MAJORONS FROM DOUBLE BETA DECAY∗
C.P. BURGESS and JAMES M. CLINE†

Physics Department, McGill University, 3600 University Street
Montréal, Québec, Canada H3A 2T8

ABSTRACT

We explain the existence of excess events near the endpoints of the double beta decay (ββ) spectra of several elements, using the neutrinoless emission of massless Goldstone bosons. Models with scalars carrying lepton number −2 are proposed for this purpose so that ordinary neutrinoless ββ is forbidden, and we can raise the scale of global symmetry breaking above the 10 keV scale needed for observable emission of conventional Majorons in ββ. The electron spectrum has a different shape, and the rate depends on different nuclear matrix elements, than for the emission of ordinary Majorons.

Double beta decay (ββ) is an extremely rare process in which two neutrons in a nucleus simultaneously decay into protons. It has now been observed in seven elements: 76Ge, 82Se, 100Mo, 128Te, 130Te, 150Nd and 238U. One reason ββ has received so much attention is the possibility of detecting small Majorana masses for the neutrinos if the neutrinoless mode ββ0ν should be observed. The amplitude for ββ0ν is proportional to an effective Majorana mass, $m_{\text{eff}} = \sum \theta_i^2 m_i$, where $\theta_i$ is the mixing angle between $\nu_e$ and the $i$th mass eigenstate. Experimentally it is known that

$$m_{\text{eff}} < 1 \text{ eV}.$$  

It is also possible that ββ0ν proceeds by the annihilation of the virtual neutrinos, through a vertex

$$g \varphi \bar{\nu}_e \gamma_5 \nu_e,$$

where a scalar particle $\varphi$ is emitted. This can naturally occur if lepton number is broken spontaneously, so that neutrinos couple to a massless Goldstone boson called the Majoron. The sum-energy electron spectrum for this process, which we refer to as $\beta\beta_M$, is skewed toward higher electron energies than that for $\beta\beta_{2\nu}$, because it is a three-body decay. Previous experiments searching for this mode obtained a limit of $g < 2 \times 10^{-4}$, following indications of a positive signal at the $3 \times 10^{-4}$ level.

It is therefore intriguing that the UC Irvine group has for the last few years been seeing excess events near the endpoint in three different elements: 82Se, 100Mo and 150Nd. A similar excess is also plainly visible in the 76Ge spectrum of Avignone et al. In all four

∗ Talk given at Beyond the Standard Model III, Ottawa, 1992
† speaker
elements, the excess events become evident starting from 0.5 to 1 MeV below the endpoint, and constitute between 2 and 3% of the total signal. Because the expected signal from $\beta\beta_{2\nu}$ is negligible near the endpoint, the chance of these events being due to statistical fluctuations is only 1 in $10^5$.

Let us write the partial rate for $\beta\beta_M$ as

$$d\Gamma_{0\nu,M} = 2\pi^{-5}(g_AG_F)^4g_{\text{eff}}^2|\mathcal{M}|^2d\Omega,$$

(3)

where $g_{\text{eff}}$ is a model-dependent effective coupling, $\mathcal{M}$ is the usual combination of Gamow-Teller and Fermi nuclear matrix elements, and $d\Omega$ is the phase space. We have analyzed the data by matching the number of excess events above some threshold energy where visually the anomalies seem to begin for each element. For $^{238}\text{U}$ and the Te isotopes, only the total rate is observed, so we omit thresholds for these elements to see how large a Majoron coupling would be needed for $\beta\beta_M$ to saturate the total observed signal.

The results are shown in Table 1. We find that the required values of $g_{\text{eff}}$ are in the range $5 \times 10^{-5}$ to $2 \times 10^{-4}$ except for $^{100}\text{Mo}$, which needs a somewhat larger coupling. This presents a problem for the ratio of the rates of $^{130}\text{Te}$ to $^{128}\text{Te}$ decays, which has recently been measured$^7$ to be $(2.41 \pm 0.06) \times 10^3$. The $\beta\beta_M$ prediction implied by Table 1 (and the phase space calculations not shown there) would be much smaller, 93, assuming $g_{\text{eff}} = 1 \times 10^{-4}$ so that the endpoint anomalies could be explained.

Table 1: couplings for emission of ordinary Majorons in double beta decay, assuming total rate of geochemically observed decays was saturated by $\beta\beta_M$, and using the nuclear matrix elements of ref. [8].

| Element | $^{76}\text{Ge}$ | $^{82}\text{Se}$ | $^{100}\text{Mo}$ | $^{150}\text{Nd}$ | $^{128}\text{Te}$ | $^{130^*}\text{Te}$ | $^{238}\text{U}$ |
|---------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $g_{\text{eff}}$ | $1 \times 10^{-4}$ | $8 \times 10^{-5}$ | $4 \times 10^{-4}$ | $8 \times 10^{-5}$ | $4 \times 10^{-5}$ | $1 \times 10^{-4}$ | $2 \times 10^{-4}$ |

Even if we ignore the Te problem, there is a more serious difficulty: no Majoron models exist which seem able to give such a large coupling. The favored theory for $\beta\beta_M$ was the triplet Majoron model$^{2,9}$ but it has been ruled out by LEP’s measurement of the invisible $Z$ width. Alternatively, the singlet Majoron model$^{10}$ has couplings which are highly suppressed, since generically

$$g_{\text{eff}} \sim m_{\text{eff}}/v,$$

(4)

where $v$ is the scale of lepton number breaking. The most natural choice for $v$ is of order the weak scale, leading to $g_{\text{eff}} \sim 10^{-11}$. But we need $v < 10$ keV to explain the size of the $\beta\beta$ anomalies, given the limit (1) on $m_{\nu}$. This unnaturally small value requires extreme fine tuning to keep it so far below the weak scale.* However, suppose we could remove the relation between $v$ and $m_{\text{eff}}$ by maintaining lepton number as an exact symmetry, thus insuring $m_{\text{eff}} = 0$ to all orders. Perhaps we would be able to have a higher scale of breaking for whatever symmetry it is that gives us the Goldstone bosons. But if lepton number is conserved, then the Majorons must carry $-2$ units of $L$. We will call these charged Majorons since they carry the global lepton charge.

* Moreover, even if $v \sim 10$ keV, one must introduce some neutrinos heavier than the 100 MeV scale of the nucleon Fermi momentum; otherwise the mixing of the light neutrinos is ineffective and $g_{\text{eff}}$ becomes the bare coupling of the Majoron to $\nu_e$, which is zero for singlet Majorons.
This situation reminds us of the standard model, where a would-be Goldstone boson, the longitudinal component of $W^-$, has electric charge. It is therefore quite easy to construct a model of charged Majorons; we just mimic the standard model by introducing an SU(2)$_s \times$U(1)$_L'$ global symmetry that acts on the leptons and on some new sterile neutrinos and Higgs multiplets. Although we could give a renormalizable model of this sort, it is clearest to present an effective Lagrangian with the dimension-five operators

$$
\frac{g}{\Lambda}(L_e H)(1 + \gamma_5)(\phi^\dagger s) + \frac{h}{\Lambda}(\phi^\dagger s)(1 + \gamma_5)(\bar{\phi}^\dagger s) + \text{h.c.},
$$

where $\bar{\phi} \equiv \tau_2 \phi^*$. These particles transform as $L_e \sim (1, 1, 2, -1/2)$, $s \sim (2, 0, 1, 0)$, $\phi \sim (2, 1, 1, 0)$ and $H \sim (1, 1, 2, -1/2)$ under the SU(2)$_s \times$U(1)$_L' \times$SU(2)$_L \times$U(1)$_y$ symmetries. When $\phi$ gets a VEV $v$, a Dirac mass term between $s_1$ and $s_2$ results. $\nu_e$ is mostly a massless state $\nu$, but mixes with a heavy Dirac neutrino $\Psi$ with mixing angle $\theta = \tan^{-1}(g \langle H \rangle / hv)$. One finds that the charged Majorons couple to the neutrinos via

$$
\frac{1}{4v} \left( \sin \theta \bar{\nu} - \cos \theta \bar{\Psi} \right) (1 - \gamma_5) \frac{\partial}{\partial \phi} \bar{\nu} + \text{h.c.}
$$

They have charge $-2$ under the residual lepton symmetry which is given by $L = L'$ plus the diagonal generator of SU(2)$_s$.

We note that the sum-energy electron spectrum for emission of charged Majorons differs from that of ordinary Majorons. Consider singlet Majorons: in the language where they are derivatively coupled ($v^{-1} \bar{\nu} \partial \phi \gamma_5 \nu$), there exist extra Feynman diagrams with Majoron Brehmsstrahlung off the electron lines via the vector coupling $v^{-1} \bar{\phi} \frac{\partial}{\partial \phi} \bar{\phi}$, while the neutrinos annihilate through their Majoron mass. These survive at zero Majoron momentum $q$ because the internal electrons are going on shell as $q \to 0$. But these graphs don’t exist for charged Majorons because there are no Majorona masses; hence the amplitude vanishes as $q \to 0$, and consequently the sum energy spectrum of the electrons has a shape intermediate between that of $\beta \beta_{2\nu}$ and generic $\beta \beta_M$ spectra. Our spectrum behaves like $(Q - E)^3$ near the endpoint energy $Q$, compared with $(Q - E)^5$ for $\beta \beta_{2\nu}$ and $(Q - E)$ for ordinary $\beta \beta_M$. The different spectra are shown in Fig. 1.

Furthermore $\beta \beta_M$ with charged Majorons depends on different nuclear matrix elements than the usual ones, because the amplitude contains an extra factor of the neutrino momentum $p$. The leptonic part has the Lorentz structure $\bar{e}_L \gamma_{\mu}[\bar{\nu}_L \gamma_5 \nu]_{\nu e}$, and for kinematic reasons only the spatial components of $p$ contribute significantly. Since they have odd parity, but the nuclear transition is $0^+ \to 0^+$, we require nuclear operators with odd parity also, unlike the Gamow-Teller and Fermi terms. Such operators come from the recoil corrections to the hadronic weak currents, as well as p-wave Coulomb corrections to the electron wave functions. For example, we have matrix elements of the form

$$
\mathcal{M}_R = (g_\nu / g_A) \langle p^2 \rangle^{-1} \left( \hat{r} \cdot (D_n \times \sigma_m + \sigma_n \times D_m) \frac{\partial}{\partial r} h(r) \right),
$$

using the notation of Doi et al.\textsuperscript{11} or Tomoda.\textsuperscript{12}

\textsuperscript{*} For example, introduce sterile neutrinos $N_1$ and $N_2$ having U(1)$_{L'}$ charges +1 and -1 respectively, and include all their allowed couplings to $L_e$, $H$, $\phi$ and $s$. Integrating out $N_i$ results in the effective theory. Even if $N_i$ is no heavier than the remaining particles, the resulting prediction is quantitatively similar to that of the renormalizable theory.
We find our process is roughly comparable to the ordinary Majoron process with a coupling of
\[ g_{\text{eff}} \sim (M/2v)\theta^2 \frac{\langle qp \rangle M^2}{\left(\langle p^2 \rangle + M^2\right)^2} \frac{\mathcal{M}_{\text{new}}}{\mathcal{M}}, \]  
where \( M \) is the mass of the heavy sterile neutrino, \( \theta \) is its mixing angle to \( \nu_e \), and \( \langle qp \rangle \) is of order 100 MeV since the Majoron only has energy \( q \sim 1 \) MeV, whereas the virtual neutrino momentum is \( \sim 100 \) MeV. \( \mathcal{M}_{\text{new}} \) denotes the new nuclear matrix elements associated with charged Majorons. Note that \( M \sim v \) is also the scale at which our global symmetry breaks. To get \( g_{\text{eff}} \) as large as possible we need \( \theta \sim 0.1 \), which forces us to take \( M > 500 \) MeV due to the constraints from peak searches in meson decays.\(^{13,14}\) With these values, we would also need the new nuclear matrix elements (some of which have not been computed, to our knowledge) to be 25 times larger than the usual ones. Another possibility is that the coupling producing the heavy neutrino mass is very strong (\( M > v \)), and that it decays so quickly into light neutrinos and Majorons that it is really a rather broad resonance. Then its mass might be closer to 100 MeV and yet escape notice in \( \pi \) and \( K \) decay peak searches, even with a large mixing angle.

More exactly, the rate for charged Majoron emission is,
\[ d\Gamma_{\text{M}} = 2\pi^{-5}(g_A G_F \theta)^4(M/2v)^2|\mathcal{M}_{\text{new}}|^2d\Omega_{\text{new}}. \]  

The detailed forms of \( \mathcal{M}_{\text{new}} \) and \( d\Omega_{\text{new}} \) are given in ref. [15]. Some of the contributions to \( \mathcal{M}_{\text{new}} \), such as the one displayed in eq. (7), have been calculated. In Table 2 we compare the value of \( \mathcal{M}_{\text{new}} \) needed to account for the anomalous events with the computed value\(^{16}\) of \( \mathcal{M}_R \), one of the actual contributions to \( \mathcal{M}_{\text{new}} \). (For this comparison we have assumed a large mixing angle and strong coupling, as indicated in the table caption.) The agreement is remarkably good except for \(^{128}\)Te. However, the absolute rate \(^{128}\)Te is known less well than the ratio to \(^{130}\)Te, and our model predicts \( \Gamma(\text{^{130}Te})/\Gamma(\text{^{128}Te}) = 770 \), in much better agreement with the experimental value quoted above than in the singlet Majoron model.

In summary, we have presented a model of Goldstone boson emission in double beta decay which may be able to explain excess events near the endpoints of several elements. The boson has lepton number \( L = -2 \) and the light neutrinos remain exactly massless so
Table 2: matrix elements for charged Majoron emission in double beta decay. We give the values of the matrix elements which are needed to explain the data (assuming all events were $\beta\beta_M$ for $^{238}\text{U}$ and Te) for $\theta^2(\langle p^2\rangle/M^2)/(\langle p^2\rangle+M^2)=3\times10^{-2}$, and values of some representative matrix elements that have been calculated.\textsuperscript{16}

| Element | $^{76}\text{Ge}$ | $^{82}\text{Se}$ | $^{100}\text{Mo}$ | $^{150}\text{Nd}$ | $^{128}\text{Te}$ | $^{130}\text{Te}$ | $^{238}\text{U}$ |
|---------|------------------|------------------|------------------|------------------|------------------|------------------|------------------|
| $|\mathcal{M}_{\text{new}}|_{\text{needed}}$ | 1.1 | 0.99 | 1.4 | 1.3 | 0.32 | 0.48 | 0.61 |
| $|\mathcal{M}_R|_{(\text{Muto et al.)}}$ | 1.1 | 0.95 | 1.1 | 1.3 | 0.92 | 0.78 | ? |

that ordinary neutrinoless double beta decay is forbidden. The electron spectrum for these double beta decays with “charged Majoron” emission is less peaked toward high energies than that for ordinary Majorons. The model requires a sterile Dirac neutrino with a large ($\sim 0.1$) mixing angle and hence (due to experimental constraints) a mass $M \sim 500$ MeV. Moreover, the $\beta\beta_M$ rate in this case depends on unknown nuclear matrix elements $\mathcal{M}_{\text{new}}$, different from those appearing in the usual amplitudes for $\beta\beta_M$, which must turn out to be larger than the usual ones in order for the rate to be as large as is seemingly observed.

We warmly thank M. Moe and F. Avignone for information about their experiments, E. Takasugi and T. Kotani for helpful correspondence, W. Haxton, M. Luty and P. Vogel for valuable discussions, and R. Fernholz for producing plots of the spectra. This work was supported in part by the Natural Sciences and Engineering Research Council of Canada and les Fonds F.C.A.R. du Québec.

References

1. M. Doi, T. Kotani and E. Takasugi, Phys. Rev. D\textbf{37} (1988) 2572.
2. H.M. Georgi, S.L. Glashow and S. Nussinov, Nucl. Phys. B\textbf{193} (1981) 297.
3. P. Fisher et al., Phys. Lett. B\textbf{192} (1987) 460; D.O. Caldwell et al., Phys. Rev. Lett. 59 (1987) 1649.
4. F.T. Avignone III et al., in Neutrino Masses and Neutrino Astrophysics, proceedings of the IV Telemark Conference, Ashland, Wisconsin, 1987, edited by V. Barger, F. Halzen, M. Marshak and K. Olive (World Scientific, Singapore, 1987), p. 248.
5. M.K. Moe, M.A. Nelson, M.A. Vient and S.R. Elliott, preprint UCI-NEUTRINO 92-1 (1992).
6. F.T. Avignone III et al., Phys. Lett. B\textbf{256} (1991) 559.
7. W. Haxton, quoting experimental results of Bernatowicz and Holenberg, Neutrino ‘92, Granada, Spain. It was Haxton who first used the Te decays in order to constrain Majoron couplings.
8. A. Staudt, K. Muto and H.V. Klapdor-Kleingrothaus, Europhys. Lett., 13 (1990) 31.
9. G.B. Gelmini and M. Roncadelli, Phys. Lett. 99B (1981) 411.
10. Y. Chikashige, R.N. Mohapatra and R.D. Peccei, Phys. Lett. 99B (1981) 411.
11. M. Doi, T. Kotani and E. Takasugi, Prog. Theo Phys. Suppl. 83 (1985) 1.
12. T. Tomoda, Rept. Prog. Phys. 54 (1991) 53.
13. R.E. Shrock, Phys. Rev. D\textbf{24} (1981) 1232; T. Yamazaki et al., Proceedings of the XIth International Conference on Neutrino Physics and Astrophysics, eds. K. Kleinknecht and E.A. Paschos (World Scientific, Singapore, 1984), p. 183.
14. D.I. Britton et al., Phys. Rev. Lett. 68 (1992) 3000.
15. C.P. Burgess and J.M. Cline, McGill preprint 92-22 (1992).
16. K. Muto, E. Bender and H.V. Klapdor, Z. Phys. A 334 (1989) 187, as tabulated in ref. 12.