Neutrino self-interactions and XENON1T electron recoil excess

Andreas Bally,† Sudip Jana,‡ and Andreas Trautner

Max-Planck-Institut für Kernphysik, Saupfercheckweg 1, 69117 Heidelberg, Germany

The XENON1T collaboration recently reported an excess in electron recoil events in the energy range between $1 - 7\,\text{keV}$. This excess could be understood to originate from the known solar neutrino flux, if neutrinos couple to a light vector-mediator with strength $g_{\nu N}$ that kinetically mixes with the photon with strength $\chi$, and $g_{\nu N} \chi \sim 10^{-13}$. Here, we show that such coupling values can naturally arise in a renormalizable model of long-range vector-mediated neutrino self-interactions. The model could be discriminated from other explanations of the XENON1T excess by the characteristic $1/T^2$ energy dependence of the neutrino-electron scattering cross section. Other signatures include invisible Higgs and $Z$ decays and lepto-philic charged Higgses at a few 100 GeV. ALPS II will probe part of the viable parameter space.

A recently publicized search for low-energy electronic recoil events performed with the XENON1T detector yielded some unexpected excess over background, statistically significant at the level of $3.2\sigma$ [1]. New physics explanations of this excess could include axions produced in the Sun [1], an unexpectedly sizable neutrino magnetic moment [1], inelastic semi-annihilation recoils [2], axion-like particle warm dark matter [3], or a fast Dark Matter component [4]. The solar axion and neutrino magnetic moment explanations are already practically ruled out, respectively by stellar cooling [5], or white dwarfs [6] and globular cluster cooling constraints [7, 8]. A more conventional origin of the excess would be an unaccounted Tritium background in the detector [1], or simply a statistical fluctuation. Notwithstanding this perception, here we proclaim a new physics explanation of the excess based on a UV-complete neutrino self-interaction model recently proposed by M. Berbig and two of the authors [9].

Elastic scattering of solar neutrinos off electrons is a subdominant background for low-energy electron recoil events in XENON1T with around 220 expected events [1]. The inclusion of a light mediator that couples neutrinos to electrons could change this conclusion and explain the observed excess, just as in the case of an enhanced neutrino magnetic moment. However, generic light mediators, especially when coupled to electrons, face severe constraints, see e.g. [10–12]. As in [9], we circumvent the most rigorous constraints from cosmology by relying on a low temperature phase transition in the neutrino sector [13]. However, different from [9], where mediator masses of $m_{Z'} \sim O(10)\,\text{eV}$ were considered, we here focus on a parameter region with $m_{Z'} \lesssim 10^{-4}\,\text{eV}$. The reason is that in order to accommodate the observed excess at XENON1T there needs to be a sufficiently strong interaction between electrons and the new mediator, which is excluded for the former region. In the new parameter region, the phase transition after which neutrinos start mixing with the hidden sector happens after recombination, i.e. below $T \sim 1\,\text{eV}$.

The UV-complete and renormalizable model has been presented in some detail in [9] and we will only highlight the most important features here. For the present analysis, relevant terms of the low energy effective Lagrangian in the mass basis read

$$L_{\text{eff}} = \frac{\varepsilon g_X}{\sqrt{2}} \bar{Z}'^\mu N \gamma^\mu \nu_L + \chi c_W Z'_\mu J_{\text{e.m.}}^\mu. \quad (1)$$

Here, $Z'$ is the gaugino boson of a new $U(1)_X$ gauge symmetry with coupling $g_X$, $N$ and $\nu_L$ denote respectively the new hidden and left-handed SM neutrinos while $\varepsilon$ is the neutrino-hidden-neutrino mixing, $c_W$ is the cosine of the electroweak mixing angle [58], $\chi$ the strength of gauge-kinetic mixing of $U(1)_X$ and $U(1)_Y$ induced by an operator $L_{\chi} = -s_\chi/2 B^\mu\nu X_{\mu\nu}$ [14, 15], and $J_{\text{e.m.}}^\mu$ the standard electromagnetic current. We only state (1) for a single flavor of neutrinos $\nu_L$ here, which is to be understood as a template for solar electron-neutrinos, while the extension to other flavors is straightforward.

FIG. 1: Feynman diagram for inelastic neutrino-to-hidden-neutrino up-scattering on electrons mediated by a light $Z'$. The couplings are generated by active-hidden neutrino mixing on the one side, and $U(1)_X - U(1)_Y$ kinetic-mixing on the other.

*andreas.bally@mpi-hd.mpg.de
†sudip.jana@mpi-hd.mpg.de
‡trautner@mpi-hd.mpg.de
From (1) it is possible to inelastically up-scatter neutrinos to hidden neutrinos on electrons in the XENON1T detector, see Fig. 1 (there is also elastic neutrino electron scattering, but this is suppressed by another insertion of $\epsilon$). Taking $m_{Z'} \ll m_e T$ and neglecting neutrino masses, the differential cross section for this process is given by (see e.g. [16])

$$\frac{d\sigma_{\nu e \to Ne}}{dT} = \frac{\varepsilon^2 g_X^2 \chi^2 e^2_{W}}{16 \pi m_e T^2} \left[ 1 + \left( \frac{T}{E_{\nu}} \right)^2 - \frac{m_e T}{E_{\nu}} \right].$$

(2)

Here $T$ is the electron recoil energy, $E_{\nu}$ the incident neutrino energy and $m_e$ the electron mass. Note the $1/T^2$ enhancement of the differential cross section at low recoil energies. This is akin to the $1/T$ enhancement in the scattering induced by a neutrino magnetic moment. With more data it should become possible to discriminate between the energy dependence of our model and a magnetic moment scattering, for example. Also note that without assuming $m_{Z'} \ll m_e T$, the actual form of the propagator relevant in (2) is $(2m_e T + m_{Z'}^2)^2$ and there is no enhancement for low-energy recoils, which practically rules out an explanation of the excess for mediators heavier than $\sqrt{m_e T}$.

The low-energy solar neutrino flux consists essentially of the continuous pp and discrete $^7$Be flux components [17],

$$\phi_{pp} = 5.94 \times 10^{10} \text{cm}^{-2}\text{s}^{-1},$$

(3)

$$\phi_{\text{Be}} = 4.86 \times 10^9 \text{cm}^{-2}\text{s}^{-1}.$$  

(4)

These neutrinos are affected by vacuum-dominated flavor oscillations resulting in a survival probability of [18]

$$P_{ee} = \cos^4 \theta_{13} \left( 1 - \frac{1}{2} \sin^2 2\theta_{12} \right) + \sin^4 \theta_{13},$$

(5)

for electron neutrinos arriving at Earth.

To compute our signal prediction, we take into account XENON1T detection and selection efficiency [1] $\epsilon(T)$ as well as the finite detector energy resolution by a gaussian smearing [19, 20]. Similarly to the XENON1T analysis we use the Free Energy Approximation

$$\frac{d\sigma_{\text{tot}}}{dT} = \sum_{i=1}^{54} \Theta(T - B_i) \frac{d\sigma_{\nu e \to Ne}}{dT},$$

(6)

to take into account Xenon electron binding energies. At keV energies this is a good approximation to more sophisticated computations [21, 22]. The differential event rate is then computed by the convolution

$$\frac{dN(T_\text{r})}{dT_\text{r}} = N_0 \times t \times \int dT dE_{\nu} \frac{d\phi(E_{\nu})}{dE_{\nu}} P_{ee} \frac{d\sigma_{\text{tot}}}{dT} \Theta \left( \frac{2E_{\nu}^2}{m_e + 2E_{\nu}} - T \right) \epsilon(T_\text{r}) g\text{Gauss}(T_\text{r},T),$$

(7)

where $T$ and $T_\text{r}$ are the actual and reconstructed electron recoil energies, respectively.

Fitting this to the observed excess we obtain Fig. 2. The best fit point has

$$\varepsilon g_X \chi = \left( 2.0^{+0.7}_{-0.5} \right) \times 10^{-13} \quad (95\%\text{C.L.}),$$

(8)

and is statistically preferred over the background-only hypothesis by $3\sigma$.

Note that nothing in our analysis prevents us from considering flavors other than electron-neutrinos in (1). This allows the possibility that also the subdominant non-electron-flavor solar neutrino flux contributes to the excess. If the absolute relevant flux changes by a factor $f$, it is straightforward see that our result in (8) should be rescaled by $1/\sqrt{f}$.

We now introduce our complete model in which $\mathcal{L}_\text{eff}$ and the parameter region (8) is naturally obtained. New particles and their charges under the new U(1)$_X$ gauge symmetry are shown in Tab. 1. We introduce a pair of SM-neutral but U(1)$_X$ charged chiral fermions $N_{1,2}$ and two new scalars $\Phi$ and $S$ [59]. New interaction terms for
the SM lepton doublet $L = (\nu_L, e_L)^T$ are given by

$$\mathcal{L}_{\text{new}} = -y \bar{L} \Phi N_1 - M N_1 N_2 + \text{h.c.} ,$$

where $\Phi := i \sigma_2 \Phi^*$, $y$ is a dimensionless Yukawa coupling, and $M$ has mass-dimension one. We only discuss the one-generation case here, with the extension to three generations of SM leptons and multiple generations of hidden fermions being straightforward. We consider the most general possible scalar potential, cf. [9] for details, and decompose the scalars as

$$H = \left( \frac{1}{\sqrt{2}} (h + i a_h) \right) , \quad \Phi = \left( \frac{1}{\sqrt{2}} (\phi + i a_\phi) \right) , \quad \text{and} \quad S = \frac{1}{\sqrt{2}} (s + i a_s) .$$

We assume that all neutral scalars obtain vacuum expectation values (VEVs) $v_3 := \langle \sigma \rangle$ for $\sigma = h, \phi, s, v_\theta$ spontaneously breaks EW symmetry, $v_3$ breaks $U(1)_X$, while $v_\theta$ breaks both. We will always assume the hierarchy $v_\theta \gg v_3, v_\phi$, which is required by the assumption of a light $Z'$. Next to the SM Higgs boson $H$, the physical scalar spectrum consists of a pair of heavy ($>100 \text{GeV}$) charged scalars $\Phi^\pm$, a pair of mass-degenerate heavy neutral scalar and pseudo-scalar $\Phi$ and $A$ ($|m_{\Phi, \pm} - m_{\Phi}| \lesssim 120 \text{ GeV}$ by electroweak precision), as well as a sub-keV light scalar $h_S$, all with very lepton-specific couplings and practically no coupling to quarks.

The photon is exactly the same massless combination of EW bosons as in the SM. The very SM-like $Z$ boson contains a miniscule admixture of the new gauge boson $X$,

$$Z_\mu = c_X (c_W W^3_\mu - s_W B_\mu) + s_X X_\mu ,$$

with an angle

$$s_X \approx -2 c_W \frac{g_b}{g_2} \left( \frac{v_\phi}{v_\theta} \right)^2 \ll 1 \quad \text{and} \quad c_X \approx 1 .$$

Gauge-kinetic mixing from the operator $\mathcal{L}_X$ shifts the $Z'$ coupling to the SM neutral current only by a negligible amount proportional to $\chi O(m_{Z'}^2/m_{Z}^2) [23, 24]$ (given $m_{Z'} \ll m_Z, \chi \ll 1$). However, it will introduce the coupling of $Z'$ to the electromagnetic current shown in Eq. (1), which is instrumental for our explanation of the XENON1T excess. The masses of the physical neutral gauge bosons are

$$m_Z \approx \frac{g_2 v_\theta}{2c_W} \quad \text{and} \quad m_{Z'} \approx g_X \sqrt{v_\phi^2 + v_\theta^2} =: g_X \bar{v} .$$

After $\phi$ assumes its VEV $v_\phi$, the yukawa coupling $y$ will introduce bi-linear mixing between SM neutrinos and the new hidden neutrinos. For exactly massless neutrinos this implies a mixing of $\nu_L$ and $N_2$ by an exact angle

$$\tan \epsilon = (g y v_\phi)/(\sqrt{2} M) .$$

Depending on the specific neutrino mass generation mechanism (see for example [27–29] for mechanisms compatible with the model), and in particular whether or not there is violation of lepton number, this will slightly change for massive neutrinos in which case there can also be a slight admixture of $N_1$ and (15) becomes approximate. In any case, $N$ in Eq. (1) should be understood as the resulting hidden-neutrino mass eigenstate which will have a (Dirac-)mass $M_N \approx (M^2 + y^2 v_\phi^2/2)^{1/2}$. By this mixing, SM neutrinos pick up a coupling to $Z'$ from the gauge coupling of $N_{1,2}$. This gives rise to the first term in (1) but also to a pure SM neutrino $Z'$ interaction proportional to $g_X \epsilon^2$. For temperatures $T \gg m_{Z'}$, while neutrino mixing via $\epsilon$ is relevant, $Z'$ will be effectively massless giving rise to an induced long-range four-neutrino interaction with thermally averaged rate $\Gamma \sim \epsilon^4 g_X^4 T$. Requiring this rate not to surpass the Hubble rate $H \sim T^2/M_P$ before recombination, but before today, yields

$$10^{-8} \lesssim \epsilon^2 g_X \lesssim 10^{-7} .$$

Clearly, this includes the assumption $m_{Z'} \lesssim 1 \text{ eV}$. We even focus on the region $m_{Z'} \lesssim 10^{-4} \text{ eV}$.

### Table I: New fields and their charges under SM and new $U(1)_X$ gauge symmetry.

| Field | $\Phi$ | $N_1$ | $N_2$ | $S$ | $X_\mu$ |
|-------|--------|-------|-------|-----|-------|
| SU(2)$_L \times U(1)_Y$ | $(2, \frac{1}{2})$ | 0 | 0 | 0 | 0 |
| $U(1)_X$ | +1 | +1 | -1 | +1 | 0 |
heavier \( m_{Z'} [10–12] \). Parametrizing \( m_{Z'} = g_X \bar{v} \) we can constrain the size of the effective \( U(1)_X \) breaking VEV \( \bar{v} \) to

\[
\bar{v} = \frac{m_{Z'}}{g_X} \lesssim \varepsilon^2 \times 10 \times \text{keV} .
\]

This fixes the necessary hierarchy between the relevant scales of the model to

\[
\xi := \bar{v}/v_h \lesssim \varepsilon^2 \times 4 \times 10^{-8} ,
\]

where \( v_h = 246 \text{ GeV} \) is the SM Higgs VEV. Stabilizing these hierarchies might require tuning in scalar quartic couplings which would not change any of our conclusions. Combining the requirements Eq. (8) and (16) and \( \chi \lesssim 5 \times 10^{-8} \) we obtain (the lower bound arises from \( g_X < 1 \))

\[
10^{-4} \lesssim \varepsilon \lesssim 2.5 \times 10^{-2} .
\]

This implies we automatically, obey constraints arising from violation of PMNS unitarity [30, 31] or direct search bounds [32–35]. The bounds on \( g_X \) and \( \varepsilon \) also imply a lower bound on \( \chi \gtrsim 10^{-10} \), which would be excluded for \( m_{Z'} \gtrsim 10^{-3} \text{ eV} \). This explains why it is not possible to fit the present excess for the parameter region considered in [9]. Note that we assume lower scales for \( v_\phi \) and \( v_s \) than in [9] but we do not strive to change the relative hierarchy of \( v_\phi \) and \( v_s \), parametrized by the angle

\[
\tan \gamma := v_\phi/v_s .
\]

This implies that despite our changes in \( m_{Z'}, g_X \) and \( \chi \) other very characteristic details of this model are exactly the same as in [9]. This includes key signatures \( H \to h_SH, H \to Z'Z', H \to ZZ' \) and \( Z \to Z'h_S \) whose rates are independent of \( m_{Z'}, g_X \), and \( \chi \) because they are fixed by Goldstone Boson Equivalence. These decays contribute to invisible \( H \) and \( Z \) decays at potentially observable levels, which already constrains \( s_\gamma \lesssim 0.2 [9] \).

Regarding big bang nucleosynthesis (BBN) constraints, none of the new light states \(( m_{Z'}, h_S, N) \) was in thermal equilibrium with the SM sufficiently before BBN, and between BBN and recombination, as required by BBN [36] and CMB constraints [37]. While thermal abundances of the light states is generated by heavy scalar exchange at temperatures above the electroweak scale, any such abundance would be depleted by reheating in the SM for example at the QCD phase transition. Still dangerous is the process \( e^+ e^- (\nu \bar{\nu}) \to NN \) via \( t \)-channel \( \Phi \Phi (\Phi, A) \) exchange. Absence of this process after QCD (EW) epoch requires \( g \lesssim 6 \times 10^{-4} v \xi_{h_S} / 100 \text{ GeV} \). Other BBN constraints related to \( Z' \) coupling to neutrinos do not apply here, simply because the new gauge interactions become important only after recombination. The now sizable up-scattering process of Fig. 1 is cosmologically irrelevant.

The leading direct constraint on the effective neutrino-\( Z' \) coupling arises from allowing unperturbed propagation of SN1987A neutrinos through the cosmic neutrino background (CνB) and implies \( \varepsilon^2 g_X \lesssim 5 \times 10^{-4} [38] \) (see also [39]). Furthermore, our constraint Eq. (16) already warrants that we are not violating the requirement of free-streaming neutrinos during CMB formation [40]. Laboratory constraints on \( m_{Z'} \) and \( g_X \) are not very limiting for light mediators (see references collected in [9]) and become even less relevant here as compared to [9] as the effective coupling to neutrinos here is smaller by an order of magnitude.

There are strong constraints on dark photon models and kinetic mixing from stellar cooling if the Higgs mode associated to \( U(1)_X \) breaking becomes light [41, 42]. The relevant Higgs mode in our model is \( h_S \) which has a sub-keV-scale mass \( m_{h_S} \approx \xi v_h \sqrt{2\lambda_S} \), but it is certainly much heavier than \( m_{Z'} \). So the stronger bounds of [41] (for the Higgsed case), which assume \( m_{h_S} \sim m_{Z'} \) do not apply at face value. A dedicated analysis in the context of our model would be required, which we expect to give the leading constraint on \( g_X \chi \) directly. We stress that it is the stellar cooling bounds that matter here, not the direct detection constraints, as the latter can always be avoided if \( Z' \) decays to (keV energy) neutrinos before arriving at the Earth, which is what generically happens in our model if \( m_{Z'} > 2m_\nu \), i.e. when \( Z' \) is not a dark matter candidate.

We note there is a parameter region around \( m_{Z'} \sim 5 \times 10^{-4} \text{ eV} \) and \( \chi \sim 10^{-3} \) (and \( m_{Z'} < 2m_\nu \)) where it is not excluded that our \( Z' \) could make up the entirety of the Dark Matter (see [12] and references therein). For smaller \( m_{Z'} \) this possibility is excluded. We stress though, that nothing in our resolution of the present XENON1T excess depends on the possibility of \( Z' \) being the Dark Matter. The parameter region \( \chi \gtrsim \approx 10^{-9} \) and \( m_{Z'} \gtrsim 5 \times 10^{-5} \text{ eV} \) will be probed by ALPS II [43].

Taking the expression for the mixing angle (15) one can show that

\[
M \lesssim (y/\sqrt{2}) \varepsilon s_\gamma \times 10 \times \text{keV} .
\]
However, since this is a practically invisible decay this should not leave an observable signature.

Note that for $m_{\nu_i} > m_{Z'} + m_{\nu_j}$ two- and for higher masses also three-body decays of SM neutrinos become possible, depending on the flavor structure of $y$. This fact, and the $Z'$-mediated four-neutrino interaction could substantially modify the cosmic neutrino background. Sufficiently fast decays would render it mono-generational, while the long-range neutrino self-interaction would modify the clustering even leading to neutrino condensation. In this case the possible coincidence of $m_{Z'} \sim m_{\nu} \sim T_{\text{CMB}}$ could become meaningful for the cosmological “why-now” problem.

In summary, we have outlined a possible explanation for the excess in electronic recoil events observed in XENON1T. In our scenario, the events are caused by inelastic neutrino up-scattering on electrons in the detector medium, induced by the standard solar neutrino flux. Effectively our explanation is based on a light vector mediator that, on the one hand, has a coupling of SM neutrinos to hidden neutrinos, and on the other hand, couples to the electromagnetic current. More specifically, these couplings are understood to originate from neutrino mass-mixing and gauge-kinetic mixing between a new $U(1)\times$ gauge symmetry with SM hypercharge. This model could be discriminated from other explanations of the XENON1T excess by the $1/T^2$ recoil-energy dependence of the differential cross section.

We have also presented an explicit gauge invariant, renormalizable and UV-complete model which realizes this explanation of the XENON1T excess without conflicting with observational constraints. The model would also lead to long-range neutrino self-interactions that could substantially modify the appearance of the cosmic neutrino background. More accessible signatures of the explicit model are new invisible Higgs and $Z$ decays and the presence of lepto-philic charged and neutral scalars with masses $O(100 \text{ GeV})$ which all could be searched for at the LHC and future colliders. Sterile neutrinos are required and they should mix with the SM neutrinos with an angle $\varepsilon > 10^{-4}$. The parameter space of the model can further be tested by searches for dark vectors and kinetic mixing, which is actively pursued, for example at ALPS II [43]. Leading limits on electron-neutrino scattering are set by TOXONO [44], Borexino [25] and GEMMA [26] experiments [45], while complementary regions of parameter space will also be probed by electron recoils at ongoing reactor neutrino experiments [16, 46] like CONUS [47], CONNIE [48] or $\nu$-eclus [49].

If the excess and our explanation holds up, this may open the exciting possibility of using neutrino-electron scattering to learn more about the mechanism behind neutrino mass generation and potentially strong neutrino self-interactions.

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Note added.—After our paper was submitted to, but before it appeared on the arXiv, Refs. [50–52] appeared which also treat the XENON1T excess. Especially [51] also considers neutrino-electron scattering by a light vector mediator, despite in a somewhat effective picture without complete model. Our results are consistent where overlapping, but we stress that mediator masses heavier than $m_{Z'} \gtrsim 1 \text{ eV}$ are excluded by Dark Photon constraints; which unavoidably brings long-range neutrino self-interaction into focus.

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[58] We abbreviate trigonometric functions of all angles by \( \sin \theta_i \equiv s_i \), \( \cos \theta_i \equiv c_i \), and \( \tan \theta_i \equiv t_i \) in this work.

[59] Similar models albeit in a completely different range of parameters have been conceived in [53–55] (see also [56, 57] for similar but somewhat incomplete models).