Oscillations of Dirac Neutrinos coupled to a Light Higgs Field and the Short Baseline Anomalies

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1 Motivation

The nature of neutrino mass is widely recognized as one of the most important open theoretical and experimental questions in particle physics. The masses of most particles in the Standard Model (SM) are understood to be generated via the Higgs mechanism, a claim which is is now supported experimentally by both the complete particle content of the SM electroweak and SM Higgs sectors, and also measurements of the mass-dependent coupling between many fermions and Higgs bosons at the LHC. Within the standard model, the structure of the electroweak sector relates the masses of the W, Z and Higgs boson masses to electroweak couplings and the two parameters λ_H, ν_H which define the Standard Model Higgs potential. Yukawa couplings to fermions generate masses $m_i = g_i v_H$ after symmetry breaking. They have experimentally measured values spanning from 7×10^{-1} (top quark) to 2×10^{-6} (electron). Their distribution appears consistent with random sampling from a scale-invariant distribution - that is, if the weighting distribution is assumed to be of the form dm/m^{δ}, the best fit exponent is $\delta = 1.02 \pm 0.08$, where $\delta = 1$ would be complete scale invariance [5].

Neutrinos present a critical exception to this rule. The masses of the neutrinos are not yet known, but they are all certainly too small to be consistent with this scale invariant distribution. That neutrinos have mass at all was inferred from oscillation experiments, which measure squared mass splittings $\Delta m^2_{21} \sim 7.37 \times 10^{-5}\mathrm{eV}^2$ and $\Delta m^2_{23} \sim 2.5 \times 10^{-3}\mathrm{eV}^2$. To know the full mass spectrum, two further pieces of information are required: the absolute mass scale (or lightest neutrino mass m_1), and the mass ordering. Upper limits on the mass scale from tritium beta decay give $m_{\bar{\nu_e}} < 2.05\,\mathrm{eV}$ from the Troitzk experiment and $m_{\bar{\nu_e}} < 2.3\,\mathrm{eV}$ from the Mainz experiment, with possible sensitivity as low as $m_{\bar{\nu_e}} \sim 0.2\,\mathrm{eV}$ expected from the Katrin experiment. Cosmological observations of the CMB and BAO suggest the sum of neutrino masses is still smaller, $\sum_i m_i < 0.23\,\mathrm{eV}$. One neutrino mass state may still be massless $m_{\nu 1} \geq 0$, but oscillations place lower limits on the other two masses of $m_{\nu 2} > 8.5\,meV$ and $m_{\nu 3} > 50\,meV$.

The inconsistency of neutrino masses with the other particles suggests another mechanism is responsible for neutrino mass. Because neutrinos are unique in the SM in carrying no conserved gauge charges, they can admit a Majorana mass term. This is of significant theoretical interest, not least because it can be generated under symmetry breaking of the Weinberg operator, only dimension-five operator consistent with the symmetries of the SM. If the Standard Model is to be understood as a low energy effective theory with corrections from higher energy scales, the effects arising from this operator are suppressed by only one power of the new (presumably high) energy scale Λ, whereas all other effects would be suppressed by more. The Weinberg operator can generate neutrino masses through the Seesaw mechanism, with the experimentally observable Majorana neutrino masses m_ν being related to the (SM-scale) neutrino Dirac masses m_D and the new energy scale Λ by:

$$m_\nu = \frac{m_D^2}{\Lambda} \tag{1}$$

where m_D is assumed to be a "natural" scale Dirac mass from a SM Higgs coupling and $\Lambda$is a large new energy scale, thus explaining the smallness of neutrino mass relative to the other fermions. Although the Majorana picture appears a natural explanation, the Majorana or Dirac nature of the neutrino is not yet known, and no evidence of this larger mass scale Λ yet exists.

The possibility remains that neutrinos are Dirac particles, without a Majorana mass term. The challenge in this case is explaining why the Yukawa coupling to neutrinos are so much smaller than to the other particles. this would be particularly strange since the neutrinos are members of SU(2) doublets with the charged leptons, all of which have much larger masses. Allowing Dirac neutrinos suggests there must be some nontrivial SU(2) structure in the Higgs sector. This is often codified in two-Higgs-doublet models [3]. In particular, some authors have suggested that new, light Higgs fields with a smaller vacuum expectation value could generate the light neutrino masses, while the standard model Higgs generates the masses of the other Fermions [12, 7, 2]. In this work we consider phenomenological implications of such a scheme.

We assume a minimal model with a single new Higgs boson with a standard-model type potential:

$$L_{\nu Higgs} = \epsilon^2 \Phi^\dagger \Phi + \zeta \left(\Phi^\dagger \Phi\right)^2 \qquad \epsilon^2 < 0 \quad \zeta > 0 \tag{2}$$

Generating a new Higgs mass and VEV:

$$v_h = \frac{|\epsilon|}{\sqrt{\zeta}} \quad > 0 \tag{3}$$

$$m_h = \sqrt{2\zeta} v_h \quad > 0 \tag{4}$$

The masses of the neutrinos are then given, as usual, by Yukawa couplings to new Higgs VEV:

$$L_{\nu mass} = \sum_i \left(g_i v_h\right) \bar{\nu}\nu \tag{5}$$

Both the Higgs mass and vacuum expectation value are independent free parameters. To generate natural neutrino masses using Yukawa couplings of similar order to those in the standard Higgs sectors, the VEV of the new Higgs field should be in the range $10\, meV < v_h < 100\, keV$. It is worth noting that the "natural" range of Yukawas in this new sector may be entirely unrelated to that in the known Higgs sector, so this range should be treated only as a rough guide, and naturalness (as usual) cannot be argued rigorously. With the VEV established, the mass of the new Higgs boson is an independent and unconstrained parameter which is, in principle, experimentally measurable. If prejudices developed from the SM Higgs sector were directly applicable, the Higgs self coupling ϵ would be $O(1)$ and the new Higgs mass would be of the same order as the new VEV. Again, this should only be taken as a rough guide. Finally we note that the neutrinos in our model are strictly Dirac fermions with no Majorana term present, and we do not rely on the seesaw to explain the low mass scale.

A particularly interesting phenomenological implication of this model is that exchange of this new boson with the cosmic microwave background can, under some circumstances, generate a resonance observable in neutrino oscillation experiments. In particular, for a Higgs mass of O(10keV), this resonance may be observable in short baseline neutrino oscillation experiments. In fact, for some parameter points, a signature very similar to the MiniBooNE low energy excess may be expected. This requires a local over-density of the cosmic microwave background which is large, but not experimentally excluded. In the following sections we elaborate on this mechanism.

2 New Higgs exchange with the Cosmic Neutrino Background

We consider a beam of relativistic neutrinos with a narrow four momentum spread centered at $\{E, p\}$ as test-particles. The new Higgs boson introduces a new force between the test particles and other neutrinos, but does not couple to any other Standard Model particle. Interactions between the test particles and the all-pervasive bath of relic big bang neutrinos can introduce new and measurable effects via exchange of this new boson.

The Higgs couplings to neutrinos are determined by the Yukawa coupling. Were Dirac neutrino masses generated by a standard model Higgs and no Majorana term, this Yukawa coupling would be of order $10^{-15} < g < 10^{-13}$, which is why standard Higgs-neutrino interactions are always negligible in our calculation. In fact, the effect of the SM Higgs is far sub-dominant to the effects of Z exchange. The effects of Z exchange between neutrinos and the relic neutrino background have been studied [4]. The new oscillation effects introduced are negligible in almost all cases of experimental interest. In our model, on the other hand, the VEV of the Higgs is much smaller, and so the Yukawa coupling used to generate the neutrino masses are much bigger. This means that exchange of Higgs bosons with the relic neutrino background may no longer be strongly suppressed.

The effect we consider here coherent forward scattering off the cosmic neutrino background via light Higgs exchange. In this process, no momentum is transferred and the scattered neutrino wave interferes coherently with the incident unscattered wave, generating a refractive effect. Because this refractive effect depends on the mass composition of the beam in a non-flavor-diagonal manner, it introduces new oscillation physics.

The relevant Feynman diagrams are shown in Figure 1. We label each with an amplitude $\mathcal{M}^a_{ij}(p, k)$ where the index a specifies whether a background ν or $\bar{\nu}$ is involved, p labels the four-momentum of the test neutrino with mass Eigenstate i, and k labels the four-momentum of the relic neutrino or antineutrino with mass eigenstate j. We restrict our consideration to neutrino beams only, but exactly analogous effects are present for antineutrinos.

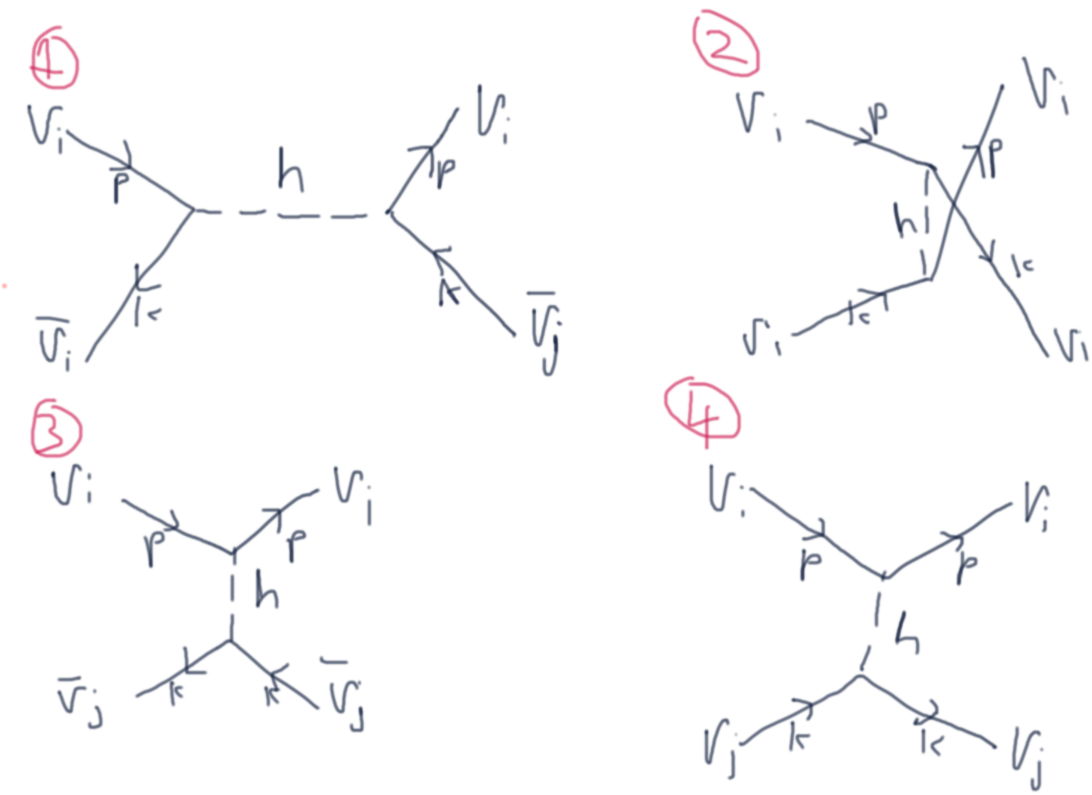

Figure 1: Feynman diagrams of Higgs exchange between a test neutrino of four-momentum p and flavor i and the CNB neutrino or antineutrino with momentum k and four-momentum k.

To maintain coherence with the incident beam we are limited to interactions where the incoming and outgoing momenta and masses of both the test particle and the relic neutrino are unchanged by the interaction. Given that the mass basis of the neutrinos is defined by diagonally of the Higgs coupling in our model, it is clear that diagrams 1 and 2 only give nonzero contributions when $i = j$

Applying the Feynman rules, allowing for a finite Higgs width, and assuming the CNB neutrinos to be at rest and unpolarized, we find the spin-averaged matrix elements:

$$i\langle\mathcal{M}_{ij}^{\bar{\nu},1}\rangle = i\frac{g_i^2\left[k.p - m_i^2\right]}{(p+k)^2 - m_h^2 - im_h\Gamma_h}\delta_{ij} \tag{6}$$

$$i\langle\mathcal{M}_{ij}^{\nu,2}\rangle = i\frac{g_i^2\left[k.p + m_i^2\right]}{(p-k)^2 - m_h^2 - im_h\Gamma_h}\delta_{ij} \tag{7}$$

$$i\langle\mathcal{M}_{ij}^{\bar{\nu},3}\rangle = 2i\frac{g_i^2 g_j^2 \nu_h^2}{m_h^2 - im_h\Gamma_h} \tag{8}$$

$$i\langle\mathcal{M}_{ij}^{\nu,4}\rangle = -2i\frac{g_i^2 g_j^2 \nu_h^2}{m_h^2 - im_h\Gamma_h} \tag{9}$$

To account for the effects of the full relic background, we must sum over background momentum and flavor distribution. This distribution is determined by a density function f_a^i, which in the standard cosmological model is the Fermi Dirac function, redshifted from decoupling time to a temperature $T \sim 1.95\,\mathrm{K}$ in the present era. F_i is a normalization constant, such that [8]:

$$f_a^i(p,T) = \frac{F_i}{1 + e^{(E+\mu_a)/kT}} \tag{10}$$

$$\int d^3k\, f_a^i(k) = N_i^a V \tag{11}$$

Where N_i is the number density of background species and V is a finite normalization volume. In the absence of a chemical potential driving an over-density of either neutrinos or antineutrinos in the early universe, $\mu_i = 0$ and the distributions are equivalent for ν and $\bar{\nu}$. In the absence of significant gravitational clustering, the number density of all mass states are equivalent, with $\langle n_i \rangle = 56\,\mathrm{cm}^{-3}$ ($4.3 \times 10^{-13}\mathrm{eV}^3$ in natural units), or for each flavor in each of neutrinos and antineutrinos.

Deviations from this baseline model are expected from the effects of gravitational clustering of non-relativistic neutrinos [11]. Although [11] predicts only small over-densities, it has been suggested in the context of direct detection searches [6, 10] that the over-density may be as large as 10^6, in order to match the local baryon over-density. Over-densities as large as 10^9 have been discussed in connection with direct CNB detection searches [9], and these remain outside experimental limits. Finally, over-densities of 10^{13} have been suggested in order to explain the knee of the cosmic ray spectrum. Here we parameterize the local over-density through the dimensionless number $\eta = n_i/\langle n_i \rangle$.

With equal number densities of relic neutrinos as antineutrinos, the contributions of Feynman diagrams 2 and 3 cancel in the forward scattering amplitude. We assume this case to simplify our model, although in practice the effects of any asymmetry, if present, would be highly sub-leading in our final expressions.

A question that will be of some importance is whether the neutrinos of the relic background can be assumed to be mostly relativistic, $k \gg m_i$, or mostly non-relativistic, $k \ll m_i$. At temperatures of 2K, the energies of the relic neutrinos are in the $0.1 - 1$ meV range. Thus all except perhaps the lightest are non-relativistic. The dominant effects in our model will be caused by scattering from the species with the largest Yukawa coupling, which are the heaviest neutrinos, and these are necessarily non-relativistic. This allows us to write the four vector $k = (m, 0)$ and replace the sum over the Fermi Dirac distribution with a simple multiplication by number density. With all these considerations, the only important scattering amplitude comes from diagrams (1) and (2), with each background neutrino and antineutrino of the same mass as the test particle giving a contribution:

$$\langle \mathcal{M}_i^\nu \rangle = \delta_{ij} g_i^2 \left[\frac{E m_i}{2Em_i - m_h^2 + i m_h \Gamma_h} + \frac{E m_i}{-2Em_i - m_h^2 + i m_h \Gamma_h} \right] \tag{12}$$

Where it has been assumed that $E \gg m_i$. If we further make the assumption that the only allowed decay mode of this new boson is to neutrinos, we may substitute in the width:

$$\Gamma_i = \frac{1}{2} m_h \sum g_i^2 \tag{13}$$

To yield real and imaginary parts of the scattering matrix:

$$Re\langle M \rangle = \delta_{ij} \frac{g_i^2 E}{2m_i^2} \left[\frac{E - E_{res}}{(E - E_{res})^2 + \left(\frac{1}{2} E_{res} g_i^2\right)^2} + \frac{-E - E_{res}}{(E + E_{res})^2 + \left(\frac{1}{2} E_{res} g_i^2\right)^2} \right] \tag{14}$$

$$Im\langle M \rangle = -\delta_{ij} \frac{g_i^4 E}{4m_i^2} \left[\frac{E_{res}}{(E - E_{res})^2 + \left(\frac{1}{2} E_{res} g_i^2\right)^2} + \frac{E_{res}}{(E + E_{res})^2 + \left(\frac{1}{2} E_{res} g_i^2\right)^2} \right] \tag{15}$$

Where $E_{res} = m_h^2/2m_i$. It is also instructive to consider also a more limited model invoking the Zero Width Approximation (ZWA), which involves setting $\Gamma \to 0$ in eq. 12 to yield:

$$Re\langle M_{ZWA} \rangle = \delta_{ij} \frac{g_i^2}{4m_i} \left[\frac{-E_{res}}{E^2 - E_{res}^2} \right] \tag{16}$$

$$Im\langle M_{ZWA} \rangle = 0 \tag{17}$$

In both the ZWA and finite-width cases, a clear resonance is observed at $E = E_{res}$, corresponding to resonant production of new Higgs bosons at rest in the center of mass frame. Although this resonance only has a very small effect in terms of real particle production, it makes a significant contribution to the real part of the forward scattering amplitude, and thus contributes an additional oscillation phase near resonance.

In all parameter space considered here, the effects of the imaginary part will be negligible, with all observable phenomenology being contained in the real part, which affects the propagation phase in a non-negligible way, creating a new form of neutrino oscillation.

3 Connection to neutrino refractive properties and oscillations

The scattering matrix calculated above has both real and imaginary parts, which contribute to the refractive and absorptive behaviors of the neutrino beam, respectively. It will be most convenient in what follows to work with the T matrix normalized with single-particle wave functions, rather than the Lorentz-invariant M matrix. These are related by $T = M/4Em_i$.

For relativistic scattering we can incorporate the effects of the forward scattering amplitudes calculated above to the neutrino refractive properties via the modification:

$$\psi = u_0 Exp\left[ipx\right] \rightarrow \psi' \tag{18}$$

$$\psi' = u_0(p) Exp\left[i\left(p + in_iT_i\right)x\right] \tag{19}$$

$$= u_0(p)\, e^{i(p+nRe[\langle M\rangle/4Em_i])x}\, e^{-nIm[\langle M\rangle/4Em_i]x} \tag{20}$$

Although usually described in terms of matter potentials, it is easy to verify that this treatment reproduces the standard MSW effect if M is chosen to be that for the weak interactions. In our case, the matrix elements of interest are given by 12.

Following the standard derivation of the neutrino oscillation formula, we find the probability for conversion from a neutrino flavor α to neutrino flavor β is given by:

$$P_{\alpha\beta} = \delta_{\alpha\beta} - 4\sum_{i>j} Re(U^*_{\alpha i}U_{\beta i}U_{\alpha j}U^*_{\beta j})sin^2\left[\frac{\Delta_{ij}}{2}L\right] + 2\sum_{i>j} Im(U^*_{\alpha i}U_{\beta i}U_{\alpha j}U^*_{\beta j})sin^2\left[\frac{\Delta_{ij}}{4}L\right] \tag{21}$$

With:

$$\Delta_{ij} = (p_i - p_j) + nRe\left(T_i - T_j\right) \tag{22}$$

Noting that:

$$p_i - p_j = \left(\sqrt{E^2 - m_i^2} - \sqrt{E^2 - m_j^2}\right) = \frac{\Delta m^2_{ij}}{2E} \tag{23}$$

We can substitute for Δ explicitly to find the oscillation phase in this model:

$$\frac{\Delta_{ij}}{2}L = \left(\frac{\Delta m^2_{ij}}{4E} + \frac{1}{2}nRe[T_i - T_j]\right)L \tag{24}$$

We consider two cases. First, when E is very far from the resonant energy $m_h^2/2m_i$, both the new terms will be negligible relative to the standard oscillation phase, and we retrieve the usual neutrino oscillation formula. Thus, at energies far from resonance, ordinary neutrino oscillations are observed.

$$P_{\alpha\beta}|_{|E-E_{res}|\gg 0} \sim P^0_{\alpha\beta} \tag{25}$$

$$P^0_{\alpha\beta} = \delta_{\alpha\beta} - 4\sum_{i>j} Re(U^*_{\alpha i}U_{\beta i}U_{\alpha j}U^*_{\beta j})sin^2\left[\frac{\Delta m^2_{ij}}{4E}L\right] + 2\sum_{i>j} Im(U^*_{\alpha i}U_{\beta i}U_{\alpha j}U^*_{\beta j})sin^2\left[\frac{\Delta m^2_{ij}}{2E}L\right] \tag{26}$$

Second, at short baselines, the standard oscillation phase will be negligible. At energies when near resonance for one mass state k, $E \sim E_{res,k}$, and one new term will be large and the other small. In what follows, we will, for simplicity, neglect CP-violation. Then, near resonance and at short baseline there is a new oscillation effect:

$$P_{\alpha\beta}|_{L\ll L_{osc}} \sim -4\sum_{i>j}(U_{\alpha i}U_{\beta i}U_{\alpha j}U_{\beta j})sin^2\left[\frac{1}{2}nRe[T_i\delta_{ik} - T_j\delta_{jk}]L\right] \tag{27}$$

$$sin^2 2\theta\, sin^2\left[\frac{1}{2}nRe[T_k]L\right] \tag{28}$$

The effective angle $sin^2 2\theta$ is determined entirely from the PMNS matrix, assumed to be well measured by off-resonance experiments:

$$sin^2 2\theta_{\alpha\beta} = \left(4U_{\alpha k}U_{\beta k}\sum_{j\neq k}U_{\alpha j}U_{\beta j}\right) \tag{29}$$

$$= 4U_{\alpha k}U_{\beta k}(1 - U_{\alpha k}U_{\beta k}) \tag{30}$$

For example, for $\nu_\mu \to \nu_e$ oscillations near resonance of mass state ν_3, which is the heaviest under the normal ordering, the effective mixing angle is:

$$sin^2 2\theta^{NO}_{\mu e} = 4s_{13}s_{23}c_{13}\left(1 - s_{13}s_{23}c_{13}\right) = 0.33 \tag{31}$$

On the other hand, given an inverted mass ordering, the heaviest mass state (and thus the one with the largest coupling to the new Higgs field) is mass state ν_2, and then:

$$sin^2 2\theta^{IO}_{\mu e} = 4s_{12}c_{13}(c_{12}c_{23} - s_{12}s_{23}s_{13})\left[1 - s_{12}c_{13}(c_{12}c_{23} - s_{12}s_{23}s_{13}\right] = 0.89 \tag{32}$$

Allowing for non-zero CP violation introduces further dependencies on the CP-phases of the PMNS matrix which have not been measured, thus introducing extra fit parameters and reducing the predictiveness of this model. These considerations would modify our results quantitatively by O(1) numbers, though they do not change our primary conclusions.

Based on the above discussions we have constructed a model which introduces one new particle into the SM to explain neutrino mass at a natural energy scale. Assuming only one resonance is accessible in a given experiment, this model can be parametrized in terms of three phenomenological parameters, X, E_{res} and g_i, via:

$$P_{\alpha\beta} = sin^2 2\theta sin^2\left[\frac{Y}{2}\left(\frac{E - E_{res}}{(E - E_{res})^2 + \left(\frac{1}{2}E_{res}g_i^2\right)^2} + \frac{-E - E_{res}}{(E + E_{res})^2 + \left(\frac{1}{2}E_{res}g_i^2\right)^2}\right)\right] \tag{33}$$

$$Y = \frac{g_i^2 n}{8m_i} \qquad E_{res} = \frac{m_h^2}{2m_i} \tag{34}$$

It is also possible to consider a reduced expression that applies at low values of g_i, in practice less than around $g \sim 0.1$, given by the ZWA. Under such an approximation:

$$P^{ZWA}_{\alpha\beta} = sin^2 2\theta sin^2\left[Y\frac{E_{res}}{E^2 - E^2_{res}}\right] \tag{35}$$

Some examples of the expected oscillation behavior are shown in Fig. 2 for the ZWA. A narrow peak which oscillates rapidly is visible above a baseline of no-oscillations, off resonance. Fig. [] shows oscillation probabilities for some finite-width models. For narrow widths, the character of the resonance remains clear. However, for much larger values of the coupling, the resonance becomes sufficiently wide that nontrivial oscillation structure becomes visible as the oscillation phase changes across the resonance.

4 Signatures in SBN Experiments

Here we explore whether this new signature can explain the oscillation anomaly observed in the MiniBooNE short baseline neutrino experiment [1]. We perform fits to the oscillation hypothesis described above in both ZWA and finite width models, and compare with fits made to the more standard 3+1 sterile neutrino model.

Our fit uses public MiniBooNE data release accompanying [1], and employs the χ^2 minimization recipe described along with that data release. We use the full MiniBooNE covariance matrix to account for systematic error in both neutrino and antineutrino modes. Fits to the 3+1 model are shown in Figure 4, and are in qualitative agreement with the official MiniBooNE results.

For the ZWA we calculate the χ^2 value for points on a grid in (Y,E_{res}) and then use Wilks theorem to draw 68%, 90% and 99% confidence intervals in this 2D parameter space. There are two allowed regions. A sharp resonance around 360 MeV well describes the data in both neutrino and antineutrino modes, with χ^2 values that indicate large p-values for the fit. A very low-mass resonance with large ample also provides a good fit. These two cases are shown in Fig. 6.

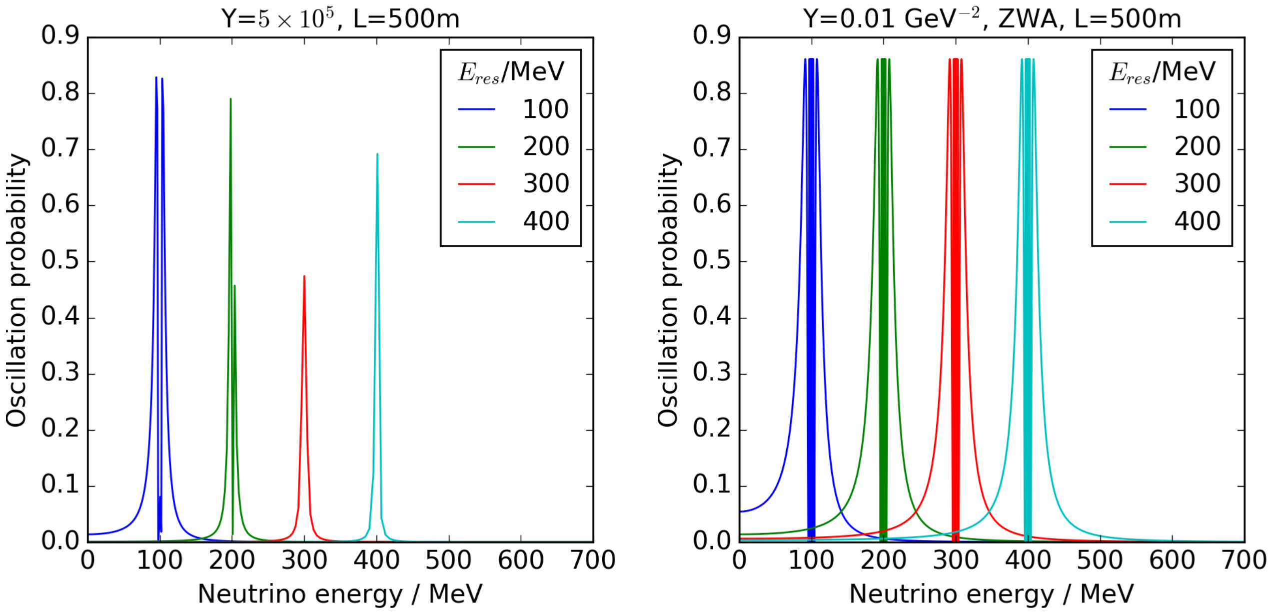

Figure 2: Some example oscillation probabilities under the Zero Width Approximation.

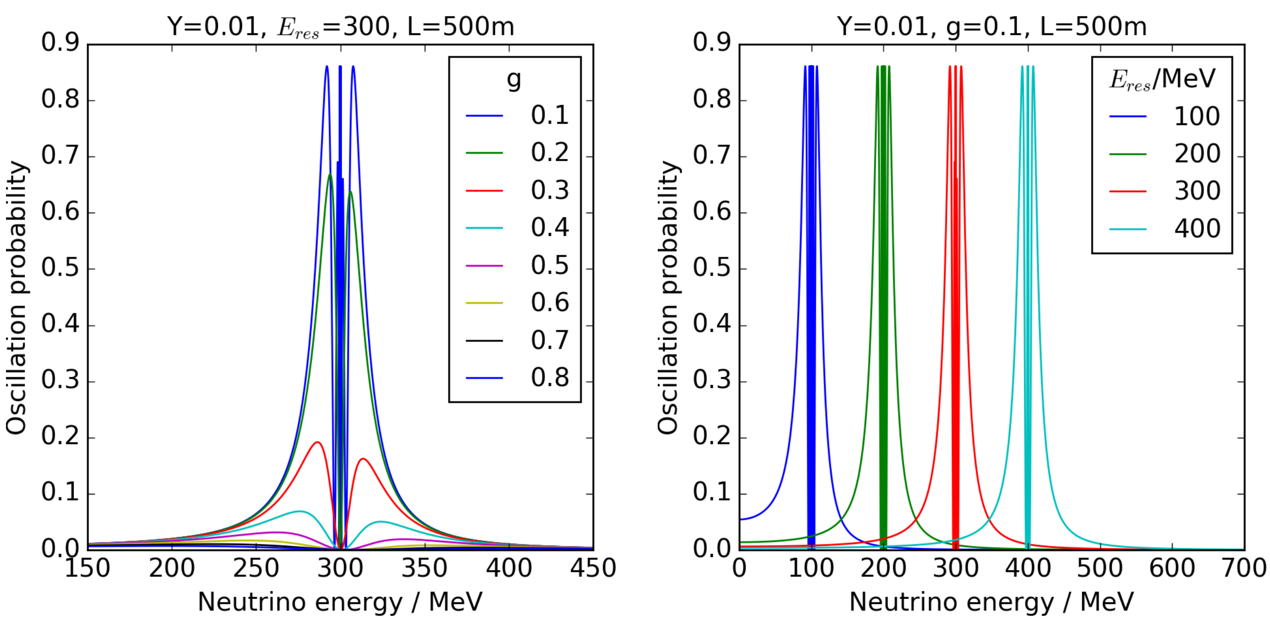

Figure 3: Some example oscillation probabilities given finite new Higgs widths. On the right we see that the oscillation behavior with g=0.1 is very similar ot the ZWA.

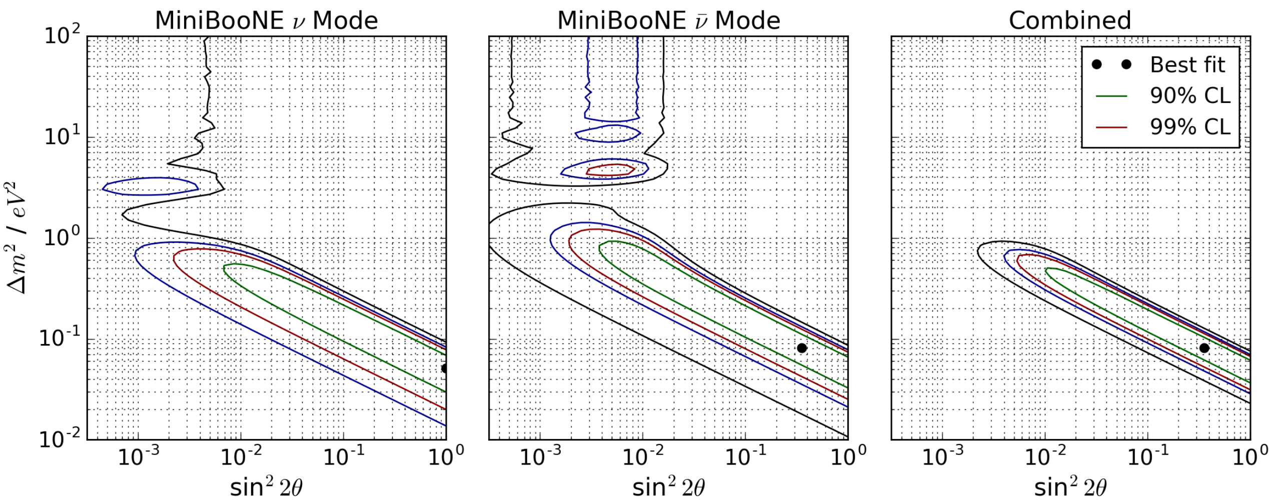

Figure 4: MiniBooNE fit to the 3+1 sterile neutrino oscillation hypothesis made using the same χ^2 minimization approach as used in our new Higgs model fits.

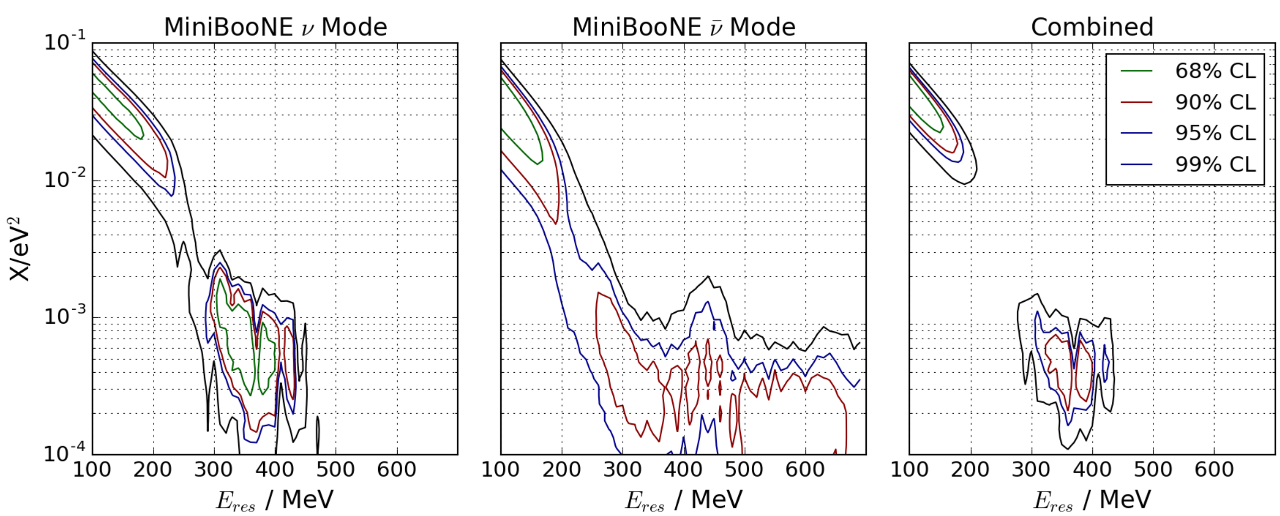

Figure 5: Allowed regions in the ZWA. Left shows fits to neutrino mode only; Center antineutrino mode only; Right, combined fit.

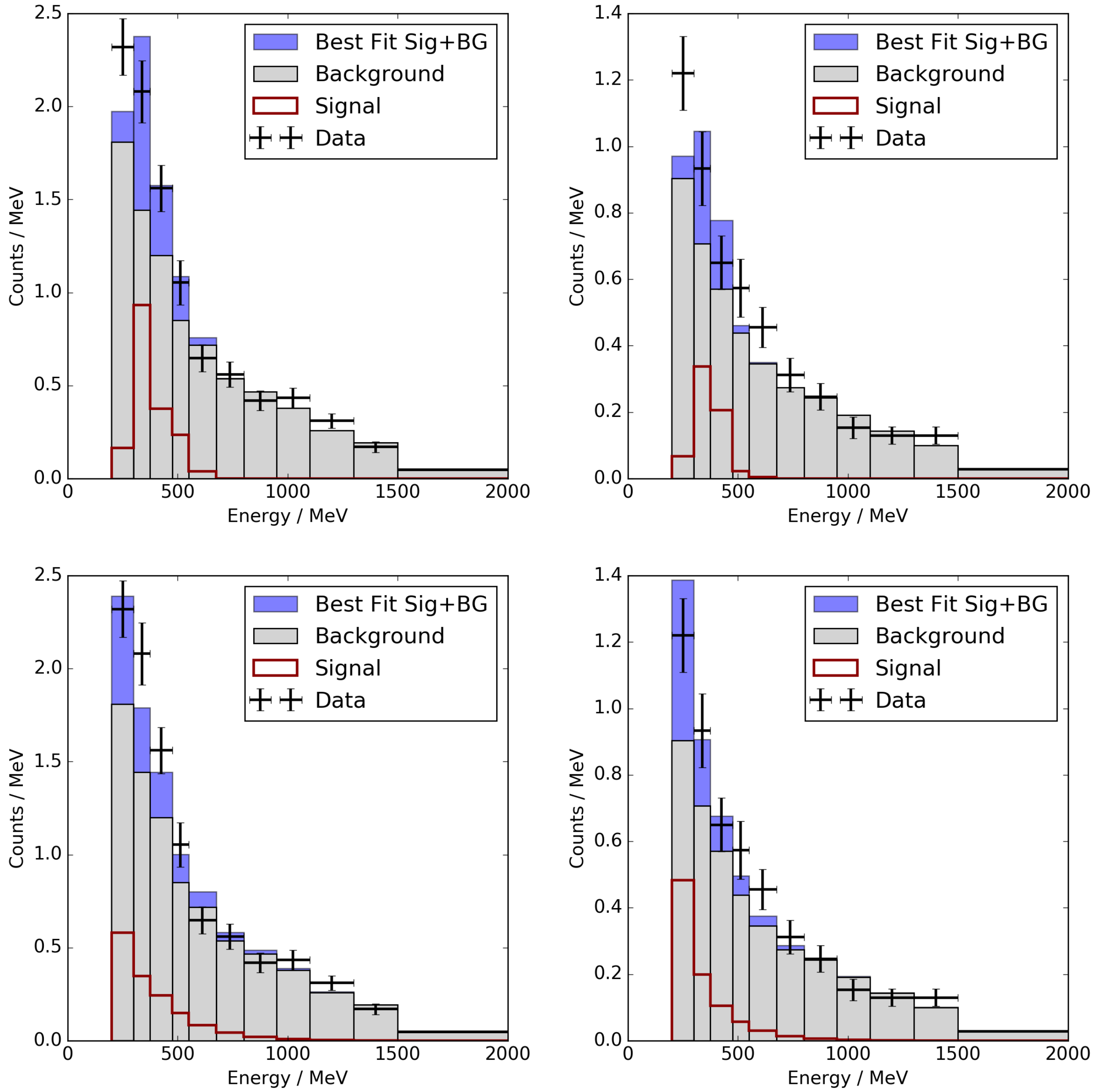

Figure 6: Typical points in the allowed region for the ZWA. Top: low-mass best fit; bottom: High mass region (E_{res} =360 MeV, Y=5×10^{-4}); low mass region (E_{res} = 100 MeV, Y=4×10^{-2}). Left: neutrino mode; Right: antineutrino mode.

Hypothesis	**DOF**	χ^2_ν	$\chi^2_{\bar{\nu}}$	$\chi^2_\nu + \chi^2_{\bar{\nu}}$	$\Delta\chi^2_{null-bf}$ (dof)
No Oscillations	22	18.9	15.26	34.16	N/A
3+1 sterile neutrino, MB Best fit	20	11.41	4.68	16.11	18.1 (2)
3+1 sterile neutrino, global best fit	20	23.0	12.45	35.44	-1.28 (~)
h_ν, ZWA, low mass (abs)	20	6.79	7.38	14.17	20.0 (2)
h_ν, ZWA, high mass	20	4.75	12.43	17.17	17.0 (2)
h_ν, finite width	19	6.06	6.08	12.14	22.0 (3)

Table 1: χ^2 values and p-values for null hypothesis, the best fit in the 3+1 model, and the best fit in the neutrino-Higgs model. In the latter case we consider both the absolute best fit (low mass), the best fit point in the high mass island, and the best fit in the finite width case.

Relaxing the ZWA and extending the fit into the larger, 3-parameter space allows for more flexibility to match the hypothesis at the cost of an extra fitted degree of freedom. Slices at a few values of g are shown in Figure 7. Again, there are two allowed regions, with larger g tending to favor higher Y as the peak of the resonance becomes suppressed by its finite width. The best fit χ^2/DOF in the 3-parameter fit is 14.17 / 19DOF, to be compared with 12.14 / 20DOF under the ZWA. Hence it does not appear that the additional parameter significantly improves the quality of the model beyond simply adding an additional degree of freedom. We thus consider the best fit from the ZWA, which is a 2 parameter model which can well match the data, to be our primary result.

In both ZWA and finite-width models, this hypothesis gives a drastic improvement over the 3+1 model and the null hypothesis, for fitting the MiniBooNE excess. Table 1 gives a comparison at the best fit point (fitted using both neutrino and antineutrino datasets) between null, 3+1, new Higgs with ZWA, and new Higgs with finite width.

5 Discussion

We have shown that excellent fits to the MiniBooNE anomaly can be achieved using a model with a new, light boson that couples to neutrinos. Resonances in the neutrino oscillation probability are achieved where the forward scatter of neutrinos from the CNB through the S channel has a resonance. The allowed regions in our fits suggest that resonances around $E_{res} \sim 100$ MeV and $E_{res} \sim 380$ MeV provide the optimal fit point. This resonance energy is related to the mass scale of the new boson by:

$$E_{res} = \frac{m_h^2}{2m_i} \tag{36}$$

If the MiniBooNE data is to be explained by this coupling, we expect that it must be the heaviest neutrino mass state which is on resonance, in order to avoid detection in other, higher energy experiments. The mass of this state is constrained from oscillation experiments and cosmology, and leads us to conclude that the mass of this new boson would be either 6-9 keV in the case of the higher energy island, or 3-5 keV for the lower energy one.

Fitted values of Y of 2×10^{-4} and larger are allowed at 90% CL, which imply a required local CNB over-density, via the relation:

$$n = \frac{8m_iY}{g_i^2} \tag{37}$$

Taking the heaviest neutrino mass at its smallest allowed value of 0.05 eV, a CNB density $\geq 8 \times 10^{-5}$ eV^3 is implied. This exceeds the standard expected global relic density by a factor of 1.8×10^8. This is large, but not excluded by direct measurements, and smaller than over-densities that have been discussed in theoretical literature, and are being probed experimentally, for example at KATRIN .

It is important to ask, if this resonance were really the explanation for the MiniBooNE anomalies, would it have been seen by other experiments? Fig. 8 shows the energy vs baseline of various neutrino oscillation experiments, with the location of these possible resonances marked. For the 360 MeV resonance, the possible experiments that may see effects include T2K and SuperK. At 100 MeV, the only relevant experiment is LSND, which also finds an excess. Thus we expect the low mass fit point to be relatively unconstrained by world data. On the other hand, the higher mass best fit point may be in tension with other experiments, and a global fit with all world data is required to fully understand its validity.

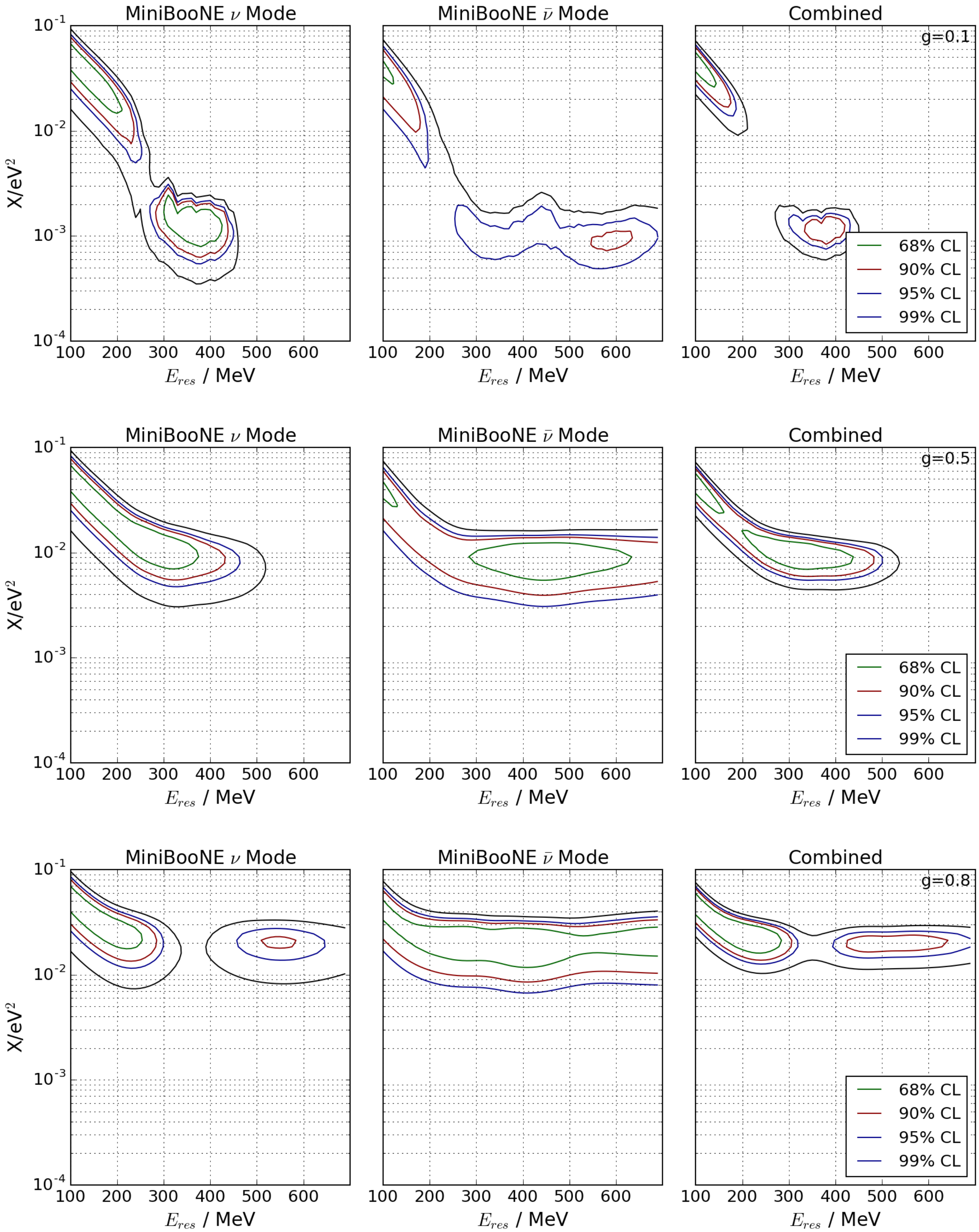

Figure 7: Slices in the allowed region at fixed g given the finite width approach.

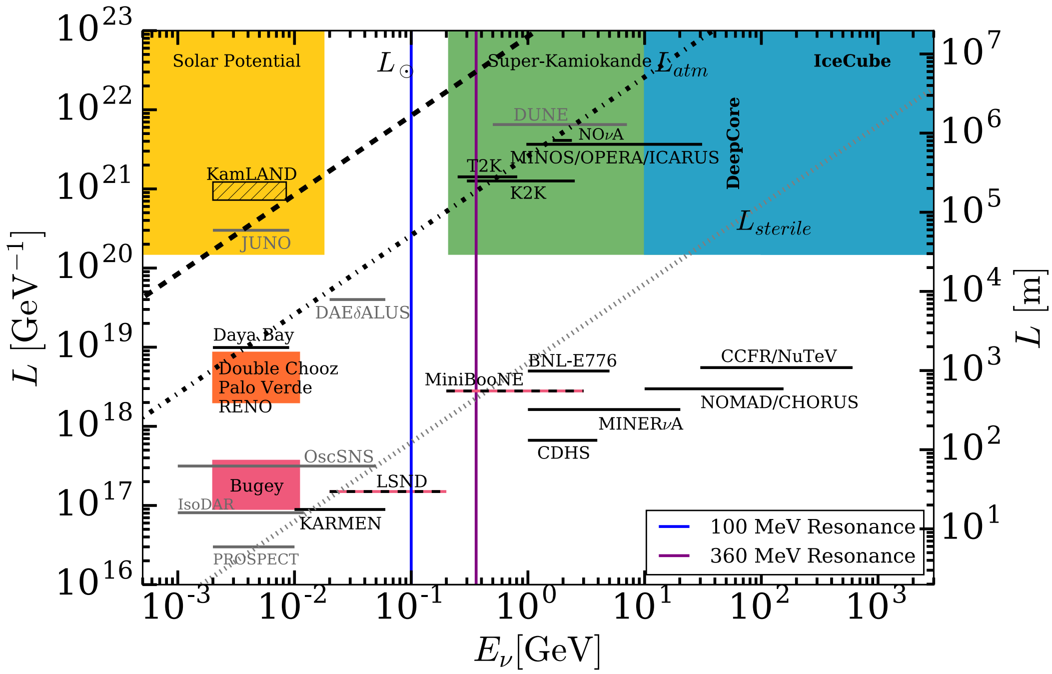

Figure 8: Energies and baselines of present and proposed neutrino oscillation experiments.

6 Conclusions

We have presented a model whereby neutrinos couple to a new light boson with a strength proportional to their mass. The mass scale of this new boson is in the keV range. This new boson, if it couples to neutrinos, may introduce a new resonant oscillation effect due to interactions with the relic neutrino background. With a sufficiently over-dense relic background, these oscillations provide a good fit to the MiniBooNE anomalies. Although the large over-density required appears unlikely, it is still within experimental limits and smaller than other that have been considered in the literature.

References

[1] A. A. Aguilar-Arevalo et al. Improved Search for $\bar{\nu}_\mu \to \bar{\nu}_e$ Oscillations in the MiniBooNE Experiment. *Phys. Rev. Lett.*, 110:161801, 2013.

[2] Seungwon Baek and Takaaki Nomura. Dark matter physics in neutrino specific two Higgs doublet model. 2016.

[3] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, Marc Sher, and Joao P. Silva. Theory and phenomenology of two-Higgs-doublet models. *Phys. Rept.*, 516:1–102, 2012.

[4] J. S. Diaz and F. R. Klinkhamer. Neutrino refraction by the cosmic neutrino background. *Phys. Rev.*, D93(5):053004, 2016.

[5] John F. Donoghue. The Multiverse and Particle Physics. 2016.

[6] Amand Faessler, Rastislav Hodak, Sergey Kovalenko, and Fedor Simkovic. Can one measure the Cosmic Neutrino Background? *Int. J. Mod. Phys.*, E26(01n02):1740008, 2017.

[7] S. Gabriel and S. Nandi. A New two Higgs doublet model. *Phys. Lett.*, B655:141–147, 2007.

[8] C. Giunti and C. W Kim. *Fundamentals of neutrino physics and astrophysics.* Oxford University Press, 2007.

[9] A. Kaboth, J. A. Formaggio, and B. Monreal. Sensitivity of Neutrino Mass Experiments to the Cosmic Neutrino Background. *Phys. Rev.*, D82:062001, 2010.

[10] R. Lazauskas, P. Vogel, and C. Volpe. Charged current cross section for massive cosmological neutrinos impinging on radioactive nuclei. *J. Phys.*, G35:025001, 2008.

[11] Andreas Ringwald and Yvonne Y. Y. Wong. Gravitational clustering of relic neutrinos and implications for their detection. *JCAP*, 0412:005, 2004.

[12] Fei Wang, Wenyu Wang, and Jin Min Yang. Split two-Higgs-doublet model and neutrino condensation. *Europhys. Lett.*, 76:388–394, 2006.