Lepton Number Violating Electron Recoils at XENON1T and PANDAX by the $U(1)_{B-L}$ Model with Non-Standard Interactions

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We propose an $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ model, in which the neutrino masses and mixings can be generated via Type-I seesaw mechanism after $U(1)_{B-L}$ breaking. A light mediator emerges and enables non-standard interaction that violates the lepton number. We show that the non-standard interaction leads to low energy recoil events that is consistent with the observed KeV range electron recoil events at direct detection experiments. The XENON1T electron recoil data are best explained with the light scalar mediator’s combined coupling strength to electron and neutrino $\sqrt{y_e y_\nu} = 0.9 \times 10^{-6}$. The PANDAX result is consistent with $\sqrt{y_e y_\nu} < 1.4 \times 10^{-6}$.

I. INTRODUCTION

Recently, the XENON Collaboration reported the low-energy electronic recoil data with an exposure of 0.65 ton-years$^1$. With 285 observed events over an expected background of 232±15 events, an excess is observed at electron recoil energy below 7 keV, rising towards lower energies and is most prominent between 2 and 3 KeV. Possible sources of the excess include the $\beta$-decay from a trace amount of tritium impurity ($6.2 \pm 2.0 \times 10^{-20}$ mol/mol) in the detector, which explained at 3.2$\sigma$ $^1$. Also, XENON1T found that the solar axion and the solar neutrino with magnetic moment respectively provide 3.5$\sigma$ and 3.2$\sigma$ significance fit to the excess in certain parameter ranges, which are unfortunately in tension with stellar cooling bounds$^{2,4}$. When an unconstrained tritium component is included in the fitting, both the solar axion and the solar neutrino magnetic moment hypotheses lose the substantial statistical significance, and their significance levels are decreased to 2.1$\sigma$ and 0.9$\sigma$, respectively. The bosonic dark matter has been studied as well, but the global significance is less than 3$\sigma$. Moreover, the other generic concern is the recoil energy very close to the detection threshold, which might affect the performance of the detector in these energy bins. Soon after, the PANDAX experiment reported the low-energy electron recoil results$^5$, where the rise in recoil events peaks at a slight higher 3-7 KeV range. PANDAX performed systematic analyses on potential experimental backgrounds, including trace amount of tritium, and xenon, krypton isotopes$^5$, and derived constraints on solar axion and anomalous neutrino dipole moments.

Following the XENON1T result, this excess has been extensively discussed in a number of emerging papers on the non-standard neutrino-electron interactions with light mediators$^6, 14$, and with solar axions$^{15, 16}$, axion-like dark matter$^{17}$, hidden dark photon dark matter$^{18, 20}$, warm or fast moving dark matter$^{21}$ (also see$^{22}$), boosted$^{23, 27}$, inelastic or multi-component dark matter$^{28, 33}$, decaying dark matter$^{34, 36}$, Migdal effect$^{37}$, luminous or shining dark matter$^{38, 39}$, inverse Primakoff effect$^{40}$, hydrogen decay$^{41}$, dark fluxes from accreting black holes$^{42}$, as well as re-examining detector backgrounds$^{43}$, collider searches$^{44}$, neutrino magnetic moment$^{45}$, and stellar cooling$^{46, 47}$ limits.

In this paper, we propose a $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ model to derive low-energy electron recoils via non-standard neutrino-electron interactions with a light mediator. After the $U(1)_{B-L}$ breaking, the neutrino masses and mixings can be generated via the Type I seesaw mechanism. In particular, a light mediator exists in the model, as well as the non-standard interactions between the light mediator and leptons which violate the lepton number. We show that these recoil events fits in well with the observed KeV results. When various residue backgrounds are included, an upper limit on the non-standard neutrino interactions can be obtained by fitting to XENON1T and PANDAX electron recoil data.

II. MODEL SETUP

The particles in the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ model follow the conventional notations that the Standard Model (SM) quark doublets $Q_i$, right-handed up-type quarks $U_i$, right-handed down-type quarks $D_i$, lepton doublets $L_i$, right-handed charged leptons $E_i$, and right-handed neutrinos $N_i^c$, with $i = 1, 2, 3$ for three generations. Then we introduce new scalar fields, including one $SU(2)_L$ triplet $\Phi$, two $SU(2)_L$ doublets $H$ and $H'$, and two SM singlets $S$ and $T$. These particles and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ gauge symmetry are summarized in Table I.

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The particles in the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ model follow the conventional notations that the Standard Model (SM) quark doublets $Q_i$, right-handed up-type quarks $U_i$, right-handed down-type quarks $D_i$, lepton doublets $L_i$, right-handed charged leptons $E_i$, and right-handed neutrinos $N_i^c$, with $i = 1, 2, 3$ for three generations. Then we introduce new scalar fields, including one $SU(2)_L$ triplet $\Phi$, two $SU(2)_L$ doublets $H$ and $H'$, and two SM singlets $S$ and $T$. These particles and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ gauge symmetry are summarized in Table I.

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As we know, the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ gauge symmetry can be obtained from the $SO(10)$ gauge symmetry breaking. However, heavy $X_E$, $X_E^c$ fields, and the scalars $\Phi$, $H'$, $S$, $T$ are not new particles in the traditional $SO(10)$ models. Interestingly, they can be obtained via the tensor products of the 10 fundamental representation, as well as $\mathbf{16}$ and $\overline{\mathbf{16}}$ spinor representations of $SO(10)$, i.e., the higher representations of $SO(10)$ as follows

$X_E \subset \mathbf{16} \otimes \mathbf{16} \otimes \mathbf{16}$,

$X_E^c \subset \mathbf{16} \otimes \mathbf{16} \otimes \mathbf{16}$,

$\Phi \subset \mathbf{126} \subset \mathbf{16} \otimes \mathbf{16}$,

$H' \subset \mathbf{10} \otimes \mathbf{16} \otimes \mathbf{16}$,

$S/T \subset \mathbf{120} \subset \mathbf{16} \otimes \mathbf{16}$.

We consider the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ groups as an intermediate stage of symmetry breaking sequence after the breaking of $SO(10)$.

The $U(1)_{B-L}$ gauge symmetry spontaneously breaks after $T$ obtains a vacuum expectation value (vev), leaving out the SM $SU(3)_C \times SU(2)_L \times U(1)_Y$ at lower energy scales. $H$ is the SM Higgs doublet whose vev finally breaks the electroweak gauge symmetry, and we assume that $H$, $H'$ and $S$ do not acquire vevs. In particular, the effective Yukawa couplings between $H'$ and charged leptons can be generated if we introduce a pair of vector-like particles ($X_E, X_E^c$) as heavy mediators and $\Phi$ can couple to lepton doublets as well. The $X_E$ mass are generated by some UV symmetry breaking above the $U(1)_{B-L}$ scale. Because the CP-even neutral components of $\Phi$, $H'$, and $S$ can mix with each other, their lightest mass eigenstate $s$ can couple to the charged leptons as well as neutrinos. In short, $T$ and $H$ are introduced to break the $U(1)_{B-L}$ gauge symmetry and electroweak gauge symmetry, respectively, while $\Phi$, $H'$, and $S$ are introduced to generate the Yukawa couplings between the lightest CP-even neutral scalar and charged leptons/neutrinos.

$$
\begin{array}{c|c|c|c}
& SU(3)_C & SU(2)_L & U(1)_Y & U(1)_{B-L} \\
\hline
Q_i & 3 & 2 & 1/6 & 1/6 \\
U_i & 3 & 1 & 2/3 & 1/6 \\
D_i & 3 & 1 & -1/3 & 1/6 \\
L_i & 1 & 2 & -1/2 & -1/2 \\
E_i & 1 & 1 & -1 & -1/2 \\
N_i & 1 & 1 & 0 & -1/2 \\
X_E & 1 & 1 & 1 & 3/2 \\
X_E^c & 1 & 1 & 1 & 3/2 \\
\Phi & 1 & 3 & 1 & 1 \\
H & 1 & 2 & -1/2 & 0 \\
H' & 1 & 2 & -1/2 & -1 \\
S & 1 & 1 & 0 & -1 \\
T & 1 & 1 & 0 & -1 \\
\end{array}
$$

TABLE I. Particles and their quantum numbers under the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ gauge group. $T$ and $H$ sequentially develop vev that break $U(1)_{B-L}$ and $SU(2)_L \times U(1)_Y$, respectively.

The Lagrangian sector involving the fermions is

$$
\mathcal{L} = y_{ij}^U Q_i U_j^c \bar{H} + y_{ij}^D Q_i D_j^c H + y_{ij}^L L_i E_j^c H + y_{ij}^L L_i N_j \bar{H} + y_{ij}^N T N_i^c N_j^c + y_{ij}^T L_i \Phi L_j + y_{ij}^{T} H^i L_i X E^c + y_{ij}^{T} T E_i^c X E + M_{XE} XE^c XE + \text{H.C.}.
$$

(1)

With the $y_{ij}^T L_i N_j \bar{T}$ and $y_{ij}^{N} T N_i^c N_j^c$ terms, we can generate the neutrino masses and mixings via Type I seesaw mechanism after $T$ acquires a vev and breaks the $U(1)_{B-L}$ gauge symmetry. Now we can see non-standard neutrino interactions emerge from this model setup. The vector-like ($X_E, X_E^c$) masses are assumed heavy as the result of UV symmetry breaking. Integrating out ($X_E, X_E^c$), we can obtain effective operators,

$$
\mathcal{L} \supset -\frac{1}{M_{XE}} y_{ij}^{H'} y_j^{H'} T L_i E_j^c + \text{H.C.}.
$$

(2)

After $U(1)_{B-L}$ gauge symmetry breaking, we get

$$
\mathcal{L} \supset -\frac{\langle T \rangle}{M_{XE}} y_{ij}^{H'} y_j^{H'} T L_i E_j^c + \text{H.C.}.
$$

(3)

For simplicity with phenomenology, we can assume that only $y_{ij}^{H'}$ and $y_{ij}^{T}$ are non-zero. From the terms $y_{ij}^{\Phi} L_i \Phi L_j$ and Eq. (3), we obtain

$$
\mathcal{L} \supset \frac{y_{ij}^{\Phi}}{2} \sin \alpha \cos \beta \bar{s}_i \nu_{ij} e + y_{s} \bar{e} \nu_{e} e + \text{H.C.}
$$

(4)

where $\nu_{ij}$ is the neutrino mass eigenstate, $e$ is the electron, $s_1$ is a CP-even mass eigenstate from $S, \Phi, H'$ mixing, and

$$
y_{e} = \frac{\langle T \rangle}{M_{XE}} y_{ij}^{H'} y_j^{H'} T s_i \sin \alpha \sin \beta.
$$

(5)

For a heavy $M_{XE}$ mass much above $U(1)_{B-L}$ breaking scale the effective coupling $y_{es}$ can be naturally small. $\alpha, \beta$ denote the mixing angles from the neutral components in $S, \Phi, H'$,

$$
\begin{pmatrix}
 s_1 \\
 s_2 \\
 s_3
\end{pmatrix} = \begin{pmatrix}
 \cos \alpha & \sin \alpha \cos \beta & \sin \alpha \sin \beta \\
 -\sin \alpha & \cos \alpha \cos \beta & \cos \alpha \sin \beta \\
 0 & -\sin \beta & \cos \beta
\end{pmatrix} \text{Re} \left( \frac{S}{\Phi^0} \right)
$$

(6)

$s_1, s_2, s_3$ are their mass eigenstate and we can take $s_1$ to be the lightest state, $s_1 = \cos \alpha \text{Re} S + \sin \alpha \cos \beta \text{Re} \Phi^0 + \sin \alpha \sin \beta \text{Re} H^0$. 

![FIG. 1. Feynman diagram for the NSI $\nu \rightarrow \nu e^{-} e^{-}$ scattering.](image-url)
from diagonalizing their mass matrix. Due to the number of scalar fields, this model has an extended scalar potential, and its general form can be written as

\[ V = m_3^2 |S|^2 - m_3^2 |T|^2 - m_3^2 |H|^2 + m_3^2 |H'|^2 + m_3^2 |\Phi|^2 + \lambda_S |S|^4 + \lambda_T |T|^4 + \lambda_H |H|^4 + \lambda_H' |H'|^4 + \lambda_\Phi |\Phi|^4 + \lambda_{ST} |S|^2 |T|^2 + \lambda_{SH} |S|^2 |H|^2 + \lambda_{SH'} |S|^2 |H'|^2 + \lambda_{S\Phi} |S|^2 |\Phi|^2 + \lambda_{T\Phi} |T|^2 |\Phi|^2 + \lambda_{H\Phi} |H|^2 |\Phi|^2 + \lambda_{H\Phi'} |H'|^2 |\Phi|^2 + (A_1 \Phi \Phi' H' + A_2 \Phi \Phi' H + \Phi \Phi' \Phi' + H' + \Phi + \Phi' + H + H') , \]

where \( \Pi = i \sigma_2 H^* \) and \( \Pi' = i \sigma_2 H'^* \) with \( \sigma_2 \) the second Pauli matrix, and we neglect the \( \Phi \Pi' \) and \( \Phi \Pi H \) terms, which will induce the tadpole terms for \( H, T \) and \( \Phi \).

Here we need a light \( s_1 \) to mediate non-standard soft neutrino scattering that is relevant to electron recoils. In principle, because \( S \) and \( T \) carry the same quantum numbers, without loss of generality, we can make a \( U(1)_{B-L} \) rotation so that only one linear combination of them has a vev in case they both have vevs. After \( H \) and \( T \) obtain vevs, the neutral components of \( \Phi, S \), and \( H' \) fields will mix with each other, and we assume the lightest CP-even mass eigenstate to be very light in this paper. This typically would assume \( V \) to be flat in some direction of Re\{\( S, \Phi^0, H^0 \)\}. Before analysing the electron recoil phenomenology, we would like to emphasize the difference from the \( SU(3)_C \times SU(2)_L \times U(1)_{B-L} \) model, i.e., the traditional \( U(1)_{B-L} \) model \( \| \)\( \| \). The main point is that \( A_1 \Phi H H', A_2 \Phi \Pi H, \) and \( \lambda_\Phi \Pi H \) terms are necessary to generate the mixings among the neutral components of \( \Phi, S, \) and \( H' \), as well as \( g_i L_i \Phi L'_{j} \) to generate the \( g_i \Phi \Sigma \nu_j \nu_j \) terms. In the traditional \( SU(3)_C \times SU(2)_L \times U(1)_{B-L} \) model, the lepton doublets \( L_i \) are charged under \( U(1)_{B-L} \) while the Higgs field \( H \) is neutral under \( U(1)_{B-L} \). Thus, our model cannot be realized in the traditional \( SU(3)_C \times SU(2)_L \times U(1)_{B-L} \) model.

III. ELECTRON RECOIL

The \( s_1 \overline{\nu} \nu \) and \( s_1 \overline{\nu} e \) vertices in Eqs.3 and 4 allows \( s_1 \) to mediate 'non-standard' electron scattering process \( \nu \nu \rightarrow \nu \nu e \) as shown in Fig.1. Since \( s_1 \) carries lepton number, this process violates the lepton number by two units, and has no corresponding diagrams in the SM. The scattering amplitude-square is

\[ |M|^2 = \frac{y_\nu^3 y_\nu'^3 (4 M^2_{\nu} - t)^2}{(M^2_{\nu} - t)^2} , \]

where \( t = (p_\nu - p_\nu')^2 (p_\nu - p_\nu) \) is the Mandelstam \( t \) variable. Here we will denote \( y_\nu' = \frac{y_\nu^3}{2} \sin \alpha \cos \beta \) and neglect the flavor indices for convenience. For negligible neutrino masses, \( t = -2 M_{\nu} E_\nu \) for neutrinos scattering off a free electron, and the differential cross-section is

\[ \frac{d \sigma^\nu_e}{dE_\nu} = \frac{y_\nu^2 y_\nu'^2 M_{\nu} E_\nu (2 M_{\nu} + 2 M_e)}{8 \pi E_{\nu}^2 (M^2_{\nu} + 2 M_{\nu} E_\nu)^2} . \]

\( E_\nu \) is the electron's acquired kinetic energy after scattering, and \( E_\nu \) is the incident neutrino energy, see [3] for detail. Noted the \( s_1 \overline{\nu} \nu \) vertex leads to different scattering kinematics comparing to that from a lepton number conserving \( s_\nu \nu \) vertex [52] at large momentum exchange, where the NSI scattering spectrum with the \( s_1 \overline{\nu} \nu \) vertex is harder at very large recoil energy \( E_\nu \sim E_\nu \) if \( M_{\nu} \) were assumed heavy. However, since we are interested in KeV recoils in this paper, the soft recoil spectra from these two types of vertices would converge and demonstrate a similar \( E^{-1} \) behavior as long as \( M_{\nu} \) is assumed smaller than the transferred momentum. In particular, it features a kinematic region:

\[ M_{\nu} \ll \sqrt{2 M_{\nu} E_\nu} \ll M_{\nu} . \]

Here low momentum transfer dominates the scattering as the cross-section in Eq. 3 behaves as \( d\sigma/dE_\nu \propto E_\nu^{-1} \). For \( s_\nu \nu \) mass below the KeV scale, Eq. 10 is typically satisfied by near-threshold (KeV) energy transfer in electron recoil events with solar/reactor neutrinos at direct-detection experiments. Eq. 9 is for a free electron, and the differential rate for recoil energy \( E_{\nu} \) would be

\[ \frac{dN}{dE_\nu} = N \cdot T \cdot \epsilon(E_\nu) \int dE_{\nu} \mathcal{G}(E', E_{\nu}) \int dE_{\nu} \mathcal{F}(E') \frac{d \phi_{\nu}}{dE'} \frac{d \sigma^\nu_{ee}}{dE' \ dE'}, \]

where \( N \) and \( T \) are the number of targets and exposure time. For comparison with XENON1T results, \( \epsilon \) is the detector efficiency [11]. \( \mathcal{G} \) is a Gaussian smearing on \( E_\nu \) that accounts for detector resolution,

\[ \mathcal{G} = \frac{1}{\sqrt{\sigma^2}} \exp \left( -\frac{(E_\nu - E_{\nu})^2}{\sigma^2} \right) \]

\( \delta_E = \sqrt{0.3 E} + 0.0037 E, \) where \( E \) is in KeV, as given in Ref. [11]. \( \phi_{\nu} \) is the Solar neutrino flux model that we adopt from Ref. [53]. \( \mathcal{F}(E) = \sum_i \theta(E - B_i) \) is a sum of step-functions with threshold at the th atom's binding energy, which has been shown to be a good approximation of the corrections from atomic binding [51, 52].

The \( E_{\nu}^{-1} \) dependence from light scalar mediated NSI leads to a relative moderate spectrum rise at the lowest energy bins, if compared to a more steep \( E_{\nu}^{-2} \) dependence observed in light vector-boson mediated scenarios, as studied in Ref. [39, 40, etc. Note that photon-mediated BSM neutrino dipole interaction would also give an \( E_{\nu}^{-1} \) dependence thus should yield similar explanation to data. In following sections we show the \( s_1 \overline{\nu} \nu \) NSI gives good fit the low energy electron recoils, and its significance is subject to effects from several SM radiative backgrounds.
IV. FIT TO XENON1T

The solar neutrino’s $\nu e \rightarrow \nu' e$ event rate rises towards low energy, which helps explaining the near-threshold excess reported by XENON1T. We make a likelihood fit to the 29 binned data [1] below 30 KeV, by combining these NSI-induced events with XENON1T’s best-fit background modeling $B_0$,

$$\chi^2 = \sum_i (\eta B_0 + N_{i^e} - N_{i^e}^{\text{data}})^2 / (\delta N_i)^2 + (1 - \eta)^2 / (\delta \eta)^2,$$  

(13)

where last term accounts for a small but crucial normalization uncertainty in the background model. In the low $E_R$ range, the detector background $B_0$ is primarily the flat $^{214}$Pb component, which is a calibrated in the entire 1-210 KeV range and has a 2\% statistic uncertainty. Detector efficiency modeling would contribute another 1\% normalization uncertainty, and we take a combined $\delta \eta = 3\%$.

Best-fit spectra to XENON1T data is shown in Fig. 2. Taking the low $M_\nu$ limit ($M_\nu < \text{KeV}$), a minimal $\chi^2 = 41$ is obtained at $\sqrt{y_{\nu e} y_{\nu e}} = 0.96 \times 10^{-6}$ with the background being slightly down-scaled at $\eta = 1 = -4.5\%$. The best fit point yields a $\Delta \chi^2 = -6.7$ improvement over fixed $B_0$ fit ($\eta = 1$). The $\chi^2$ dependence on $\sqrt{y_{\nu e} y_{\nu e}}$ is plotted in Fig. 3 and the 2$\sigma$-preference threshold around the best-fit $\sqrt{y_{\nu e} y_{\nu e}}$ is shown as the dotted curve. With 28 degrees of freedom, the minimal reduced $\chi^2 / \text{dof} \# = 1.46$ corresponds to 93\% credence level (C.L.), which is not a perfect fit to the data below 30 KeV. This is due to fluctuations above 10 KeV that are still unaccounted for by the flat $^{214}$Pb background and the NSI contribution close to the detector threshold.

Note a trace abundance of tritium below calibrated level is also a possibility to account for the KeV range recoils [1]. The tritium induced recoil spectrum shape is similar to that from NSI and leads to degeneracy in explaining the low-energy rise. In the next session, we will show the significance of NSI contribution decreases in the fit to PANDAX results, by including experiment-reported tritium and other SM background at low-energy recoils.

V. FIT TO PANDAX

The PANDAX collaboration reported the latest data with 100.7 ton-day exposure and 2121 events selected in December last year [5]. There is also a rise at 3-7 KeV. In contrast to the XENON experiment, more background were taken into account. Including tritium, $^{127}$Xe, $^{85}$Kr and $^{222}$Rn. We also make a likelihood fit to the 24 binned data below 25 KeV, by combining these NSI-induced events with PANDAX’s best-fit background modeling,

$$\chi^2 = \sum_{ij} (\eta_j B_{0j} + N_{i^e}^{j\nu e} - N_{i^e}^{j\text{data}})^2 / (\delta N_i)^2 + \sum_j (1 - \eta_j)^2 / (\delta \eta_j)^2,$$  

(14)

where $i$ denotes data point and $j = 1, 2, 3, 4$ denote $^{127}$Xe, tritium, $^{85}$Kr and $^{222}$Rn background. Their best-fit value are respectively 80.8, 202.9, 1095, 735.6 [5]. $\eta$ is background floating parameters. Among this background, $^{127}$Xe and tritium were considered a major factor in low energy range electron recoil excess and their statistic uncertainty are $\delta \eta_1 = 21\%$, $\delta \eta_2 = 35\%$ [51]. The rest of them are considered mostly flat background which are calibrated at full energy region. Their statistic uncertainty is relatively much smaller than $^{127}$Xe and tritium so we don’t consider their float and set $\eta_3 = \eta_4 = 1$. The
PANDAX detector efficiency curve is taken from [58], which is degenerate with flat $^{222}$Rn background so the efficiency error is not also considered. Detector smearing curve is from [59], and we parameterized this curve as 

$$\delta_\epsilon = 0.2626 \sqrt{E} + 0.0426E$$

where $E$ is in KeV.

Fitting PANDAX data is shown in Fig. 4. The gray dotted line represent new physical signal with best-fit result. The background $^{127}$Xe is up-scaled at $\eta_1 - 1 = 2.3\%$ and tritium is down-scaled at $\eta_2 - 1 = -6.7\%$. The $\chi^2$ dependence on $\sqrt{y_e y_e}$ is plotted in Fig. 3. A minimal $\chi^2 = 21.2$ is obtained at $\sqrt{y_e y_e} = 1.1 \times 10^{-6}$ and this point yields a $\Delta \chi^2 = -1.6$ improvement over background-only fitting results ($\chi^2 = 22.8$). The $2\sigma$-preference threshold around the best-fit $\sqrt{y_e y_e}$ is shown as the dotted curve. As PANDAX considered more background channels, the low energy range electron recoil excess can be explained well with background-only fit. Our new physical contribution is consistent with background-only fit, and we can constrain the NSI coupling to $\sqrt{y_e y_e} < 1.4 \times 10^{-6}$.

![FIG. 4. Best-fit new physics event distribution (gray dotted) with $\sqrt{y_e y_e} = 1.1 \times 10^{-6}$, background $^{127}$Xe with $\eta_1 = 102.3\%$ and tritium with $\eta_2 = 93.3\%$. Electron recoil events are the total 100.7 ton-day data [3].](image)

A few comments are due for comparing the result from lepton-number violating scattering to those existing constraints at solar and reactor neutrino experiments. As long as Eq. (10) is satisfied, $d\sigma/dE_R \propto E_R^{-\frac{1}{2}}$ is a common kinematic feature that is also observed in lepton-number conserving ($s\bar{\nu}$) scalar and neutrino magnetic dipole moment [40, 42] scenarios. As most signal events are expected to be near-threshold, $s\bar{\nu}$, $s\bar{\nu} \nu$ and $\nu - \nu - \mu$ operators will lead to almost identical event distributions. This allows the scalar coupling bounds to be directly scaled from the existing $\nu_e$ limits. For instance, the signal events with $\nu_\mu = 2.8 \times 10^{-11} \mu_B$, which corresponds to BOREXINO [50] bounds, can be reproduced with $\sqrt{y_e y_e} = 1.2 \times 10^{-6}$ and it is shown as the shaded exclusion in Fig. 3.

As reactor neutrino is predominantly $\nu_0$ at short distances, reactor neutrino constraints can be circumvented by assuming $y_e^0$ in Eq. (4) only involve $\nu_\mu$ and $\nu_\tau$, at the

![FIG. 5. Minimal $\chi^2$ after marginalizing over $\eta_1$. $\sqrt{y_e y_e} < 1.4 \times 10^{-6}$ is $2\sigma$ favored range and consistent with background-only fit. The shaded exclusion region is inferred from the BOREXINO bound [50].](image)

VI. CONCLUSION

In light of the 2-7 KeV electron recoil excess at the XENON1T experiment, we propose a model in the $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B-L$ framework that provides a light scalar $s_1$ after $U(1)_B-L$ and electroweak symmetry breaking, and also generates $s_1 \bar{\nu} \nu$, $s \nu_0 \nu$ couplings via heavy fields above the $U(1)_B-L$ breaking scale. These couplings violate lepton number and lead to non-standard $\nu e \rightarrow \nu' e$ scattering. Solar MeV neutrinos may scatter off detector’s electron via NSI, and enhance low $E_R$ electron recoil event rate that explains the observed excess.

We calculate the recoil spectrum for Solar neutrino’s $\nu e \rightarrow \nu' e$ process, and compare the NSI spectrum in the light mediator limit ($M_s$ below KeV) to other scattering processes, which share a common $E_R^{-\frac{1}{2}}$ kinematic feature near the recoil threshold energy. A preferred coupling range $5.8 \times 10^{-7} < \sqrt{y_e y_e} < 1.1 \times 10^{-6}$ is obtained by fitting to the binned XENON1T data below 30 KeV, and
allowing the $^{214}\text{Pb}$ dominated background and detector efficiency to fluctuate at an 3% uncertainty. The optimal case $\sqrt{y_\nu y_\nu} = 0.96 \times 10^{-6}$ improves over the $^{214}\text{Pb}$ background fit to XENON1T by more than 2σ. The preferred $\sqrt{y_\nu y_\nu}$ range that explains the XENON1T excess is consistent with existing solar neutrino measurements, and by assuming flavored $\nu_\mu, \nu_\tau$ couplings, avoid the constraint from reactor neutrino experiments. The $y_\nu$ coupling however may require more sophisticated design to evade tension with stellar cooling bounds.

For PANDAX results, SM backgrounds including tritium and $^{127}\text{Xe}$ provide excellent fit to the shape of measured electron recoil spectrum. The inclusion of light $s_1$ mediated NSI signal only yield a slight improvement and is consistent with the background only fit. The NSI-induced event rise towards the lowest energy bin allows constraint to be placed the NSI couplings. PANDAX data requires $\sqrt{y_\nu y_\nu} < 1.4 \times 10^{-6}$ that is close to reactor experiment limits.

**Appendix A: $s\bar{\nu}\nu$ Scattering Amplitude**

For the $s\bar{\nu}\nu$ scattering on electrons, the amplitude is

$$iM = i \frac{y_\nu y_\nu}{(p_4 - p_2)^2 - M^2} \bar{u}(p_4)u(p_2)\bar{u}(p_3)u(p_1), \quad (A1)$$

$p_1, p_3$ are the initial and recoil electron 4-momenta. $p_2, p_4$ are the incident and ejected neutrino 4-momenta. The amplitude-square is

$$|M|^2 = \frac{y_\nu^2 y_\nu^2}{4(M_e^2 - t)^2} \text{tr}[(\rho_4 \cdot \rho_2)\text{tr}((\rho_3 + M_e)(\rho_1 + M_e))]$$

$$= \frac{y_\nu^2 y_\nu^2}{4(M_e^2 - t)^2} (4p_2 \cdot p_4)(4p_1 \cdot p_3 + 4M_e^2). \quad (A2)$$

Using $p_2 \cdot p_4 = -\frac{1}{2} t, p_1 \cdot p_3 = M_e^2 + M_E k = M_e^2 - \frac{1}{2} t$, where $t$ is the Mandelstam variable which is defined in the main text, then we get

$$|M|^2 = \frac{y_\nu^2 y_\nu^2}{t} (M_e^2 - t)^2. \quad (A3)$$

In the lab frame, the total cross-section is

$$\sigma = \int \frac{d^4 \rho_3 d^4 \rho_4}{(2\pi)^4 E_3 E_4} (2\pi)^4 \delta^4(p_1 + p_2 - p_3 - p_4)|M|^2. \quad (A4)$$

Integrating out $p_4$, we have

$$\sigma = \frac{1}{4M_e E_\nu} \int \frac{dE_3}{8\pi \rho_2^2} |M|^2. \quad (A5)$$

$E_3 = M_e + E_k, |\rho_2| = E_\nu$, where $E_k$ is the electron’s acquired kinetic energy after scattering. The differential cross-section is

$$\frac{d\sigma}{dE_k} = \frac{1}{4M_e E_\nu} \frac{1}{8\pi \rho_2^2} |M|^2$$

$$= \frac{y_\nu^2 y_\nu^2 E_k M_e (E_k + 2M_e)}{8\pi E_\nu^2 (M_e^2 + 2M_e E_k)^2}, \quad (A6)$$

as in Eq. [7]

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