the neutrino has in fact Dirac or Majorana character? In other words: how do we define the proper field-theoretical description of neutrino mass terms?

For massive neutrinos, Majorana masses present the more general concept: we showed above that it takes a special symmetry to establish Dirac mass terms.[1] As long as the light neutrinos are not truly massless, an overall picture of compatibility of the elementary fermion masses with higher symmetry multiplet concepts is usually (and only loosely) illustrated in terms of a "see-saw" mechanism, which invokes the existence of a heavy neutrino mass (m_N) beyond our present experimental sensitivity for a relation

$$m_\nu m_N \approx m_\ell^2 \ . \qquad (4)$$

Eq. (4) relates the light and heavy neutrino masses to a typical charged-fermion mass in the same multiplet. In the first generation, call the latter mass 1 MeV, and, with m_ν = 1 eV, find 1 TeV for the heavy neutrino mass. A proper treatment, which extends the argument to zero values for the light neutrino mass, includes a full neutrino mass matrix [1].

The question that poses itself today is this: are there clear-cut experimental signals that tell us about the existence of heavy Majorana neutrinos? The classic channel in which to study

[1] A Dirac neutrino can be thought of as a pair of degenerate two-component Majorana neutrinos.

the overall Majorana character has been what is commonly referred to as "neutrinoless double beta decay" of certain nuclei of charge number Z and mass number A,

$$(A, Z) \rightarrow (A, Z+2) + 2e^- \ . \tag{5}$$

It can occur when the appropriate energy levels for neighboring nuclei of charge numbers $Z, Z+1$ energetically forbid single beta decay. It is obviously a very rare process, $\sim G_F^2$, and has been studied in promising isotopes for a number of years, with a great deal of experimental ingenuity [1]. The telling signal that sets it apart from double beta decay accompanied by neutrino emission, is a line spectrum for the total energy of the two emerging electrons. While several instances of the latter process have been observed in recent years, there is no evidence to date for process (5). It has long been realized that the same lepton-number-violating effect could also be observable in electron–electron scattering, where like-sign W pairs would saturate the final state [2]. This effect is a natural for high-energy electron colliders, where the exchanged Majorana neutrinos could be heavy. But detailed calculations, mostly employing left-right-symmetric models, did not lead to promising numbers until we introduced a scenario where two highly polarized e_L beams exchange one of at least two very massive (on the order of 1 TeV) neutrinos N_i, to yield two longitudinally polarized W^- bosons in the final state [3]:

$$e_L^- e_L^- \rightarrow W^-_{\mathrm{Long.}} W^-_{\mathrm{Long.}} \ . \tag{6}$$

P. Minkowski [4] has shown that the occurrence of such massive neutrino states is natural in the decomposition of SO(10) multiplets to single-family left-handed fermion 16plets.

How does all this relate to the mass problem in a quantitative way? Let us consider process (5) first: the relevant quantity in the matrix element for process (5) is

$$\mathrm{Amp} \sim \frac{q m_\nu}{q^2 + m_\nu^2} \ , \tag{7}$$

where q is the neutrino momentum, usually taken to be of order 50 MeV. This makes it clear that for light neutrinos, where $q^2 \gg m^2$, the mass can be neglected in the denominator, whereas the opposite case of heavy neutrinos with $m^2 \gg q^2$ similarly makes the momentum negligible in the denominator:

$$\mathrm{Amp} \sim \begin{cases} m/q \text{ for } q^2 \gg m^2, \\ q/m \text{ for } q^2 \ll m^2 \ . \end{cases} \tag{8}$$

In fact, the present experimental limit on the observation of neutrinoless $\beta\beta$ decay of ^{76}Ge $\rightarrow^{76}$Se, for a half-life $> 5 \times 10^{24}$ years, translates into a light-mass upper limit of about 1 eV. Can the experimental number also be used to set a lower limit on heavy Majorana mass exchanges? In the framework of eq. (8) we could expect to do just that – but I will below caution you from drawing a faulty conclusion along these lines, because the high degree of locality implied by the massive exchange mechanism brings in other, much more restrictive, factors. Let us keep in mind at this point that $\beta\beta_{0\nu}$ decay may be due to either light or heavy neutrino exchanges, but not to massless ones. However, a single measurement will never tell us whether the exchanged mass was large or small.

II. TeV-MASS MAJORANA NEUTRINOS IN ELECTRON–ELECTRON SCATTERING

The experimental breakthrough for the observation of effects that can only be due to the exchange of very heavy Majorana neutrinos may well come with the realization of the next generation of electron colliders: the simple fact that LEP2 is most certainly the highest-energy cirular electron collider that remains economically feasible, and that therefore future colliders at higher CM energy must be linear configurations where two independent accelerators point at the interaction region. This makes it trivial to run these machines in either an e^+e^- or in an e^-e^- mode. Technically, there is no serious problem in using the same layout for interaction region and detector, either, once allowances are made for differences in beam–beam interactions and beam disposal after the interaction [5]. Also, luminosity losses due to the electrostatic repulsion of like-sign beams can probably be made up for by different bunch-train configurations of the incoming beams [6]. The very obvious advantage of operating in the e^-e^- configuration is the possibility, unneeded for a number of asymmetry measurements but of vital importance for the suppression of backgrounds in others, that both beams can be highly polarized. This is necessary for the determination of the desired chiral couplings of a number of rare processes. SLC routinely runs at 80% polarization now with no limitation in luminosity; what is more, polarization can be changed instantaneously. Although the quantitative understanding of the stressed GaAs-surface guns that deliver these electrons upon irradiation by appropriately polarized high-powered lasers is incomplete, recent estimates see no limitation below 91% [7], so that by the time of an NLC turn-on we may well expect that additional edge.

Given the availability of electron–electron collisions with 80% polarization at center-of-mass energies of 0,5 (1.0, 1.5) TeV and with luminosities of 50 (100,100 fb^{-1}y^{-1}), the event numbers to be expected can be read from the cross section expression [3]

$$\sigma(e^-e^- \rightarrow W^-W^-) = \frac{1}{M^2[\mathrm{TeV}^2]} \left(\frac{s}{M^2}\right)^2 \left|\frac{h_\ell}{4\pi}\right|^2 \times 4 \times 10^5 [\mathrm{fb}] \ . \tag{9}$$

Here, M is the reduced mass of the two (lightest, if more than two) heavy neutrinos; h_ℓ is a measure for the mixing of the heavy neutrinos with the incoming electrons, and is experimentally bounded by information from rare lepton decays [3]. This cross section leads, within those bounds, to interesting counting rates for $m_N = 1$ TeV: 0.1 to 30 (2 to 900; 5 to 2,400) events per year, where the two numbers for each energy correspond to the lower and upper limits on h_ℓ, respectively. Since the final states are rather spectacular – back-to-back W^- pairs – there is every assurance that they will stand out above all Standard Model backgrounds [8]. An important capability of our method is a definitive background check: if an observed signal is due to process (6), a change in incoming electron helicity (which decouples the electrons from the W^-, will necessarily make the signal vanish.

Given the fact that this is a rather unique signal, the observation of which could do away with two major conundrums of our

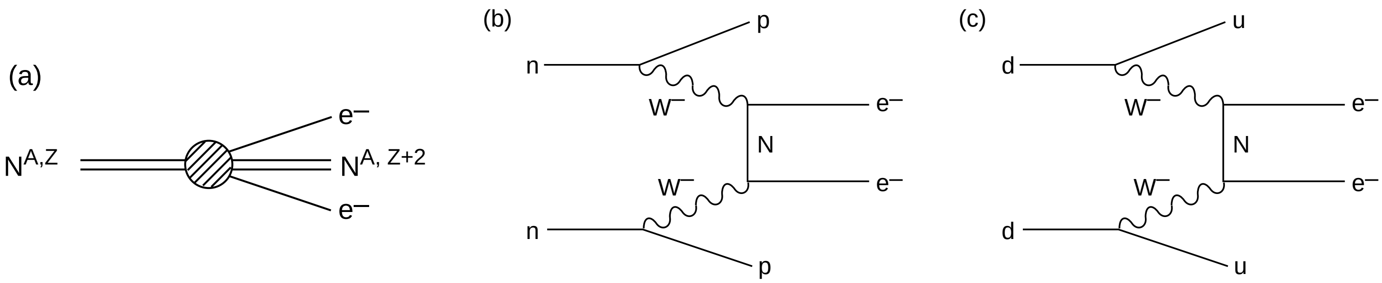

Figure 1: a) Generic graph for neutrinoless double beta decay; b) details of heavy neutrino exchange in a), shown on the nucleon level; c) same as b), but on the quark level. Given that the locality scale is set by m_W^{-1}, this is the most relevant graph.

otherwise astonishingly successful Standard Model, its promise should strengthen the case for the electron linear collider in the TeV range considerably: there is no other method known to us that will give an equally unequivocal answer – not even at energies that permit pair production of these heavy neutrinos, or s-channel production via the charged weak current in ep scattering.

III. COMPATIBILITY WITH EXISTING EVIDENCE FROM NEUTRINOLESS DOUBLE BETA DECAY

In section I, above, we described dependency of the $\beta\beta_{0\nu}$ process on light and heavy neutrino masses [see eq. (8)], in the absence of further constraints. It is unfortunate that much of the relevant literature disregards the very special context within which very heavy exchanges in nuclei have to be examined. As a consequence, the above-mentioned limits on the observation of process (5) in the ^{76}Ge $\rightarrow^{76}$Se system have been taken to imply lower limits on heavy Majorana neutrino masses on the order of 10^3 TeV [9]. It is quite obvious, however, that heavy mass exchanges occur over distances that require calculations (or reasonable estimates) of nucleon- and/or quark-level correlations, as we previously pointed out [10]. In fact, the nuclear double beta decay process shown in Fig. 1a has to be estimated in terms of constraints illustrated by the subprocesses contained in Figs. 1b and c. TeV mass exchanges imply a localization of order 10^{-16}cm; QCD calculations on this scale will influence the quark-level as well as the nucleon-level Hamiltonian densities leading to process (5). Only then will a quantitative comparison of limits on the observation of $\beta\beta_{0\nu}$ with process (6) become meaningful.

In a recent paper, Belanger *et al.* [11] explicitly make the mistake of quantitatively disregarding all the constraints imposed by this strict localization; as a result, they maintain that process (6), illustrated for clarity of comparison in Fig. 2, should be interpreted as the inverse of process (5). If that were the case, the failure of observation of the latter would already exclude all possibilities of observing heavy N_M exchange in e^-e^- scattering at TeV electron colliders, as they illustrate in Fig. 3. P. Minkowski and the author have therefore estimated [12] the effect of all the neglected correlations stemming from nuclear and quark-level factors in the relevant Hamiltonian density operator – stressing the fact that all of these make the definition of the claimed "inverse" relationship incongruous.

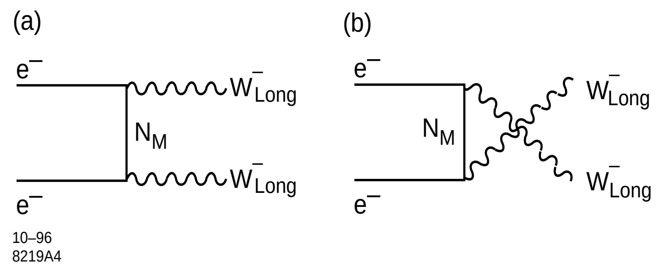

Figure 2: The process $e^-e^- \rightarrow W^-W^-$, mediated by heavy Majorana neutrino exchange.

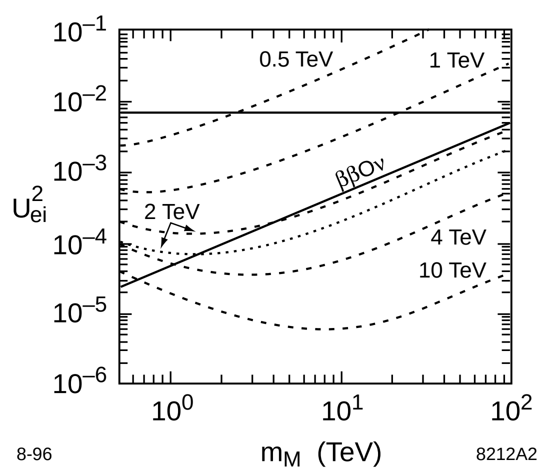

Figure 3: Exclusion plot from Belanger *et al.*, [10]. The dashed curves show the sensitivity of reaction (6) for different e^-e^- CM energies: only the region above them is accessible. The straight line marked $\beta\beta_{0\nu}$ is the (misunderstood) upper limit of what can be seen given the existing limits on the observation of process (5). See the text for its correct interpretation.

Briefly, the argument runs like this: The color Coulomb interaction of the two d quarks will have to happen such that the two final-state electrons emerge in an overall S state, as a spin singlet (to satisfy the Pauli principle). This imposes a spin singlet configuration on the interacting dd pair.

We then need a symmetrical product of space and color SU(3) wave functions, in order to produce an overall antisymmetric wave function for the $l = 0$ dd system. This can be realized only by the **6** configuration of color SU(3). This implies a suppression factor of 2/3.

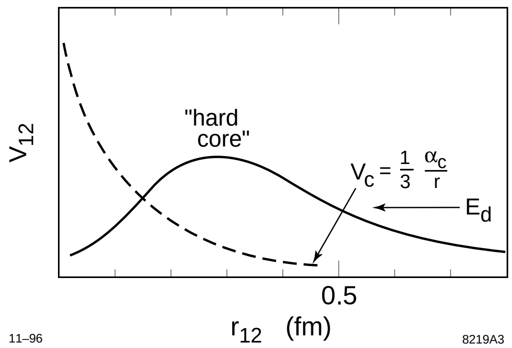

Figure 4: Schematic potential shape for the effective color Coulomb “hard core” that inhibits the free interaction of two d quarks at vanishing relative distances. For details, see text.

Next, we have to attempt to write the strong Hamiltonian density for the quark interactions in Fig. 1c. Let’s do this in terms of a modified hard-core model: Fig. 4 shows how the color Coulomb potential presents a potential barrier while allowing for asymptotic freedom at r_{12} (the relative distance between the two quarks concerned) < 0.1fm. We can do this by writing the overall matrix element in lepton–hadron-factorized form, solving the leptonic part in the conventional way, then using for heavy neutrino exchange the local Hamiltonian density with the hadronic part

$$\begin{aligned} H_q(x) &= \left[\bar{u}^b_\alpha \bar{u}^{c\alpha}(x)\right]\left[d^{c\gamma} d^b_\gamma(x)\right] \\ &= j^+_\mu(x) j^\mu(x) \qquad (10) \\ &\quad \text{with } j^+_\mu(x) = \bar{u}^c \gamma_\mu \frac{1+\gamma_5}{2} d^c \ . \end{aligned}$$

We ordered the fields for this density expression, written at the quark level where our m_W^{-1} localization puts the relevant interaction, such as to stress the incoming and outgoing quark states. The actual calculation with the color Coulomb potential and an appropriate quark momentum is best approximated by a Weizsaecker-Williams method, setting the relative distance r_{12} of the two field operators $\rightarrow 0$. This gives an inhibition (or barrier penetration) factor

$$F_B = e^{-\pi\alpha_s/3} \approx 1/3 \ , \qquad (11)$$

which is not sensitive to the nuclear environment.

The last inhibition factor to be evaluated is due to the “pull” the non-involved quarks in the two interacting neutrons, one u and one d quark each, exercise on the straying companion d which is trying to break up the color singlet nucleons. This restoring force is hard to model in detail, but an ansatz

$$F_{nn} = e^{-2\pi\alpha_s/3} \approx 1/9 \ , \qquad (12)$$

below the value of the correct Clebsch-Gordan coefficient, appears conservative; this is particularly true since the relevant interaction happens at somewhat larger distance, where the full color potential is stronger than what we expect from the r^{-2} level prevailing at fractional fermi distances.

In more traditional terms, the nucleon–nucleon repulsive vector interaction could also be modeled by omega meson exchange; this would lead to an inhibition factor stronger than our estimate above. We therefore feel quite safe multiplying the three explicit factors 2/3, $F_B \sim 1/3$, and $F_{nn} \sim 1/9$ together, for an overall quark-level inhibition factor of

$$F_{dd} = F_6 F_B F_{nn} \approx 1/40 \ . \qquad (13)$$

This entire process, including the relevant leptonic factor which contains the mixing angles between electron and heavy neutrinos, still has to be embedded in the mother-daughter nuclear configurations. That means we feel justified to modify the expectations of nuclear wave function overlaps by the above suppression factor. This applies to all systems that are candidates for neutrinoless double beta decay, in particular to the best-studied such system,

$$^{76}\mathrm{Ge} \rightarrow {}^{76}\,\mathrm{Se}\ e^- e^- \ , \qquad (14)$$

where the present half-life limit is about 5×10^{24} years.

How does this impact on the measurability of heavy neutrino exchange in high-energy electron–electron scattering at the linear colliders we are discussing for the foreseeable future? We illustrate this for two published limits that have seemed to exclude all observability of process (6): Pantis *et al.* [13] quote a limit $(1/m_N)^{-1}_L \gtrsim 6.7 \times 10^3$ TeV This limit now reduces to

$$\frac{(m_N)_L}{|U_{eN}|^2} = 0.025 \times 6.7 \times 10^3 \ \mathrm{TeV}. \qquad (15)$$

The experimental range for the mixing angles U_{eN} is bounded from data on lepton flavor-changing rare decays to values $[2 < |U_{eN}|^2 < 40] \times 10^{-4}$ [3], giving effective limits

$$(m_N)_L > \left\{ \begin{array}{c} 0.67\ \mathrm{TeV} \\ 0.033\ \mathrm{TeV} \end{array} \right\} \text{ for the } \left\{ \begin{array}{c} \text{upper} \\ \text{lower} \end{array} \right\} \text{ limit on } |U_{eN}|^2. \qquad (16)$$

We repeat that the numerical suppression factor is certainly very conservative, putting process (6) squarely in the observable range.

Similarly, let us take another look at Fig. 3: even by today’s standard, electron polarizations are routinely 80%, thus lowering the (dashed) sensitivity curves by a factor 2.5. Combine this with a raising of the line marked “$\beta\beta_{0\nu}$” by the factor of 40 we estimated above, and it becomes quite obvious that the entire parameter range discussed for the discovery process for TeV-mass Majorana neutrinos becomes generously accessible.

IV. CONCLUSION

I would be more than happy to make a completely positive statement on the experimental chances that a good deal of our Standard Model's failings in the neutral lepton sector will be brought into a more consistent focus by experimentation at linear electron colliders of the next generation. The strategy for this search that we conjure up cannot *predict*. But in the context of the most reasonable group-theoretical scenarios and of the heavy/light neutrino mixing parametrizations that are compatible with other present evidence, our scenario must appear very attractive indeed. Having an experimental method at hand that can unambiguously probe for heavy Majorana masses, must be seen as a unique capability of the electron–electron incarnation of the NLC: it does not exist within our reach elsewhere. Among many attractive contributions that electron–electron collisions are liable to make to essentially all of the physics concerns of the NLC program, this project must stand out for its importance.

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