

Searching for beyond-standard-model solar neutrino interactions using directional detectors

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Micropattern gaseous detectors (MPGDs) are a class of technologies that enable the full three-dimensional spatial reconstruction of ionization tracks from nuclear and electron recoils in gas. Anticipating near-future 30 m³-scale time projection chambers with MPGD-based readout, we forecast the sensitivity of such directionally sensitive low-energy recoil detectors to neutrino interactions beyond the standard model. We work in the framework of neutrino nonstandard interactions (NSIs), and calculate the combined recoil energy-angle distributions of the electron recoil signal generated by solar neutrinos in atmospheric-pressure He:CF₄ gas. We estimate the expected exclusion limits that such an experiment could place on various NSI parameters, as well as the mass and coupling of a new light mediator that interacts with electrons and neutrinos. We find that with an achievable background reduction of around a factor of ten from current estimates for a 30 m³ optical-readout detector using this gas mixture, sensitivity to NSI parameters would already approach Borexino's sensitivity. Directionality also allows for event-by-event neutrino energy reconstruction, which would provide a means to resolve some parameter degeneracies present in the modified neutrino cross section in this formalism. Our results strongly motivate the development of small-scale directionally sensitive gas detectors for neutrino physics.

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I. INTRODUCTION

Almost 70 years since the first experimental detection of the neutrino, its fundamental nature and connection to the rest of the otherwise well-tested standard model (SM) of particle physics is still mysterious. Although neutrinos remain challenging to detect, the pursuit of answers to questions such as the origin of their mass and the precise phenomenology of their oscillations has accelerated the proliferation of many diverse classes of large and highly sensitive detectors, see e.g., Refs. [1–21] for an incomplete sample. Detectors are able to observe and study neutrinos from a range of sources at the precision level, extending from low-energy neutrinos created by human beings, to

ultra-high-energy neutrinos generated by extreme astrophysical environments, see e.g., Ref. [22] for a review.

Although most successive generations of neutrino observatories have *increased* in size over the years, technological developments have also allowed for the size of a detector that is capable of making competitive measurements of a neutrino source to decrease—for example by trading target mass for lower energy thresholds, better signal reconstruction, and enhanced background rejection. In the case of neutrinos produced by the Sun, the first hints of detection by the liquid-xenon-based XENONnT [23] and PandaX [24] experiments are setting new records for the smallest physical size of a detector able to observe neutrinos from an astrophysical source.

These recent detections are possible thanks to the combination of the large cross section for coherent elastic neutrino-nucleus scattering (CEνNS) [25,26] as well as the excellent nuclear/electron recoil discrimination and low backgrounds of modern dual-phase xenon time projection chambers—the latter advance having primarily been driven by the (so-far unsuccessful) hunt for galactic dark matter particles interacting with nuclei [27–29]. These results, when combined with the experimental measurements of CEνNS by COHERENT [30–32] and CONUS+ [33], are inspiring explorations into the possible manifestations of

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beyond-SM neutrino interactions at lower energies than have been studied before, see e.g., Refs. [34–54]. However, at the $\gtrsim$ keV recoil energy scale searched for by these low-background detectors, the CE ν NS process is actually responsible for a subdominant number of neutrino events. In terms of total numbers, elastic neutrino-*electron* scattering (E ν eS) due to the solar pp flux is a stronger signal by around an order of magnitude. Unfortunately, at present, these events are hidden beneath a large background of electron recoils due to radioactivity in the detector and environment, making it challenging to perform precision measurements of the neutrino-electron cross section. In the long term, future xenon experiments like XLZD [55–58] should be poised to set constraints on new physics manifesting in low-energy E ν eS at an interesting level of sensitivity (given estimates for the sensitivity of the current generation like LZ, XENONnT and PandaX [34,46,59–62]), but this will depend strongly on the level of electron recoil background that is achieved. At present, the best limits still come from Borexino [63–70], whose sensitivity mostly (but not entirely) supersedes [71,72] that of experiments using terrestrial neutrino sources like beam dumps or reactors [73–77].

Nevertheless, due to the relatively high energy threshold of Borexino when compared to dark matter detectors, attempting to test beyond-SM physics in the neutrino sector using low-energy electron recoils remains a well-motivated challenge. To meet this challenge, however, we require detectors that are able to accumulate a sufficient number of reliable neutrino events that are not buried beneath backgrounds which would otherwise preclude the robust measurements of their cross section required to discern differences from the SM prediction. A promising idea to achieve this—one which was proposed many years ago in this context [78–80],¹ but has only become experimentally feasible recently—is to attempt the full spatial reconstruction of the recoil track *directions*.

Although neutrino directionality is an important and well-established observable at higher energies [4,9,18,82], there are currently only a very limited number of techniques able to reconstruct this information in the regime relevant for the low-energy solar neutrino fluxes [83–86]. So far, directional detection of low-energy recoil tracks at the necessary $\mathcal{O}(10\text{--}100)$ keV energies has only been shown to be experimentally feasible for time-projection chambers (TPCs) using gaseous targets [87,88]. This technology has seen significant advances over the last few decades, driven partially by the goal of finding a dark matter detector that is able to circumvent [89–93] the so-called “neutrino fog” limit [28,94–102] imposed upon conventional direct detection experiments. It has been shown that a cost-balanced trade-off is possible using the high-resolution capabilities

of highly segmented micropattern gas detectors (MPGDs), which enable sub-mm spatial resolution and high signal-to-noise reconstructions of fully three-dimensional ionization tracks [88,103–105].²

The reconstruction of ionization tracks in gas can provide a large amount of event-level information about the recoils that generated them. However, the detrimental effects of diffusion require that the distance over which this ionization is drifted through the gas to the readout plane must be kept to a minimum, i.e., 25–50 cm at most. This in turn implies that a large-scale experiment of, say, 100 m³ volume or more, could only be achieved in a modular configuration composed of multiple 10 m³-scale detectors. Working towards the ultimate vision of a large-scale “recoil observatory” of this kind is the goal of both the CYGNUS Consortium [103,113–116] and the CYGNO Collaboration [117–121]. The scale of an experiment necessary for a dark matter search into the neutrino fog would be naturally preceded by several stepping-stone detectors of $\mathcal{O}(10\text{ m}^3)$ volume. These experiments—which are in development now [121]—are the focus of this study as they will already have a sufficient target mass and background level to obtain sensitivity to solar neutrinos via E ν eS.

As for the science case supporting a gas detector of this scale over, say, a nondirectional liquid-based detector, there are several notable advantages a directional measurement offers. In particular, monodirectional sources like the Sun or a neutrino beam are interesting because of the fixed kinematic relationship between the neutrino’s initial energy and the energy and scattering angle of the subsequent recoil. In principle, a simultaneous measurement of the recoil energy and angle (as enabled by MPGDs) allows for the event-by-event reconstruction of the neutrino energy spectrum. On top of this, a detector able to reconstruct the full spatial distributions of recoil-induced ionization clouds comes with a powerful intrinsic method for background rejection and nuclear/electron recoil discrimination, both of which have been shown to work down to ~ 10 keV electron recoil energies [115,117,122–128].

All of the advantages of directional gas detectors mentioned above inspired Ref. [129], which evaluated the detector requirements to perform solar neutrino flux spectroscopy using E ν eS. Here, we consider the extent to which the science case for this kind of measurement extends beyond this. As well as simply reconstructing the neutrino flux, there will be additional, more subtle information about the neutrino-electron cross section accessible to an experiment measuring angular information that might uncover corrections originating from new

¹See also Ref. [81] for a similar proposal in the context of geoneutrinos.

²This type of directional detection is referred to as “recoil imaging” [88,104] to distinguish it from other, less powerful types of directional detection that seek to only measure spatial projections of ionization [106,107] or to infer directionality indirectly [108–112], e.g., via daily modulation signals.

physics. So in this paper, we extend Ref. [129] by evaluating the size and performance of a gas-based directional recoil detector required to test whether or not neutrino-electron scattering kinematics continues to be described by the SM at low energies.

The paper is structured as follows: in Sec. II we describe the theoretical calculation of the direction-dependent scattering rate of solar neutrinos with electrons, within and beyond the SM. In Sec. III we describe our model for the detector technology, and discuss the relevant effects that limit its ability to reconstruct the energy and direction of neutrino events. In Sec. IV we explain our statistical methodology, and then present results in Sec. V.

II. NEUTRINO-ELECTRON SCATTERING

We begin by describing neutrino-electron elastic scattering, both within the SM and after allowing for the possibility of so-called “nonstandard interactions” (NSIs) due to the presence of new mediators beyond the SM. We also describe the helpful density matrix formalism, which we adopt to handle the effects of NSIs on neutrino scattering and oscillations in this formalism.

A. Directional event rate

For a detector containing N_T scattering targets (e.g., electrons or nuclei) per unit detector mass, the differential event rate for elastic neutrino scattering on those targets is given by,

$$\frac{dR}{dE_r} = N_T \int_{E_r^{\min}}^{E_r^{\max}} dE_\nu \frac{d\Phi}{dE_\nu} \sum_\alpha P_{e\alpha}(E_\nu) \frac{d\sigma_\alpha}{dE_r}. \quad (1)$$

Here $d\Phi/dE_\nu$ is the neutrino flux, E_r is the recoil energy, and E_ν is the neutrino energy. The quantity $d\sigma_\alpha/dE_r$ is the differential scattering cross section for the interaction of an α -flavored neutrino with the target particle. The neutrino-energy-dependent function $P_{e\alpha}(E_\nu)$ is the probability that the electron neutrino ν_e produced in the solar core is detected as a neutrino with flavor $\alpha = \{e, \mu, \tau\}$ after oscillations. It is not possible to identify any flavor information at the level of individual recoil events, so we sum over α when computing dR/dE_r .

To describe the directional dependence of the recoil event rate, we follow the discussion in Refs. [90,130]. For a detector with directional sensitivity, we expand the space of observables to include the recoil direction $\hat{q}_r$, in addition to the recoil energy, E_r . To find the distribution of recoil energies and angles we compute the double-differential event rate, $d^2R/dE_r d\Omega_r$, where $d\Omega_r$ is a solid angle element around $\hat{q}_r$. The expression for this is

$$\frac{d^2R}{dE_r d\Omega_r} = N_T \int dE_\nu d\Omega_\nu \sum_\alpha P_{e\alpha}(E_\nu) \frac{d^2\Phi}{dE_\nu d\Omega_\nu} \frac{d^2\sigma_\alpha}{dE_r d\Omega_r}. \quad (2)$$

The first quantity inside the integral is the double-differential neutrino flux, which in the case of the Sun can be written straightforwardly as

$$\frac{d^2\Phi}{dE_\nu d\Omega_\nu} = \frac{d\Phi}{dE_\nu} \delta(\hat{q}_\nu - \hat{q}_\odot), \quad (3)$$

where $\hat{q}_\nu$ is the incoming neutrino direction and $\hat{q}_\odot$ is the unit vector pointing from the Sun to the detector. As for the double-differential cross section, we can assume the scattering is azimuthally symmetric, which means our solid angle element is $d\Omega_r = 2\pi d\cos\theta_r$, or $d\Omega_r = 2\pi d\cos\theta_\odot$ if we work in a heliocentric coordinate system. The scattering kinematics relates this angle to the neutrino energy and the energy of the recoil via

$$\cos\theta_r = \hat{q}_\nu \cdot \hat{q}_r = \frac{E_\nu + m}{E_\nu} \sqrt{\frac{E_r}{E_r + 2m}}, \quad (4)$$

so that

$$\frac{d^2\sigma_\alpha}{dE_r d\Omega_r} = \frac{1}{2\pi} \frac{d\sigma_\alpha}{dE_r} \delta\left(\hat{q}_\nu \cdot \hat{q}_r - \frac{E_\nu + m}{E_\nu} \sqrt{\frac{E_r}{E_r + 2m}}\right), \quad (5)$$

where m is the mass of the target particle. Using Eqs. (3) and (5) in Eq. (2), we find

$$\begin{aligned} \frac{d^2R}{dE_r d\Omega_r} &= \frac{N_T}{2\pi} \int dE_\nu \sum_\alpha P_{e\alpha}(E_\nu) \frac{d\sigma_\alpha}{dE_r} \frac{d\Phi}{dE_\nu} \\ &\times \delta\left(\hat{q}_r \cdot \hat{q}_\odot - \frac{E_\nu + m}{E_\nu} \sqrt{\frac{E_r}{E_r + 2m}}\right). \end{aligned} \quad (6)$$

The δ -function in Eq. (6) can be rewritten as $\delta(x + \frac{1}{\mathcal{E}})$ where we define

$$x = -\frac{1}{E_\nu}, \quad y = \sqrt{\frac{m^2 E_r}{E_r + 2m}}, \quad \frac{1}{\mathcal{E}} = \frac{\hat{q}_r \cdot \hat{q}_\odot}{y} - \frac{1}{m}, \quad (7)$$

where $\hat{q}_r \cdot \hat{q}_\odot = \cos\theta_\odot$. We can then perform the δ -function integral over this new variable x to get

$$\frac{d^2R}{dE_r d\Omega_r} = \frac{N_T}{2\pi} \frac{\mathcal{E}^2}{y} \left(\sum_\alpha P_{e\alpha} \frac{d\sigma_\alpha}{dE_r} \frac{d\Phi}{dE_\nu} \right) \Big|_{E_\nu=\mathcal{E}}. \quad (8)$$

1. Oscillation probabilities

The Sun produces electron-flavored neutrinos ν_e at its core via nuclear processes [131]. The neutrino flavor states then evolve during propagation through the solar matter and in the vacuum until they are detected on Earth. The Hamiltonian for neutrino propagation from the core of the

Sun to the Earth in the flavor basis is given by

$$H = \frac{1}{2E_\nu} U \begin{pmatrix} 0 & 0 & 0 \\ 0 & \Delta m_{21}^2 & 0 \\ 0 & 0 & \Delta m_{31}^2 \end{pmatrix} U^\dagger + \sqrt{2} G_F N_e(x) \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad (9)$$

where U is the Pontecorvo-Maki-Nakagawa-Sakata matrix that transforms flavor eigenstates ν_α to mass eigenstates ν_i , and the mass-squared differences that govern the mixing are labeled as $\Delta m_{ij}^2 \equiv m_i^2 - m_j^2$. The term $\sqrt{2} G_F N_e(x)$ is the matter potential, capturing the coherent forward scattering of neutrinos in the interior of the Sun. This term depends upon the number density of electrons in the Sun, $N_e(x)$, which is a function of position x , which is measured relative to the center of the Sun.

At the detector, the neutrino flavor state ν_α can be expressed as a superposition of ν_i , as follows:

$$P_{e\alpha} = \sum_{i=1}^3 P^m(\nu_e \rightarrow \nu_i) |U_{\alpha i}|^2, \quad (10)$$

where $P^m(\nu_e \rightarrow \nu_i)$ is the probability of adiabatic transition of ν_e to ν_i in the Sun. For a more detailed discussion and derivation of Eq. (10), we refer to Refs. [132–134].

2. SM neutrino cross section

The expression for $d^2R/dE_r d\Omega_r$ derived above applies to both electron and nuclear scattering. However, we will now focus our attention towards $E\nu_e S$ for the remainder of this study. The motivation behind this is simply that the event rate from solar pp and ${}^7\text{Be}$ neutrinos is within reach of the $\sim 30 \text{ m}^3$ -scale gas detectors currently under preparation. Nuclear recoils from the ${}^8\text{B}$ neutrino flux could also be measured, but require much larger gas volumes ($\sim 1000 \text{ m}^3$) as well as nuclear-recoil energy thresholds that still remain challenging ($\sim \text{keV}_r$). We relegate a discussion of the prospects for measuring solar NSIs through $\text{CE}\nu\text{NS}$ to Appendix B.

The SM differential cross section for $E\nu_e S$ ($\nu_\alpha e^- \rightarrow \nu_\alpha e^-$) is

$$\frac{d\sigma_\alpha}{dE_r} = \frac{2G_F^2 m_e}{\pi} \left[g_{L\alpha}^2 + g_R^2 \left(1 - \frac{E_r}{E_\nu} \right)^2 - 2g_{L\alpha} g_R \frac{m_e E_r}{2E_\nu^2} \right], \quad (11)$$

where m_e is the electron mass, G_F is the Fermi constant, and g_L and g_R are the SM couplings. The process $\nu_e e^- \rightarrow \nu_e e^-$ can occur through both a charged current via W exchange, as well as a neutral current via a Z . On the other hand the process $\nu_{\mu/\tau} e^- \rightarrow \nu_{\mu/\tau} e^-$ is purely a neutral

current interaction. As a result, the flavor-dependent SM couplings are given by

$$\begin{aligned} g_{Le} &= 1 + \left(\sin^2 \theta_w - \frac{1}{2} \right), \\ g_{L\mu} &= g_{L\tau} = \sin^2 \theta_w - \frac{1}{2}, \\ g_R &= \sin^2 \theta_w, \end{aligned} \quad (12)$$

where $\sin^2 \theta_w = 0.2387$ is the weak mixing angle [135]. The differential cross section for ν_μ interactions with e^- is the same as ν_τ interactions at tree level, but there are percent-level differences between the two with higher-loop corrections (see e.g., Refs. [133, 134]). Since an experiment of the scale we are interested in here will not be sensitive to these differences, we neglect any radiative corrections to Eq. (11).

B. NSI formalism

We have now laid out the expected rate of electron recoils due to solar neutrinos under the SM. That calculation will form our null hypothesis, which we will test against a possible extension to the SM. The extension we consider is expressed in terms of the popular NSI framework. For concreteness, we will consider NSIs due to a new vector mediator and examine two limiting regimes for its mass. A “light mediator” refers to one whose mass is comparable to, or smaller than, the typical momentum transfer q involved in the scattering process, while a “heavy mediator” is one whose mass is considerably higher than q . We expect qualitatively similar results for axial-vector interactions, so for brevity, we do not explore these cases here (compare figures for these cases in e.g., Refs. [34, 46, 48–51, 72]). We emphasize that our motivation in this study is to determine whether the size, performance and background of a 30 m^3 -scale directional detector is sufficient to distinguish NSIs from the SM at a competitive level—an exhaustive comparison of phenomenologically similar models is unnecessary to arrive at such a conclusion.

Let us assume the existence of a new vector mediator Z' of a gauged $U(1)$ symmetry extending the SM. The mediator will couple to fermions f and neutrinos ν_α . The NSI Lagrangian in this case can be written as [62, 72, 136, 137]

$$\mathcal{L}^{\text{NSI}} \subset g_{Z'} \left(\sum_\alpha Q_{Z'}^{\nu_\alpha} \bar{\nu}_\alpha \gamma^\mu \nu_\alpha + \sum_f Q_{Z'}^f \bar{f} \gamma^\mu f \right) Z'_\mu. \quad (13)$$

Here, $g_{Z'}$ is the NSI coupling and $Q_{Z'}^i$ denotes the individual model-dependent vector charges. A few specific examples of anomaly-free realizations of such models are given in Table I. In this study, we will show results for only the $U(1)_{B-L}$ case, since the results for the various models are

TABLE I. $Q_{Z'}^i$ values for a few models.

Model	$Q_{Z'}^{e/\nu_e}$	$Q_{Z'}^{\mu/\nu_\mu}$	$Q_{Z'}^{\tau/\nu_\tau}$	$Q_{Z'}^{u/d}$
$B - L$	-1	-1	-1	1/3
$L_e - L_\mu$	1	-1	0	0
$L_e - L_\tau$	1	0	-1	0

very similar. In the light mediator limit, the free parameters of the model are $(g_{Z'}, m_{Z'})$, where $m_{Z'}$ is the mediator mass.

When the mass of the mediator significantly exceeds the momentum transfer of the interaction ($m_{Z'} \gg q$), the Lagrangian of Eq. (13) can be approximated as an effective four-fermion operator interaction similar to the SM Lagrangian [72,138–140],

$$\mathcal{L}^{\text{eff,NSI}} \subset -2\sqrt{2}G_F \sum_{f,P,\alpha,\beta} \epsilon_{\alpha\beta}^{f,P} (\bar{\nu}_\alpha \gamma_\mu L \nu_\beta) (\bar{f} \gamma^\mu P f), \quad (14)$$

where $P = \{L, R\}$ is the projection operator. Here, we have introduced the dimensionless effective coupling strengths $\epsilon_{\alpha\beta}^{f,P}$, which we take to be real valued. For vector NSIs in the heavy mediator limit, we take the free parameters of the model to be the following combination of L and R couplings,

$$\epsilon_{\alpha\beta}^f \equiv \epsilon_{\alpha\beta}^{f,L} + \epsilon_{\alpha\beta}^{f,R}. \quad (15)$$

In contrast to this, the axial-vector NSI coupling is $\tilde{\epsilon}_{\alpha\beta}^f \equiv \epsilon_{\alpha\beta}^{f,L} - \epsilon_{\alpha\beta}^{f,R}$. We will focus on the pure vector case in this work by setting the axial-vector coupling $\tilde{\epsilon}_{\alpha\beta}^f$ to zero, implying in turn that $\epsilon_{\alpha\beta}^{e,L} = \epsilon_{\alpha\beta}^{e,R}$.

We label flavors as $\alpha, \beta \in \{e, \mu, \tau\}$. Cases where $\alpha = \beta$ signify “flavor-conserving” interactions while $\alpha \neq \beta$ are “flavor-violating.” The models quoted in Table I are all flavor-conserving, but it is possible to include flavor-violating terms in analogy with Eq. (14), see e.g., Ref. [141].

C. Density matrix formalism

We now move on to describe the density matrix formalism, which we adopt to handle the fact that the introduction of NSIs modifies both the neutrino oscillations and scattering at the same time. For more details on this formalism, see Refs. [46,72,142].

We first write down the labeling for the scattering processes, which will be

$$\nu_\alpha T \rightarrow \nu_f T,$$

where T represents the scattering target, ν_α refers to the neutrino flavor state taking part in the scattering, and ν_f is the final neutrino flavor state. As mentioned above,

NSIs open up the possibility of flavor-violating terms where $f \neq \alpha$.

In addition to scattering, neutrino oscillations allow the neutrino flavor at detection (ν_α) to be different from its flavor at production— ν_e for the case of solar neutrinos—which we label as

$$\nu_e \rightarrow \nu_\alpha.$$

In order to account for the combined effect of neutrino propagation and neutrino detection, one needs to combine the two processes above, summing over the intermediate state ν_α .

Let S_{int} and S_{prop} be the S -matrices for scattering and oscillation, respectively. The square of the amplitude of the combined process can be written as

$$|A_{\nu_e \rightarrow \nu_f}|^2 = \left| \sum_\alpha \langle \nu_f | S_{\text{int}} | \nu_\alpha \rangle \langle \nu_\alpha | S_{\text{prop}} | \nu_e \rangle \right|^2. \quad (16)$$

Any individual detected recoil event could be associated with any neutrino flavor. So we take the sum over amplitudes for all f when computing the rate, as follows:

$$\frac{dR}{dE_r} \propto \sum_f \left| \sum_\alpha \langle \nu_f | S_{\text{int}} | \nu_\alpha \rangle \langle \nu_\alpha | S_{\text{prop}} | \nu_e \rangle \right|^2. \quad (17)$$

We can manipulate this expression into a more convenient form, as follows:

$$\begin{aligned} & \sum_f \left| \sum_\alpha \langle \nu_f | S_{\text{int}} | \nu_\alpha \rangle \langle \nu_\alpha | S_{\text{prop}} | \nu_e \rangle \right|^2 \\ &= \sum_{\alpha\beta} \sum_{\gamma\lambda} (S_{\text{prop}})_{\alpha\gamma} \rho_{\gamma\lambda}^e(0) (S_{\text{prop}})_{\beta\lambda}^* \sum_f (S_{\text{int}})_{f\beta}^* (S_{\text{int}})_{f\alpha} \\ &\equiv \sum_{\alpha\beta} \rho_{\alpha\beta}^e(L) \left(\frac{d\xi}{dE_r} \right)_{\beta\alpha} = \text{Tr} \left[\rho^e(L) \frac{d\xi}{dE_r} \right], \end{aligned} \quad (18)$$

where $(S_i)_{\alpha\beta} = \langle \nu_\alpha | S_i | \nu_\beta \rangle$ and we have used the fact that $\sum_\alpha |\nu_\alpha\rangle \langle \nu_\alpha| = 1$. The term $\rho_{\gamma\lambda}^e(0) = \langle \nu_\gamma | \nu_e \rangle \langle \nu_e | \nu_\lambda \rangle$ is the density matrix element describing the initial neutrino system at production. The quantity $d\xi/dE_r$ is the generalized scattering cross section, which accounts for the correlation between different neutrino flavor-scattering processes. The elements of this quantity can be written as

$$\begin{aligned} \left(\frac{d\xi}{dE_r} \right)_{\alpha\beta} &= \sum_{f=e,\mu,\tau} (S_{\text{int}})_{f\alpha}^* (S_{\text{int}})_{f\beta} \\ &\propto \mathcal{M}^*(\nu_\alpha \rightarrow \nu_f) \mathcal{M}(\nu_\beta \rightarrow \nu_f), \end{aligned} \quad (19)$$

where the two matrix elements, $\mathcal{M}^*(\nu_\alpha \rightarrow \nu_f)$ and $\mathcal{M}(\nu_\beta \rightarrow \nu_f)$ describe the processes $\nu_\alpha e^- \rightarrow \nu_f e^-$ and $\nu_\beta e^- \rightarrow \nu_f e^-$, respectively.

Finally, the neutrino density matrix $\rho^e(L)$ at the detector after propagation can be written as

$$\begin{aligned}\rho^e(L) &= S_{\text{prop}} \rho^e(0) S_{\text{prop}}^\dagger \\ &= e^{-iHL} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} e^{iHL},\end{aligned}\quad (20)$$

or, if labeling the matrix elements,

$$\rho_{\alpha\beta}^e(L) = (S_{\text{prop}})_{\alpha e} (S_{\text{prop}})_{\beta e}^*. \quad (21)$$

Here, H represents the Hamiltonian of propagation (see below) while L is the distance traveled by the neutrinos. The resulting density matrix after oscillation is no longer diagonal. Physically, this implies that the incoming neutrino at the detector before scattering is in a superposition of flavor states.

1. Oscillation probability with NSIs

The inclusion of NSIs modifies the vector potential influencing neutrino oscillations in matter, cf. Eq. (9). The Hamiltonian for neutrino propagation from the core of the Sun to the Earth in the flavor basis becomes

$$\begin{aligned}H &= \frac{1}{2E_\nu} U \begin{pmatrix} 0 & 0 & 0 \\ 0 & \Delta m_{21}^2 & 0 \\ 0 & 0 & \Delta m_{31}^2 \end{pmatrix} U^\dagger \\ &+ \sqrt{2} G_F N_e(x) \begin{pmatrix} 1 + \epsilon_{ee}(x) & \epsilon_{e\mu}(x) & \epsilon_{e\tau}(x) \\ \epsilon_{e\mu}^*(x) & \epsilon_{\mu\mu}(x) & \epsilon_{\mu\tau}(x) \\ \epsilon_{e\tau}^*(x) & \epsilon_{\mu\tau}^*(x) & \epsilon_{\tau\tau}(x) \end{pmatrix}. \end{aligned}\quad (22)$$

The $\epsilon_{\alpha\beta}$ are related to the NSI free parameters $\epsilon_{\alpha\beta}^f$ defined in Eq. (15) as follows³:

$$\epsilon_{\alpha\beta}(x) = \epsilon_{\alpha\beta}^e + (2\epsilon_{\alpha\beta}^u + \epsilon_{\alpha\beta}^d) + \frac{N_n(x)}{N_e(x)} (\epsilon_{\alpha\beta}^u + 2\epsilon_{\alpha\beta}^d), \quad (23)$$

where $N_n(x)$ is the neutron number density, and we have implicitly assumed the proton number density $N_p(x) = N_e(x)$ for neutral matter. The above expression is already appropriate for describing NSIs in the heavy mediator limit. However, since the flavor conversion of neutrinos in the presence of matter gets affected by the coherent forward neutrino-matter scattering, we work in the $q^2 \rightarrow 0$ limit when accounting for the effect of NSIs on oscillations. This makes it possible to use the same expressions also in the light mediator parametrization—we simply need to

translate $(g_{Z'}, m_{Z'})$ into an equivalent set of (flavor-conserving in our case) $\epsilon_{\alpha\alpha}^f$ parameters, as follows:

$$\epsilon_{\alpha\alpha}^f = \frac{g_{Z'}^2 Q_{Z'}^f Q_{Z'}^{\nu_\alpha}}{\sqrt{2} G_F m_{Z'}^2}. \quad (24)$$

Working in the one mass dominance (OMD) approximation ($\Delta m_{31}^2 \rightarrow \infty$), the density matrix elements can be calculated using an effective three-flavor oscillation scheme. The (S_{prop}) in Eq. (20) is given by

$$(S_{\text{prop}})_{\alpha e} = O_{e1} [O_{\alpha 1} S_{11}^{(2)} + O_{\alpha 2} S_{21}^{(2)}] + O_{\alpha 3} O_{e3} e^{-i\Delta_{33}L}. \quad (25)$$

More details about how $S_{11}^{(2)}$ and $S_{21}^{(2)}$ are calculated can be found in Appendix A.

After taking the average over the distribution of production locations in the Sun (which differs for each neutrino flux), the density matrix elements for general flavor indices $\alpha, \beta \in \{e, \mu, \tau\}$ become

$$\begin{aligned}\rho_{\alpha\beta}^e(L) &= (O_{e1})^2 [(O_{\alpha 1} O_{\beta 1}) P_{ee}^{2\nu} + (O_{\alpha 2} O_{\beta 2}) P_{\text{ext}} \\ &+ (O_{\alpha 1} O_{\beta 2}) P_{\text{int}} + (O_{\alpha 2} O_{\beta 1}) P_{\text{int}}^*] \\ &+ (O_{\alpha 3} O_{\beta 3}) (O_{e3})^2.\end{aligned}\quad (26)$$

The terms $O \equiv R_{23} R_{13}$, $P_{ee}^{2\nu}$, P_{int} and P_{ext} are defined as

$$O = \begin{pmatrix} c_{13} & 0 & s_{13} \\ -s_{13}s_{23} & c_{23} & c_{13}s_{23} \\ -s_{13}c_{23} & -s_{23} & c_{13}c_{23} \end{pmatrix}, \quad (27)$$

$$P_{ee}^{2\nu} = \frac{1}{2} (1 + \cos(2\theta_{12}) \langle \cos(2\theta_m) \rangle), \quad (28)$$

$$P_{\text{int}} = -\frac{1}{2} \langle \cos(2\theta_m) \rangle \sin(2\theta_{12}) e^{i\delta_{\text{CP}}}, \quad (29)$$

$$P_{\text{ext}} = \frac{1}{2} (1 - \cos(2\theta_{12}) \langle \cos(2\theta_m) \rangle), \quad (30)$$

where $s_{ij} = \sin \theta_{ij}$, $c_{ij} = \cos \theta_{ij}$, θ_{ij} are the vacuum mixing angles and δ_{CP} is the CP -violating phase. We use the best fit values given in Ref. [143]. θ_m is the effective mixing angle in matter, calculated by finding the eigenvalues of the Hamiltonian in Eq. (22) in the OMD approximation, detailed further in Appendix A. The terms containing factors of $e^{i\Delta_{33}L}$ rapidly oscillate as a function of neutrino energy, and the resulting density matrix elements after averaging do not depend on them.

The electron flavor diagonal density matrix component is calculated to be

³Note that we take all ϵ parameters to be real valued in this study, although in general they may be complex valued.

$$\rho_{ee}^e(L) = s_{13}^4 + c_{13}^4 P_{ee}^{2\nu} = P_{ee}^{3\nu}. \quad (31)$$

Hence, the diagonal terms of the density matrices, ρ_{ee} , $\rho_{\mu\mu}$ and $\rho_{\tau\tau}$ correspond to the measurement probabilities of the system, i.e., P_{ee} , $P_{e\mu}$ and $P_{e\tau}$, respectively. These will equate to the SM case, when all $\epsilon_{\alpha\beta}$ parameters in Eq. (22) are set to zero.

2. Cross section

For E ν S we set $f = e$ in the Lagrangians of Eqs. (13) and (14). In this work, we consider purely vector couplings for which $g_L = g_R$. The effects of the light mediator NSI in the cross section can be included by modifying the couplings in Eq. (11) [62,72], as follows:

Light mediator:

$$g_{L/R}^{\text{NSI}} = g_{L/R}^{\text{SM}} + \frac{g_{Z'}^2 Q_{Z'}^e Q_{Z'}^{\nu_\alpha}}{2\sqrt{2}G_F(2m_e E_r + m_{Z'}^2)}, \quad (32)$$

where the charges Q are chosen from Table I. In this case, it is readily apparent that we recover the SM when we switch off NSIs by setting $g_{Z'} = 0$.

For the heavy mediator case, to account for flavor-violating effects, we replace the SM cross section with the generalized scattering cross section [46,72,142], as follows:

Heavy mediator:

$$\begin{aligned} \left(\frac{d\zeta}{dE_r}\right)_{\alpha\beta}^{\text{NSI}} = & \frac{2G_F^2 m_e}{\pi} \sum_f \left[G_{\alpha f}^L G_{f\beta}^L \right. \\ & + G_{\alpha f}^R G_{f\beta}^R \left(1 - \frac{E_r}{E_\nu}\right)^2 \\ & \left. - (G_{\alpha f}^L G_{f\beta}^R + G_{\alpha f}^R G_{f\beta}^L) \frac{m_e E_r}{2E_\nu^2} \right], \end{aligned} \quad (33)$$

where

$$G_{\alpha\beta}^L = g_{L\beta} \delta_{\alpha\beta} + \epsilon_{\alpha\beta}^{e,L}, \quad G_{\alpha\beta}^R = g_{R\beta} \delta_{\alpha\beta} + \epsilon_{\alpha\beta}^{e,R}. \quad (34)$$

If we turn off NSIs by setting $\epsilon_{\alpha\beta}^{e,L}$ and $\epsilon_{\alpha\beta}^{e,R}$ to zero, the generalized cross section becomes diagonal, and the diagonal components $\left(\frac{d\zeta}{dE_r}\right)_{ee}$, $\left(\frac{d\zeta}{dE_r}\right)_{\mu\mu}$ and $\left(\frac{d\zeta}{dE_r}\right)_{\tau\tau}$ correspond to the SM cross sections $\frac{d\sigma_e}{dE_r}$, $\frac{d\sigma_\mu}{dE_r}$ and $\frac{d\sigma_\tau}{dE_r}$ of Eq. (11), respectively. Note that the label $\alpha\beta$ does not mean that α is the initial and β the final neutrino state—they are both initial-state indices, as shown in Eq. (19).

3. Double-differential event rate with NSI

A change of neutrino flavor during a scattering process is disallowed in the SM. In terms of Eq. (19), this means that the final neutrino state f should be the same as the initial

neutrino states α and β ; i.e., $f = \alpha = \beta$. This implies that the generalized cross section is diagonal, and as seen in Sec. II C 2, these diagonal elements correspond to the SM cross sections. We also saw in Sec. II C 1 that the diagonal elements of the density matrix ρ^e correspond to the transition probabilities $P_{e\alpha}$. Thus, if we were to express our SM event rate in terms of the density matrix and generalized cross section, it would reduce to a product of transition probability and cross section, as follows:

$$\begin{aligned} \left(\frac{d^2 R}{dE_r d\Omega_r}\right)^{\text{SM}} & \propto \text{Tr} \left[\rho^e \left(\frac{d\zeta}{dE_r} \right) \right]^{\text{SM}} \\ & \propto \sum_{\alpha\beta} \left[\rho_{\alpha\beta}^e \left(\frac{d\zeta}{dE_r} \right)_{\beta\alpha} \delta_{\alpha\beta} \right] = \sum_{\alpha} P_{e\alpha} \frac{d\sigma_{\alpha}}{dE_r}. \end{aligned}$$

However, with NSI effects included, we can have $f \neq \alpha \neq \beta$ which means that the generalized cross section can have off-diagonal terms. To account for all these processes, we must express the NSI-modified event rate as

$$\left(\frac{d^2 R}{dE_r d\Omega_r}\right)^{\text{NSI}} \propto \text{Tr} \left[\rho^e \left(\frac{d\zeta}{dE_r} \right) \right]^{\text{NSI}}. \quad (35)$$

This expression is therefore a more general expression for event rate calculations—it can be used for both SM flavor-conserving interactions, as well as NSI flavor-violating interactions.

So in summary, we rewrite the full expression for the double-differential event rate that we wrote down in Eq. (8) in terms of a density matrix and generalized cross section, as follows:

$$\frac{d^2 R}{dE_r d\Omega_r} = \frac{N_T \mathcal{E}^2}{2\pi y} \left(\text{Tr} \left[\rho^e \frac{d\zeta}{dE_r} \right] \frac{d\Phi}{dE_\nu} \right) \Big|_{E_\nu = \mathcal{E}}. \quad (36)$$

Finally, we note that we must also calculate this rate for each solar neutrino flux separately, as opposed to defining $\Phi = \Phi_{pp} + \Phi_{\text{Be}} + \dots$ as the sum of all fluxes. This is because the density matrix depends on the neutrino production profiles within the Sun, which are different for each flux. For electron recoils, we consider both pp and ${}^7\text{Be}$ fluxes, with pp contributing around 75% of the rate. The remaining solar fluxes contribute less than a handful of events, given the detector volumes we are interested in here. For nuclear recoils (Appendix B), we consider only the ${}^8\text{B}$ flux. For the full electron and nuclear recoil rates across all of the solar neutrino fluxes, as well as prospects for their detection and spectroscopy using directional gas detectors, we refer to our previous study [129].

D. Comparison of SM and NSI total rates and angular recoil spectra

In general, the inclusion of nonzero $\epsilon_{\alpha\beta}^e$ leads to an enhancement in the cross section and so a larger event rate

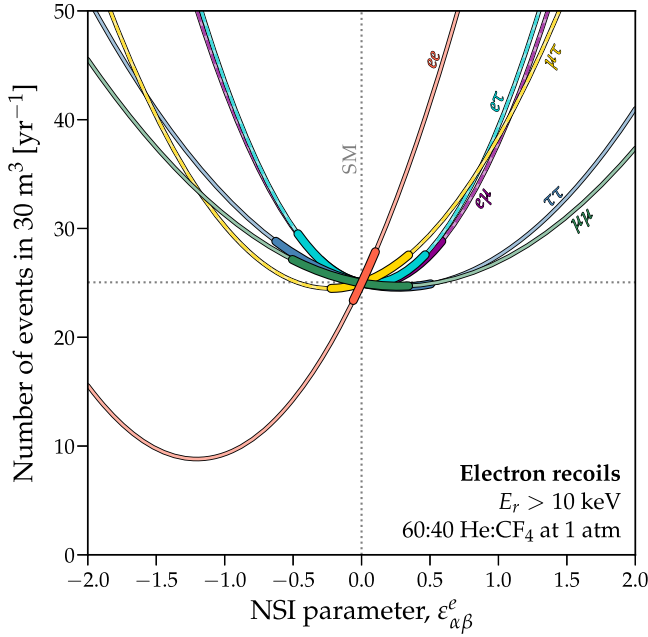


FIG. 1. Yearly expected number of electron recoil events with energies $E_r > 10$ keV, in a 30 m^3 gas detector containing a 60:40 mixture of He and CF_4 at atmospheric pressure. Each line shows the total rate as a function of a different $\epsilon_{\alpha\beta}^e$ where all other values of $\epsilon_{\alpha\beta}^e$ are kept at zero. The case for the SM is shown with a dotted gray line. The thick part of each line shows the range of ϵ allowed by Borexino at 90% CL. The event numbers shown here correspond to the sum over the pp and ${}^7\text{Be}$ fluxes.

compared to the SM prediction, roughly scaling as $R \propto (\epsilon_{\alpha\beta}^e)^2$. That said, it is possible for interference between positive and negative terms in the cross section, leading to a suppression in the rate relative to the SM.

In Fig. 1, we display the total expected number of electron recoils per year in a 30 m^3 gas volume containing a mixture of He and CF_4 at atmospheric pressure. This choice of volume and gas mixture is motivated further in the next section—for now, we focus on the dependence of the total rate on the NSI parameters, where Fig. 1 applies to the heavy mediator limit. Here we consider varying one NSI parameter at a time, with all others fixed at zero, including $\epsilon_{\alpha\beta}^{u,d}$. We see that, in general, turning the NSI parameters on leads to an enhancement in the total rate. The most obvious exception is $\epsilon_{ee}^e \approx -1.2$ for which there is a partial cancellation in the cross section. Cancellations also occur for the other couplings but are very small in magnitude and occur for small values of ϵ , so they are not easily visible on the plot. For comparison, we have highlighted the current range of $\epsilon_{\alpha\beta}^e$ allowed at 90% CL, from an analysis of Borexino’s Phase II spectral data [72], under the same assumption that only one NSI coupling is active at a time.

From Fig. 1, we can already conclude that while the event numbers are at an observable level in an experiment

of this scale, obtaining statistics needed to discriminate against the SM for $\epsilon_{\alpha\beta}^e \lesssim 0.5$ will require either a substantial increase in exposure (volume $\times$ time) or an extremely low electron recoil background. However, as we will motivate now, the ability of MPGD-based experiments to reconstruct both the recoil energy and the recoil direction independently will greatly alleviate the required background level.

To highlight the effect of the NSI parameters on the energy-angle distribution, Figs. 2 and 3 display the double-differential event rates. In this case, we are showing the double-differential event rate in heliocentric coordinates, i.e., in terms of $\cos \theta_\odot$, where $\cos \theta_\odot = 1$ corresponds to a recoil pointed in the direction of the solar vector, $\hat{q}_\odot$. The side panels of Fig. 2 also show the differential event rate as a function of $\cos \theta_\odot$ and E_r alone, i.e., after integrating $d^2R/dE_r d\cos \theta_\odot$ over energy and angle, respectively. Both panels in Fig. 2 demonstrate that the recoil distribution remains pointed towards the solar direction. Recall that energy and angle are correlated: the highest energy recoils carry more of the incoming neutrino’s forward momentum, so they will have small scattering angles or values of $\cos \theta_\odot$ close to 1. When NSIs are turned on, the event rate is generally enhanced more at low energies than at high energies, which means the angular distribution becomes slightly less focused towards the solar direction. We show two illustrative cases, one for the heavy mediator limit and one for the $B-L$ light mediator, but we have checked many different configurations and find that the behavior seen in Fig. 2 is generic.

To illustrate this latter point further, Fig. 3 shows the relative difference between the SM and the heavy mediator NSI cases, again when each coupling $\epsilon_{\alpha\beta}^e$ is turned on, while all others are fixed at zero. The blue colors show regions of this space of observables where the NSIs lead to a suppression in the rate compared to the SM, while pink regions are where it receives an enhancement. We note that while the general behavior is the same, there are subtle differences in the effects of each set of $\alpha\beta$ on $d^2R/dE_r d\cos \theta_\odot$. These differences are not present at the level of the total event rate (Fig. 1), where turning on different $\epsilon_{\alpha\beta}^e$ couplings can lead to identical values for the total expected number of events. This result motivates the use of a detector able to measure both E_r and $\cos \theta_\odot$ to attempt to search for the presence of NSIs and identify which interactions may or may not be active.

III. DIRECTIONAL ELECTRON-RECOIL DETECTION WITH GAS TPCS

We now describe how we model a directional gas-based detector. The first quantity to fix is the target medium. In line with the promising simulations and sensitivity studies performed in the context of the CYGNO experiment, we will assume the use of atmospheric pressures of He: CF_4 . This gas mixture is an electron drift gas and can allow high

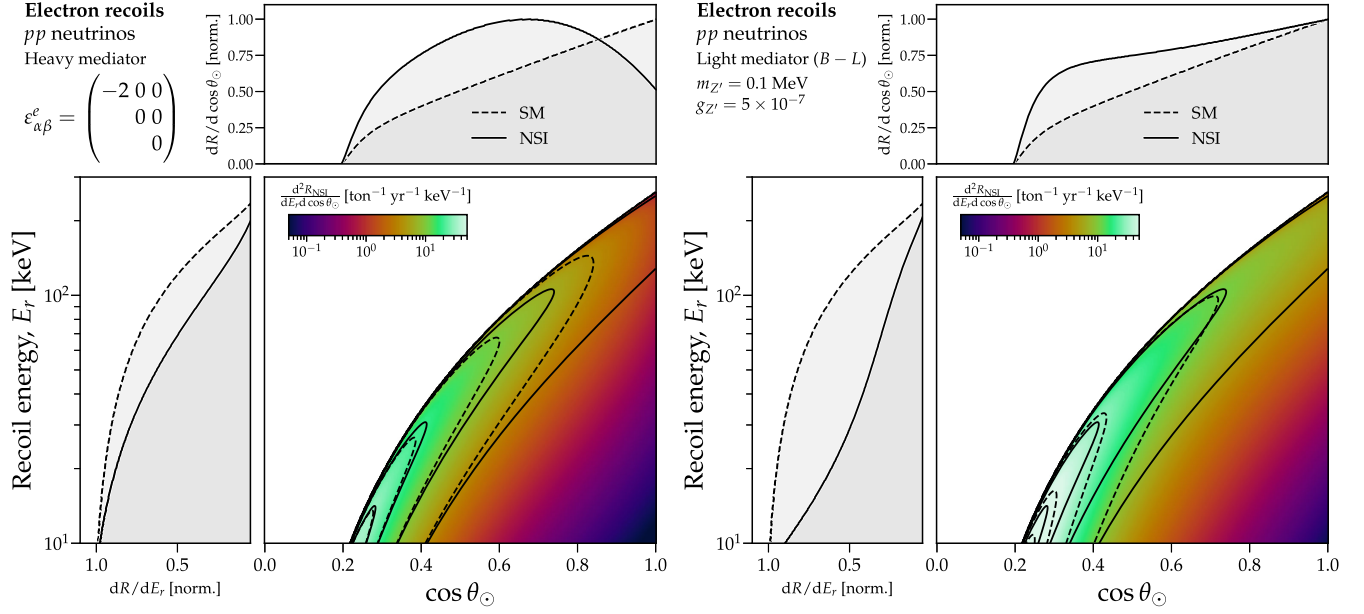


FIG. 2. Two illustrative cases of the double-differential event rates as a function of recoil energy, E_r , and recoil angle with respect to the Sun, $\cos \theta_\odot$. The central colored panels display $d^2 R / dE_r d \cos \theta_\odot$ for the NSI case. We assign a set of $\epsilon_{\alpha\beta}^e$ as labeled for the heavy mediator example (left), and $(m_{Z'}, g_{Z'})$ as labeled for the $B-L$ light mediator example (right). The smaller panels to the left and above the main panel in each case show the event rate integrated over angle and energy, respectively. In all panels, the solid lines correspond to the NSI case while the dashed lines show the SM ($\forall \epsilon_{\alpha\beta}^e = 0$ or $g_{Z'} = 0$ for heavy and light mediators, respectively). In all of these cases, the lines correspond to the rate after it has been normalized by the maximum value so as to highlight differences in shape (although the total rate also differs, as shown in Fig. 1). In the main panels, the solid and dashed contours also delineate equal values of the rate normalized by the maximum value—the four levels correspond to 0.75, 0.5, 0.25 and 0.1.

gas gains of $\mathcal{O}(10^6)$ with gas electron multipliers (GEMs), at the minor cost of worse diffusion compared with negative-ion drift gases like SF_6 . At the relevant energies, electron recoils have much longer tracks than nuclear recoils, but lower ionization densities, which means enhancing signal-to-noise is preferable over spatial resolution. The CF_4 gas also scintillates in the visible range, which as well as enabling the use of an optical readout (i.e., cameras), also means the light can be timed to use for reconstruction of a track's angle with respect to the drift direction orthogonal to the readout. This is the strategy for full three-dimensional spatial reconstruction adopted by CYGNO.

The addition of helium to the CF_4 has many advantages. In particular, it provides more scattering sites without substantially increasing the gas density, and so degrading the directional performance. It also allows the detector to operate at atmospheric pressure, simplifying the design and engineering of the gas vessel.

In Ref. [129], we motivated the use of a 60:40 ratio for $\text{He}:\text{CF}_4$ at atmospheric pressure for solar neutrino detection. Since preparations are underway by the CYGNO Collaboration for a 30 m^3 scale detector using this gas mixture, we will fix $30 \text{ m}^3 \times 5 \text{ years}$ as a near-term exposure benchmark, but will entertain further scale-ups if competitive sensitivity to some physics channel requires

it. Since ongoing design studies are now focused primarily on materials screening and background reduction, our discussion will primarily focus on the impact of the total background rate R_{bg} as opposed to total volume. Discussion of our assumed background model can be found in Sec. III B.

A. Detector performance

Having derived the input theoretical double-differential event rate $d^2 R / dE_r d\Omega_r$ in Sec. II, we now describe how these event rates can be converted into functions that more closely resemble the space of observables for a gas-based time-projection chamber.

The two main quantities a direction-sensitive detector aims to reconstruct are the total recoil energy and the initial recoil direction. This is done by reconstructing the three-dimensional ionization cloud produced when the recoil deposits energy as it slows down in the detector medium. The full three-dimensional ionization distribution is typically obtained by combining a two-dimensional projection onto the highly pixelated MPGD readout, with the third axis relative to that plane reconstructed using timing. Depending on the type of readout involved, the arriving charge can either be detected directly (electronic readout) or the scintillation produced during the amplification stage can be imaged with high-resolution cameras (optical

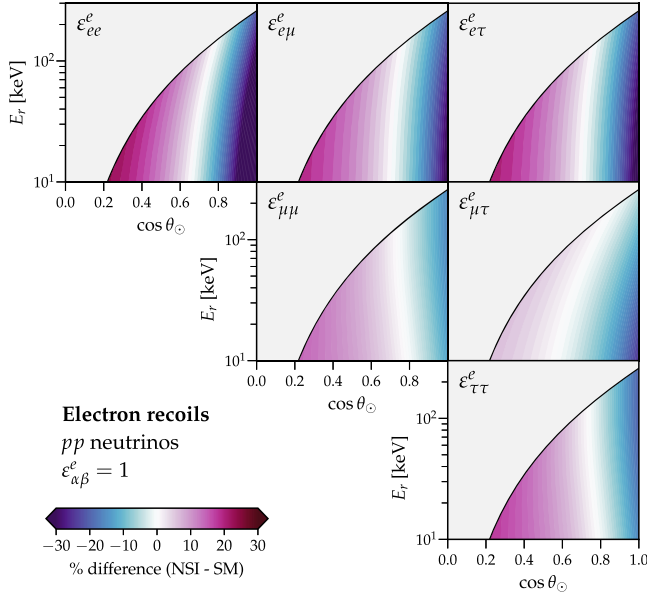


FIG. 3. Similar to Fig. 2, only now we show the relative difference in the energy-angle distributions of the recoils, which quantifies the differences compared to the SM for all six of the NSI parameters. More precisely, we are showing $(r_{\text{NSI}} - r_{\text{SM}})/r_{\text{SM}}$ where $r = (1/R)d^2R/dE_r d\cos\theta_\odot$ is the double-differential event rate normalized by the total rate. For each panel, we set all other $\epsilon_{\alpha\beta}^e$ to zero apart from the component labeled, which we set equal to 1. This figure highlights that ϵ_{ee}^e , $\epsilon_{e\mu}^e$, and $\epsilon_{e\tau}^e$ have the largest relative impact on the angular distributions.

readout). A detailed understanding of the expected ionization distributions from recoils of a given energy and particle type is then required to infer the initial energy and direction from the reconstructed ionization track. This procedure involves either fitting the distributions using traditional methods, or using dedicated machine-learning algorithms [88,114]. Since we are not performing full gas simulations here, our goal is to write down simple functional descriptions of the physical effects that limit the ability to reconstruct the initial recoil direction and energy perfectly, inspired by the performance of existing algorithms tested on experimental data. We do this following the same strategy as in Ref. [129], which we recap briefly here.

We emulate the performance of the detector by convolving the underlying theoretical event rate with several performance kernels. These describe the major effects that degrade the ability of the detector to reconstruct the correct energy and $\cos\theta$. These effects are especially important at low energies when the tracks become small and signal-to-noise decreases. The most efficient approach to apply these effects is at the event level, by running a Monte Carlo simulation. The procedure we follow is described further in Ref. [129], in short, it amounts to first generating recoil energies and angles from $d^2R/dE_r d\cos\theta_\odot$, and then resampling those energies and angles from a set of kernels that describe the energy and angular resolutions. The final

step is then to randomly assign a vector sense to the recoil direction according to a head-tail efficiency function. We describe these functions below. Then, when those effects are applied, we rebin the events into a new “observed” double-differential event rate, $d^2R_{\text{obs}}/dE_r d\cos\theta_\odot$.

The effect of finite energy resolution is captured by a Gaussian with width σ_E that depends on energy. Given a true underlying value of E_r , we draw a “measured” recoil energy E'_r from the distribution, as follows:

$$K_E(E_r, E'_r) = \frac{1}{\sqrt{2\pi}\sigma_E(E_r)} \exp\left(-\frac{(E_r - E'_r)^2}{2\sigma_E^2(E_r)}\right). \quad (37)$$

We use the following simple function for the energy spread,

$$\frac{\sigma_E}{E_r} [\%] = 11 + \frac{21}{\sqrt{E_r/\text{keV}}}, \quad (38)$$

which is fit to electron recoil data from the CYGNO-LIME detector [144]. The reference energy of 35 keV is arbitrary, but here we choose it to lie at the 10th percentile of the recoil energy distribution above the threshold of 10 keV, which means we can interpret the value of p_E as the resolution achieved by 90% of the recoil spectrum.

Next, the reconstruction of $\cos\theta$ is degraded by an angular resolution kernel taken to be a von Mises–Fisher distribution [114],

$$K_\theta(\hat{\mathbf{q}}, \hat{\mathbf{q}}', E_r) = \frac{\kappa(E_r)}{4\pi \sinh \kappa(E_r)} \exp(\kappa(E_r) \hat{\mathbf{q}} \cdot \hat{\mathbf{q}}'), \quad (39)$$

where the parameter that specifies the variance of the von Mises–Fisher distribution is κ . Rather than quoting this parameter, we will express the angular resolution in terms of the more intuitive variable, the rms axial angular resolution σ_θ , which is a monotonic function of κ that can be evaluated numerically. We find this relationship by calculating the rms axial angle between a sample of recoil vectors generated according to the distribution above, and then repeating this process over a suitable range of κ . The von Mises–Fisher distribution has the advantage over, for example, a truncated Gaussian, in that it reflects the fact that the distribution exists on the sphere.

The rms angle between two axes approaches $\sigma_\theta = 1$ radian for completely uncorrelated directions (i.e., the limit of no directional sensitivity). Following on from the discussion of $\sigma_E(E_r)$ above, we approximate the energy-dependence of the angular resolution with the following function form,

$$\sigma_\theta(E_r) = p_\theta \sqrt{\frac{E_{\text{ref}}}{E_r}}, \quad (40)$$

where $p_\theta = 10^\circ$ for our baseline detector performance benchmark, as inspired by results from gas simulations

of electron recoils at the relevant energies [144,145]. Note that the E_r appearing in this equation is the true recoil energy, not the reconstructed energy, E'_r .

One important assumption that we should emphasize here is that, by taking $\cos \theta_\odot$ as our derived measurable quantity, we are assuming that the detector has three-dimensional reconstruction, and so the track direction can be rotated into heliocentric coordinates. Although we do this for simplicity (relaxing this assumption would force us to also consider the event time and detector orientation), the ability to perform three-dimensional direction reconstruction is considered to be a major design feature in all MPGD-based detectors being developed right now. CYGNO is able to achieve this by first measuring the two-dimensional projection of the recoil track along the optical readout plane made up of highly sensitive complementary metal–oxide–semiconductor cameras, which are placed beyond an amplification stage made up of GEMs. This projection is then combined with the one-dimensional projection along the drift axis provided through the timing of scintillation light detected by photomultiplier tubes. The combination of these signals is what also enables the precise reconstruction of the recoil energy.

The final detector performance effect we account for is the limitation in reconstructing the correct $\pm$ sense of the recoil direction. This is a particularly important effect at the lowest energies, where tracks are shortest and the ionization density is low. It is possible at these energies that the track alignments could be reconstructed even while the head and the tail of the track are not. The head and tail of a recoil are primarily reconstructed by using the differences in the dE/dx along the track. Unlike nuclear recoils, which are on the stopping side of the Bragg peak, the electrons deposit *more* ionization as they slow down. So the tail (end) of the track curls up and contains the most ionization. Since in our simplified description, head-tail assignment is binary, we simply randomly assign a $\pm$ sign to each recoil vector with a probability given by the function

$$\varepsilon_{\text{HT}}(E_r) = \frac{1}{2} \left(\frac{1}{1 + e^{-a_{\text{HT}}(E_r - E_{\text{HT}})}} + 1 \right). \quad (41)$$

This functional form captures the behavior of head-tail performance well for all recoils, which, by definition, must begin at zero at $E_r = 0$, and plateau at 1 for high-enough energies. For our baseline performance we assume $a_{\text{HT}} = 0.2$ and $E_{\text{HT}} = 20$ keV. Since the neutrino recoils are restricted to a single hemisphere due to the scattering kinematics (up to the effects of angular resolution), the primary effect of the head-tail performance is to limit the ability to distinguish neutrino events from isotropic backgrounds.

The final effect we should include in principle is the event-level efficiency of the experiment. This is more difficult to quantify at present, as it will depend on the specifics of a given detector and the way the data are collected and analyzed. To keep the discussion in this study

as clear as possible for when our results are applied in real detector simulations by the CYGNO Collaboration, we simply incorporate a step-function efficiency curve by imposing a hard threshold at 10 keV. Our results are relatively weakly dependent on this threshold since the electron recoil distribution due to pp and ${}^7\text{Be}$ neutrinos is not very steeply falling at these energies, and in any case, all of our energy and angular performance curves have already degraded significantly by this energy.

B. Backgrounds

The rate of non-neutrino electron-recoil backgrounds in experiments of this type is the main quantity that will limit their sensitivity. Fortunately, the directionality will provide an excellent technique to reject backgrounds, which are expected to be close to isotropic. Nevertheless, backgrounds will appear in the signal region defined by $\cos \theta_\odot \in [0, 1]$, so we must incorporate some background model that quantifies this.

We make an educated guess as to the rate and energy spectrum of the background using recent background simulations performed in preparation for a future 30 m³ scale-up of the CYGNO experiment to be located at Laboratori Nazionali del Gran Sasso. Note that the general design parameters of CYGNO-30, in terms of gas mixture and directional performance, have already inspired most of the choices we have made in this study so far. The differential recoil spectrum for an isotropic background model can be written as follows:

$$\frac{d^2 R_{\text{bg}}}{dE_r d\Omega_r} = \frac{R_{\text{bg}}}{4\pi E_{\text{bg}} (e^{-E_{\text{th}}/E_{\text{bg}}} - e^{-E_{\text{max}}/E_{\text{bg}}})} e^{-E_r/E_{\text{bg}}}. \quad (42)$$

For our chosen model we have $E_{\text{th}} = 10$ keV and $E_{\text{max}} = 700$ keV. The model is parametrized by one free parameter, the total background rate R_{bg} . We fix the scale of the exponential fall off of the background energy spectrum to be $E_{\text{bg}} = 67$ keV, which we find to be a good fit to the spectrum of the total internal background from a simulation of CYGNO-30 [144]. This simulation includes all material electron backgrounds in the 0–1000 keV energy window, with the field cage being the dominant contributor to the background budget, followed by the GEMs and camera sensors.

As for the amplitude of the background rate, R_{bg} , to make the discussion more intuitive, we will express its value relative to the total event rate from solar neutrinos R_ν under the SM (see the gray dashed line in Fig. 1). Importantly, the simulations for CYGNO [144] suggest $R_{\text{bg}}/R_\nu \sim 500$ is *already* achievable without any dedicated background reduction which, as we will discuss, is quite promising. In particular, this study did not account for the possibility of using the ultra-low radioactivity kapton presented in Ref. [146], and so naively a further

background reduction is possible already. To see the impact of the size of the background on the scientific reach of the experiment we will explore a wide range of possible background reduction factors, while keeping in mind that $R_{\text{bg}} = 50R_\nu$ is a realistic expectation for a fully deployed experiment of 30 m³ volume. To demonstrate the power of directionality, in our likelihood analysis described in the next section, we will float the value of R_{bg} as an unknown nuisance parameter.

IV. LIKELIHOOD ANALYSIS

Having described all the necessary ingredients for calculating a reasonable approximation to the neutrino signal detectable by a gas-based directional detector, the final step is to develop a statistical methodology. Our previous analysis in Ref. [129] focused on the ability of the detector to measure the neutrino flux over some experimental background. Here, our scientific question is slightly different. Given an observed set of neutrino-induced electron recoil energies and directions, is their distribution consistent with the SM or not? Fortunately, the SM event rate can be expressed as a special case of the more general NSI formalism (i.e., $\forall \epsilon_{\alpha\beta}^e = 0$, or $g_{Z'} = 0$) which means our null hypothesis that the SM is true is nested within a more general alternative hypothesis that the event rate is described by Eq. (35). This then permits us to use the profile likelihood ratio test and apply Wilks' theorem.

We start by writing down a likelihood. To facilitate the use of Asimov statistics (see below), we will use a binned likelihood based on the event rate having been binned in energy and $\cos\theta_\odot$, i.e.,

$$N_\nu^{ij}(\epsilon) = MT \int_{\cos\theta_\odot^i}^{\cos\theta_\odot^i + \Delta\cos\theta} \int_{E_r^j}^{E_r^j + \Delta E_r} \frac{d^2 R_{\text{obs}}(\epsilon)}{dE_r d\cos\theta_\odot} dE_r d\cos\theta_\odot, \quad (43)$$

where MT is the exposure (mass $\times$ duration) of the experiment. Our bins run over the domain $\cos\theta_\odot \in [-1, 1]$ and $E_r \in [10, 700]$ keV across 50 linearly spaced bins each. The exact binning scheme does not impact our results since the finite energy and angle reconstruction are already taken care of, as discussed in Sec. III. For compactness, we use the shorthand here $\epsilon \equiv \{\epsilon_{\alpha\beta}^e\}$ to denote the set of all ϵ values upon which we are performing the statistical test. The background distribution is constructed using the same binning scheme. We parametrize the background using a normalization b whose correct value is 1. This means we can parametrize an expected number of events in bin ij as

$$N_{\text{exp}}^{ij}(\epsilon, b) = N_\nu^{ij}(\epsilon) + bN_{\text{bg}}^{ij}. \quad (44)$$

We now construct a binned likelihood for the parameters $\{\epsilon, b\}$ given some observed data N_{obs}^{ij} which we assume to

be Poisson distributed with mean N_{exp}^{ij} , as follows:

$$\mathcal{L}(\epsilon, b) = \prod_i \prod_j \text{Pois}[N_{\text{obs}}^{ij}; N_{\text{exp}}^{ij}(\epsilon, b)]. \quad (45)$$

In practice, we use the log-likelihood, which can be written as

$$\ln \mathcal{L}(\epsilon, b) = \sum_{ij} \left(N_{\text{obs}}^{ij} \ln[N_\nu^{ij}(\epsilon) + bN_{\text{bg}}^{ij}] - (N_\nu^{ij}(\epsilon) + bN_{\text{bg}}^{ij}) \right). \quad (46)$$

We neglect the $N_{\text{obs}}!$ term which will cancel when we take likelihood ratios.

We wish to construct a profile-likelihood ratio test that quantifies the extent to which the null hypothesis of the SM, i.e., $\epsilon = 0$, describes the data. To do this, we will assume that the observed data *are* described by the SM and forecast the expected exclusion limits one could then set on ϵ . We write down the standard profile likelihood ratio test statistic (TS) used for deriving exclusion limits as [147]

$$\text{TS}(\epsilon) = -2(\ln \mathcal{L}(\epsilon, \hat{\hat{b}}) - \ln \mathcal{L}(\hat{\epsilon}, \hat{b})),$$

where $\hat{\hat{b}}$ is the maximum likelihood estimator (MLE) at some fixed value of ϵ , whereas $\hat{b}$ and $\hat{\epsilon}$ are the MLEs for the full likelihood. Wilks' theorem holds in this model setup, which states that the distribution of TS asymptotically follows a χ_n^2 distribution where n is the number of parameters in ϵ .

To make an estimate of the expected exclusion limit that a future experiment could set, we employ the Asimov dataset. The Asimov dataset is a hypothetical situation in which the observed data exactly follow the expectation under some assumed set of "true" parameter values. Since we are projecting exclusion limits, where we imagine that the SM hypothesis correctly describes the data, we set $N_{\text{obs}} = N_{\text{exp}}(\forall \epsilon = 0, b = 1)$. The MLEs are then $\hat{b} = 1$ and $\hat{\epsilon} = 0$, while $\hat{\hat{b}}$ can be found numerically by solving

$$\frac{\partial \ln \mathcal{L}(\epsilon, b)}{\partial b} = \sum_{ij} \left(\frac{N_{\text{obs}}^{ij} N_{\text{bg}}^{ij}}{N_\nu^{ij}(\epsilon) + bN_{\text{bg}}^{ij}} - N_{\text{bg}}^{ij} \right) = 0. \quad (47)$$

As described in Ref. [147], the value of TS under the Asimov dataset asymptotically approaches the median of its distribution. That is to say, when we forecast an exclusion limit using the Asimov dataset, we find a good estimate of the level of sensitivity that at least 50% of hypothetical experiments could achieve. To forecast a limit at a given CL we then simply need to find values of ϵ at which the TS crosses some threshold value. For setting two-sided limits at the $1 - p$ level on the n parameters in ϵ we use the $1 - p$ percentile function of the χ_n^2 distribution.

For one-sided limits, e.g., on the parameter g_V in the light mediator case, we use the $1 - 2p$ percentile function.

V. RESULTS

We will now show estimated median 90% exclusion limits on the various $\epsilon_{\alpha\beta}^e$ and $(g_{Z'}, m_{Z'})$ parameters for a range of choices for the directional performance and background conditions. Since we are motivated by the upcoming 30 m³ detector being developed by the CYGNO Collaboration, we will present results for this fixed total volume, and provide an analytic scaling to understand how to extrapolate our results to larger volumes. Our primary goal is to motivate small-scale direction-sensitive gas detectors in general, as well as to understand the level of electron recoil background that is required to set relevant constraints on the neutrino cross section using a detector of this type.

To have something to compare our results against, it is useful to define what we call an equivalent “nondirectional experiment.” This is a hypothetical experiment that observes an identical event rate as our directional experiment, but does not have access to any $\cos \theta_\odot$ information. We can forecast limits for this experiment in the same way, but instead of N_{exp}^{ij} appearing in the likelihood, we first take the sum over $\cos \theta_\odot$ bins in the equation above, i.e.,

$$N_{\text{exp,non-dir}}^j(\epsilon) = \sum_i N_{\text{exp}}^{ij}(\epsilon). \quad (48)$$

As another point of comparison, we also define an “idealized performance” benchmark where the energy

and angular performances are not included. This amounts to calculating N_{exp} directly from the underlying $d^2R/dE_R d\cos \theta_\odot$ as opposed to $d^2R_{\text{obs}}/dE_R d\cos \theta_\odot$ which includes the performance effects. All of the figures in this section will then show three benchmark experiments: nondirectional, directional (realistic) and directional (idealized)—with the expected sensitivity improving in that order. We reiterate at this point that our goal is not to claim world-leading sensitivity to NSIs, but to highlight that—thanks to directionality—a gas-based experiment can make an interesting contribution to the global picture of the neutrino’s properties, despite being several orders of magnitude smaller and cheaper than its nearest competitors, e.g., Borexino.

First, in Fig. 4 we show a set of forecasted two-sided median 90% CL exclusion limits ($\text{TS} = 2.71$) on various $\epsilon_{\alpha\beta}^e$ parameters in the heavy mediator limit. These are one-dimensional constraints where we turn on one NSI coupling from $\alpha\beta = \{ee, e\mu, e\tau, \mu\mu, \mu\tau, \tau\tau\}$, with all others set to zero. As a direct point of comparison, we show the two-sided 90% CL exclusion limits from an analysis of Borexino Phase II spectral data [72]. In all cases, the gas mixture, exposure time, volume, and pressure are fixed as labeled in the figure. The three panels from left to right vary only the background rate, R_{bg} , from 5 to 500 times the SM neutrino rate. The latter value is the current estimate of the background level from simulations, prior to any dedicated screening or reduction efforts. We see that the constraints for the worst background condition are not at a competitive level, but after only a $\times 10$ reduction, the constraints for certain couplings like $\epsilon_{e\tau}^e$ reach within a

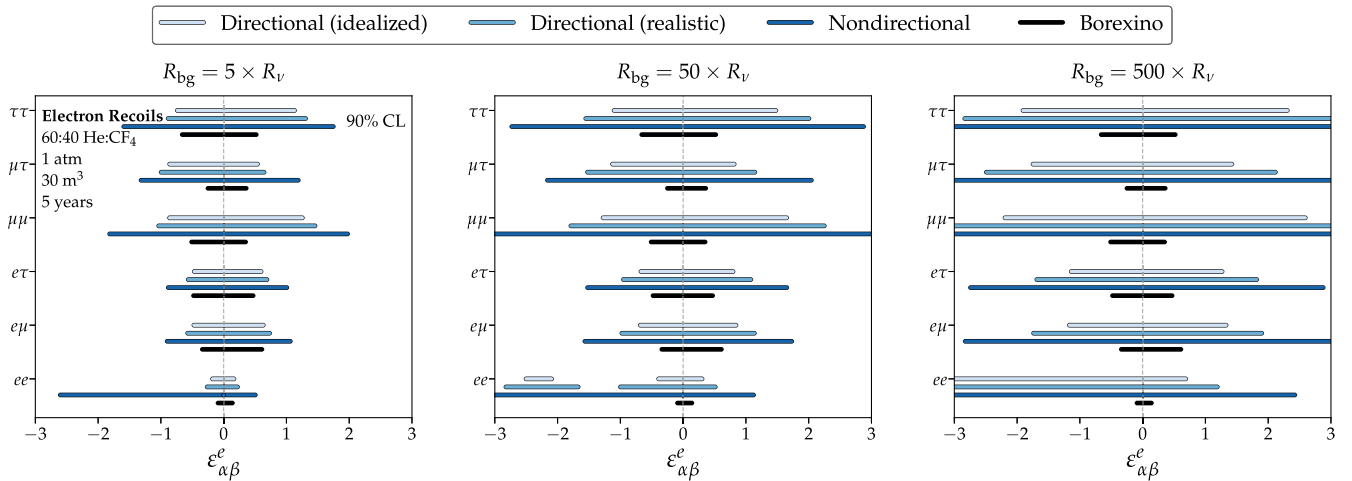


FIG. 4. Projected median 90% CL exclusion limits on each $\epsilon_{\alpha\beta}^e$ parameter (assuming all other $\epsilon_{\alpha\beta}^e = 0$). We assume a detector of 30 m³ volume with a 60:40 He:CF₄ gas mixture at atmospheric pressure. The three panels from left to right are for different background rates, measured relative to the total rate of neutrino-induced electron recoils under the SM, R_ν . The various projected limits shown in blue are for different levels of directional sensitivity: idealized performance, realistic performance, and no directionality, from lightest to darkest. We emphasize that the “nondirectional” case has an identical event rate, energy threshold and energy resolution as the “directional” cases. For comparison, in black we show the existing constraints from Borexino, demonstrating a competitive level of sensitivity despite the small size of the proposed experiment.

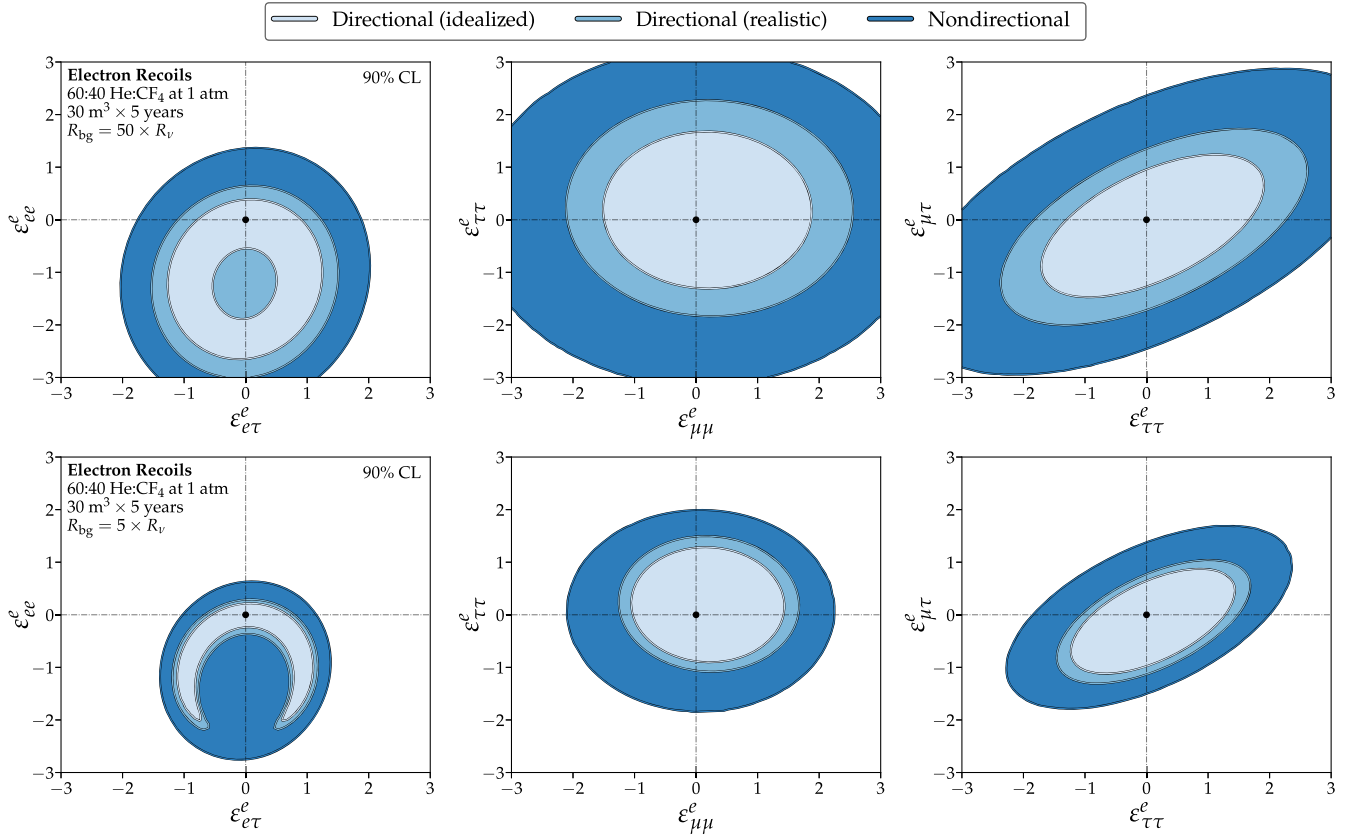


FIG. 5. Projected median 90% CL exclusion limits when two $\epsilon_{\alpha\beta}^e$ parameters are considered simultaneously (assuming all other $\epsilon_{\alpha\beta}^e = 0$). We assume a detector of 30 m^3 volume with a 60:40 He:CF₄ gas mixture at atmospheric pressure. The three columns correspond to three different pairs of NSI parameters. The top set of panels is for a background rate that is 50 times the neutrino electron recoil rate under the SM, while the bottom set of panels is for five times the neutrino rate. The various projected limits shown in blue are for different levels of directional sensitivity: idealized performance, realistic performance, and no directionality, from lightest to darkest. We emphasize that the “nondirectional” case has an identical event rate, energy threshold and energy resolution as the “directional” cases.

modest factor of Borexino’s. For an even more aggressive reduction of a factor of 100, the exclusion limits almost match those of Borexino. The case for directionality in particular is made the most apparent for ϵ_{ee}^e where the added information provided through the $\cos \theta_\odot$ distribution as well as the background rejection power allows for the region around $\epsilon_{ee}^e \approx -2$ to be excluded—this point in the parameter space is mildly degenerate with the SM because it leads to the same total number of events. Across all cases, we observe that the exclusion limit on ϵ scales as

$$\epsilon \propto (MT)^{-1/4} R_{\text{bg}}^{1/4}. \quad (49)$$

It is also possible for multiple NSI couplings to be active at once and to interplay with each other at the level of the event rate. To capture this possibility, we show three more examples in Fig. 5. Now we are showing two-dimensional 90% CL exclusion limits on 2 degrees of freedom (TS = 4.61), where two $\epsilon_{\alpha\beta}^e$ are turned on, with all others fixed at zero. The top and bottom sets of panels are again for two background assumptions: $50\times$ and $5\times$ the neutrino

rate, respectively. Again, we see from comparing the darkest blue contours (nondirectional) with the two lighter contours that the power of directionality is apparent. It is clearest again in the ee case, where the degeneracy with the SM for large negative values of ϵ_{ee}^e is broken.

Then finally, we come to the light mediator example in Fig. 6. As discussed in Sec. II we are showing only the $B-L$ model here with $\forall Q_{Z'}^{e,\nu\alpha} = -1$, but have checked that we obtain almost identical results for other models. We adopt a slightly different definition for the exclusion limits here and following the typical convention—here we have forecasted the median *one-sided* 90% CL exclusion limits on one parameter, $g_{Z'}$, with the two-dimensional constraint then built up by repeating the test over a range of values for $m_{Z'}$. The Borexino constraint is again from Ref. [72].⁴ We note that there are various other constraints from stellar

⁴We note that the Borexino constraint shown here is actually a two-sided 90% CL exclusion on $g_{Z'}$ rather than the one-sided limit on $|g_{Z'}|$ as we have shown. The differences between the two definitions are almost invisible on the plot at the scale we have shown it, so we choose not to debate this choice.

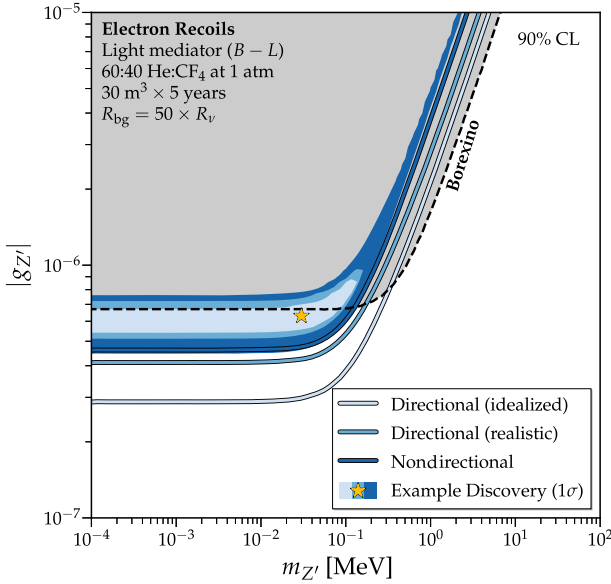


FIG. 6. Projected constraints on a light vector mediator under the $U(1)_{B-L}$ model, as a function of the coupling $g_{Z'}$ and mediator mass $m_{Z'}$. The solid lines show median 90% CL exclusion limits for the three experimental benchmarks used in the previous figures. The constraint from Borexino from Ref. [72] is shown for comparison. The shaded regions show example 1σ discovery contours if the true values of the parameters were situated at the gold star. We emphasize that the directional information (and the background rejection that comes with it) is what would allow such a discovery claim to be made. Following conventions, the shaded region encloses the *excluded* parameter space, as opposed to the previous figures, where the shaded regions were the *allowed* parameter space.

cooling, fixed target experiments, and BBN/ N_{eff} (big bang nucleosynthesis) that are also present for these light mediator models that partially overlap this region. We choose to remove these for the sake of clutter and because our main message is the comparison between these two solar neutrino experiments. We note that it is always possible to model build around any of these constraints, and not all are present for every instance of an anomaly-free gauged $U(1)'$ model.

As expected, the sensitivity shown in Fig. 6 is best at low mediator masses where the deviation in the cross section due to NSIs emerges at energies that are below Borexino's threshold. Borexino, of course, still leads at higher mediator masses. We show the case here for $R_{\text{bg}} = 50R_{\nu}$, but sensitivity scales slowly $g_{Z'} \propto (MT)^{-1/8} R_{\text{bg}}^{1/8}$, so the results for $5R_{\nu}$ are similar.

For this figure, in addition to the forecasted exclusion limits, we show another way to demonstrate the power of directionality by forecasting a possible discovery. The ability to make discoveries is another of the major motivations behind directional detection that has been appreciated for many years (see e.g., the discussion in the context of dark matter in Ref. [88]). Although experiments like

XENONnT, LZ and PandaX have the statistics to set *constraints* at a similar level to what is shown here in principle, in practice, it is challenging for them to make positive claims of a signal due to the high electron background preventing the robust identification of neutrino events. This is not the case for a directional experiment, where a much stronger statement about the probability of any given recoil as being due to a neutrino is possible, thanks to the angular domain being split into an (almost) pure background hemisphere and a signal + background hemisphere. To showcase this we show on Fig. 6 some example median 1σ discovery contours, given some arbitrarily chosen hypothetical “true” value for $(g_{Z'}, m_{Z'}) = (g_{Z'}^{\text{true}}, m_{Z'}^{\text{true}})$ shown by the gold star. We calculate this contour using a “discovery” test statistic, which essentially amounts (in the Asimov procedure) to fixing $N_{\text{obs}} = N_{\nu}(g_{Z'}^{\text{true}}, m_{Z'}^{\text{true}}) + N_{\text{bg}}$, whereas previously for exclusion limits we set $g_{Z'}^{\text{true}} = 0$. Comparing the resulting contours between the directional and nondirectional examples, we again see the benefit that comes with directionality.

VI. CONCLUSIONS

In this paper, we have motivated the continued development of 30 m³-scale gaseous detectors for the directional detection of solar neutrinos through electron recoils. In alignment with the message of Ref. [129], we argue that an experiment of this type can perform neutrino physics, trading in cost and overall size for the added signal information and background rejection that comes from directional sensitivity. We illustrated this in the context of the neutrino NSI formalism, where the interaction between neutrinos and electrons is modified by the presence of a new vector Z' mediator beyond the SM.

Our results are illustrated in Figs. 4 and 5 in the heavy mediator limit ($m_{Z'} \gg q$), where we permit both flavor-conserving and flavor-violating interactions. Then, in Fig. 6, we showed the case for a light mediator, where we allow for only flavor-preserving interactions, as expected in the gauged $U(1)_{B-L}$ model. We have not attempted an exhaustive study of all possible combinations of NSI flavor structures, alternative mediator types, or other $U(1)'$ models, as the results presented here are sufficient to convey our main message. We believe a background reduction of around a factor of ten is needed from preliminary detector background simulations to reach an interesting level of sensitivity to neutrino interactions beyond the SM using solar neutrinos.

Our results remain to be validated in a full Monte Carlo simulation of a detector like this. This work, as well as the background reduction efforts that we advocate for here, is ongoing within the CYGNO Collaboration [144]. Our results suggest a promising future for this technology. The next stage of CYGNO and the detectors of other groups within the CYGNUS Consortium [103] will be the

first in a sequence of gas-based directional detectors to bring a range of complementary capabilities to the fleet of worldwide neutrino observatories.

ACKNOWLEDGMENTS

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

APPENDIX A: EFFECTIVE NEUTRINO OSCILLATIONS

In Sec. II C 1, when describing the computation of the density matrix in the context of NSIs, we referred to the effective description of the neutrino oscillations in the Sun within the OMD approximation, as well as the averaging over the unknown neutrino production location. In this appendix, we give more details of this calculation.

Within the OMD approximation, the Hamiltonian in Eq. (22) can be rotated by $O \equiv R_{23}R_{13}$, such that the rotated Hamiltonian $\tilde{H} = O^\dagger H O$, and its evolution $\tilde{S}_{\text{prop}} = O^\dagger S_{\text{prop}} O$, is given by

$$\tilde{H} = \begin{pmatrix} H^{(2)} & 0 \\ 0 & \Delta_{33} \end{pmatrix}, \quad \tilde{S}_{\text{prop}} = \begin{pmatrix} S^{(2)} & 0 \\ 0 & e^{-i\Delta_{33}L} \end{pmatrix}. \quad (\text{A1})$$

Here, $H^{(2)}$ is the effective 2×2 Hamiltonian and $S^{(2)} \equiv \text{Evol}[H^{(2)}]$, describes the evolution in an effective two-flavor scenario. Following the conventions in Refs. [46,48,72,148] we get

$$H^{(2)} = \frac{\Delta m_{21}^2}{4E_\nu} \begin{pmatrix} -\cos 2\theta_{12} & \sin 2\theta_{12} e^{i\delta_{\text{CP}}} \\ \sin 2\theta_{12} e^{-i\delta_{\text{CP}}} & \cos 2\theta_{12} \end{pmatrix} + \sqrt{2}G_F N_e(x) \begin{pmatrix} c_{13}^2 - \epsilon_D & \epsilon_N \\ \epsilon_N & \epsilon_D \end{pmatrix}, \quad (\text{A2})$$

where the terms ϵ_D and ϵ_N are related to the combined NSI parameter $\epsilon_{\alpha\beta}$ defined in Eq. (23) by

$$\begin{aligned} \epsilon_D &= c_{13}s_{13}(s_{23}\epsilon_{e\mu} + c_{23}\epsilon_{e\tau}) - (1 + s_{13}^2)c_{23}s_{23}\epsilon_{\mu\tau} \\ &\quad - \frac{1}{2}c_{13}^2(\epsilon_{ee} - \epsilon_{\mu\mu}) + \frac{1}{2}(s_{23}^2 - (s_{13}c_{23})^2)(\epsilon_{\tau\tau} - \epsilon_{\mu\mu}), \\ \epsilon_N &= c_{13}(c_{23}\epsilon_{e\mu} - s_{23}\epsilon_{e\tau}) + s_{13}[(s_{23}^2 - c_{23}^2)\epsilon_{\mu\tau} \\ &\quad + c_{23}s_{23}(\epsilon_{\tau\tau} - \epsilon_{\mu\mu})]. \end{aligned} \quad (\text{A3})$$

We then define the effective matter potential A_{CC} and the matter-modified mixing p and q , as follows:

$$A_{\text{CC}} = 2\sqrt{2}G_F N_e(x)E_\nu, \quad (\text{A4})$$

$$p = \left| \sin(2\theta_{12})e^{i\delta_{\text{CP}}} + \frac{2A_{\text{CC}}\epsilon_n}{\Delta m_{21}^2} \right|, \quad (\text{A5})$$

$$q = \cos(2\theta_{12}) - \frac{2A_{\text{CC}}}{\Delta m_{21}^2} \left(\frac{c_{13}^2}{2} - \epsilon_d \right), \quad (\text{A6})$$

to obtain the following relations for the mass-squared difference and mixing angle in matter:

$$\Delta m_{21,m}^2 = \Delta m_{21}^2 \sqrt{p^2 + q^2}, \quad (\text{A7})$$

$$\sin(2\theta_m) = \frac{p}{\sqrt{p^2 + q^2}}, \quad (\text{A8})$$

$$\cos(2\theta_m) = \frac{q}{\sqrt{p^2 + q^2}}. \quad (\text{A9})$$

Taking the spatial average of $\cos(2\theta_m)$ over the different production zones within the Sun, weighted by the fraction of neutrinos produced in each zone $f(x)$, we get

$$\langle \cos(2\theta_m) \rangle = \int_0^1 \cos(2\theta_m(x))f(x)dx, \quad (\text{A10})$$

where $x = r/R_\odot$ is the normalized solar radial coordinate. The production site distributions, $f(x)$, are different for each flux, and are taken from Ref. [149].

APPENDIX B: COHERENT ELASTIC NEUTRINO-NUCLEUS SCATTERING

1. Standard model

Given a nucleus N of mass m_N , under the SM the $\text{CE}\nu\text{NS}$ process $\nu N \rightarrow \nu N$ occurs only through neutral current interactions via the Z exchange, which couples to all neutrino flavors equally. As a result, this process is flavor-blind at tree-level, and the cross section is given by

$$\left(\frac{d\sigma}{dE_r} \right)^{\text{SM}} = \frac{G_F^2 m_N F^2(q) g_V^2}{\pi} \left(1 - \frac{m_N E_r}{2E_\nu^2} \right). \quad (\text{B1})$$

For a given nucleus with atomic number Z , mass number A and the number of neutrons $N = A - Z$, the SM coupling $g_V = Zg_p^V + Ng_n^V$. Here, $g_p^V = 2g_u^V + g_d^V = \frac{1}{2} - 2\sin^2\theta_w$ and $g_n^V = g_u^V + 2g_d^V = -\frac{1}{2}$ which signify the couplings with protons and neutrons respectively, in terms of the u and d couplings. The form factor, $F(q)$, is used to capture the effects of nuclear structure, which become important at large momentum transfer, $q = \sqrt{2m_N E_r}$. The commonly used Helm ansatz for the form factor is [150]

$$F(q) = \frac{3j_1(qr_n)}{qr_n} \exp\left(\frac{-q^2 s^2}{2}\right), \quad (\text{B2})$$

where $j_1(q)$ is the spherical Bessel function of first kind, $r_n = \sqrt{\frac{5}{3}(R_{\min}^2 - 3s^2)}$ where $s = 0.9$ fm is a measure of the nuclear skin thickness and $R_{\min}$ is the rms radius of the proton distribution of the nucleus. The form factor plays an essentially negligible role for recoil energies and nuclear targets we are interested in here, so we do not require any more sophisticated models than this—see e.g., Ref. [151].

2. NSIs for neutrino-nucleus scattering

Considering again the pure vector NSI case, we now take the Lagrangian in Eq. (13) and set $f = \{u, d\}$ to build the modified CE ν NS process. Assuming that the new light mediator has the same coupling with both u and d quarks, the new interactions can be accounted for in the cross section in Eq. (B1) by modifying

$$g_V^{\text{NSI}} = g_V^{\text{SM}} + \frac{3g_{Z'}^f Q_{Z'}^{\nu_q}(Z + N)}{\sqrt{2}G_F(2m_N E_r + m_{Z'}^2)}. \quad (\text{B3})$$

For a heavy mediator, $m_{Z'} \gg q$, we consider the Lagrangian in Eq. (14) with $f = \{u, d\}$. Following the same procedure as described in the main text in the context of E ν eS, the generalized cross section for CE ν NS can be written as [46]

$$\left(\frac{d\zeta}{dE_r}\right)_{\alpha\beta}^{\text{NSI}} = \frac{G_F^2 m_N F^2(q)}{\pi} \left(1 - \frac{m_N E_r}{2E_\nu^2}\right) \times \left[g_V^2 \delta_{\alpha\beta} + 2g_V G_{\alpha\beta} + \sum_f G_{\alpha f} G_{f\beta}\right], \quad (\text{B4})$$

where

$$G_{\alpha\beta} = Z(2\varepsilon_{\alpha\beta}^u + \varepsilon_{\alpha\beta}^d) + N(\varepsilon_{\alpha\beta}^u + 2\varepsilon_{\alpha\beta}^d). \quad (\text{B5})$$

The double-differential event rate for CE ν NS can also be written in exactly the same way as for E ν eS—cf. Eq. (36)—up to replacing the electron mass m_e by the nucleus mass m_N , and using N_T to count the number of nuclei in the target.

3. Nuclear recoils in directional gas detectors

The directions and energies of nuclear recoils are also independently measurable in a gas detector through essentially the same procedure as with electron recoils. In the case of nuclear recoils, however, the sizes of the tracks at the relevant energies (1–15 keV $_r$ for the nuclei in our chosen gas mixture) are much shorter, at $\sim$ cm in length. This makes the effects of diffusion as the charge is transported to the readout plane a major obstacle for achieving good directionality. While directionality can still be used for background rejection, the near-disappearance of any details of the recoil tracks at low energies makes this challenging. This is coupled with the fact that we need to extend to much larger volumes to acquire sufficient statistics in the case of CE ν NS, since the only detectable flux, ^8B , is significantly weaker than that of pp neutrinos. This is why we demote this discussion to an appendix, but nevertheless, since this channel may be of interest in the longer term, we show some example results here.

Largely inspired by the extensive study of nuclear recoil directionality performed in the context of dark matter and solar CE ν NS in Ref. [103], we adopt an ambitious but not unrealistic energy threshold of 3 keV $_r$, taking all nuclear recoil energies to be those before quenching. We set the following performance parameters:

$$\left\{a_{\text{HT}}, \frac{E_{\text{HT}}}{\text{keV}}, \frac{p_\theta}{\text{deg}}, p_E\right\} = \{0.1, 10.0, 25, 0.06\}, \quad (\text{B6})$$

using the same functional forms for energy resolution, head-tail efficiency and angular resolution. Unlike in the electron recoil case, the relatively straight nature of the nuclear recoil tracks, combined with the washing-out of the ionization distribution due to the diffusion, makes head-tail much more challenging to measure.

In Fig. 7 we show forecasted 90% CL exclusion limits on the two-dimensional space of u and d -type couplings, i.e., $\varepsilon_{\alpha\alpha}^u, \varepsilon_{\alpha\alpha}^d$ across several possible flavor-conserving combinations. Results for other parameters are qualitatively similar. As can be readily observed from the plots, an ambitious volume of 1000 m 3 and an aggressive background reduction in the nuclear recoil channel down to 50 times the ^8B neutrino rate is required to obtain a level of sensitivity already met by XENONnT and PandaX via their detection of this solar flux. This conclusion is in general alignment with the sensitivity projections made in Ref. [103]. Although detecting solar CE ν NS using this kind of experiment may be somewhat impractical in the near term, we expect the situation to be much more promising for detecting CE ν NS with a higher-energy neutrino source, which we will explore in a future study.

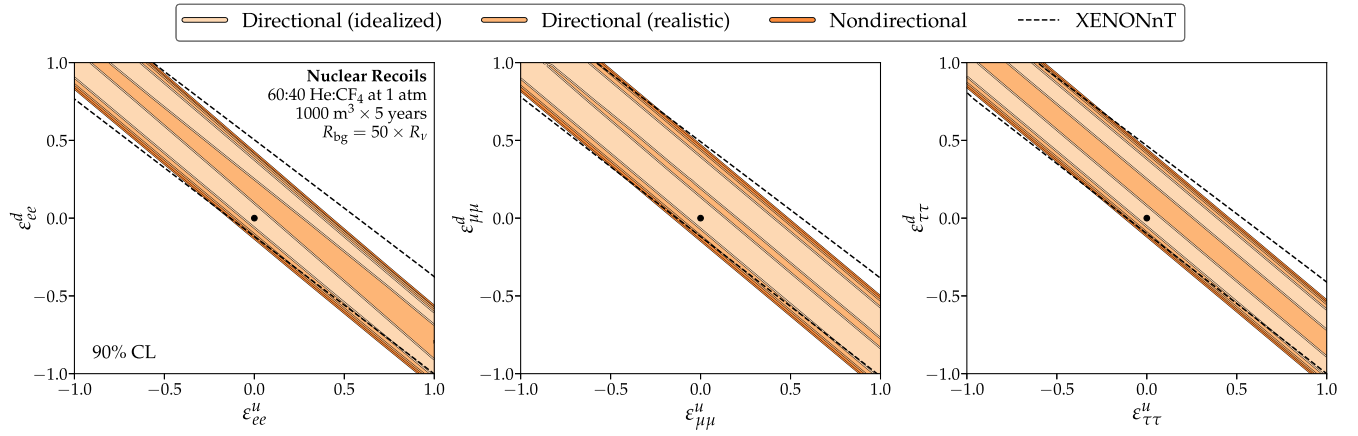


FIG. 7. Projected median 90% CL exclusion limits for a directional gas detector at the required scale needed to measure CE ν NS with solar neutrinos, i.e., a volume of 1000 m³ with atmospheric pressure gas. We assume a nuclear recoil threshold of 3 keV_r and performance as described in Appendix B 3. We project constraints onto the two-dimensional plane of up and down type NSI couplings across three example flavor-conserving interaction combinations, i.e., ee , $\mu\mu$ and $\tau\tau$. For comparison, we show the constraints from Ref. [48], for XENONnT's recent measurement of ^8B neutrinos [23]. This represents a further future and less immediate goal for a potential large-scale recoil observatory, as designed for sensitivity to dark matter [103].

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- [1] P. Huber *et al.*, Snowmass neutrino frontier report, in *Snowmass 2021* (2022), [arXiv:2211.08641](#).
 - [2] J.R. Klein *et al.*, SNOWMASS neutrino frontier NF10 topical group report: Neutrino detectors, [arXiv:2211.09669](#).
 - [3] K. Abe *et al.* (Hyper-Kamiokande Collaboration), Hyper-Kamiokande design report, [arXiv:1805.04163](#).
 - [4] Y. Fukuda *et al.* (Super-Kamiokande Collaboration), Evidence for oscillation of atmospheric neutrinos, *Phys. Rev. Lett.* **81**, 1562 (1998).
 - [5] B. Abi *et al.* (DUNE Collaboration), Deep underground neutrino experiment (DUNE), far detector technical design report, volume I introduction to DUNE, *J. Instrum.* **15**, T08008 (2020).
 - [6] V. Albanese *et al.* (SNO+ Collaboration), The SNO + experiment, *J. Instrum.* **16**, P08059 (2021).
 - [7] P. Adamson *et al.* (MINOS Collaboration), Measurement of the neutrino mass splitting and flavor mixing by MINOS, *Phys. Rev. Lett.* **106**, 181801 (2011).
 - [8] C. Arpesella *et al.* (Borexino Collaboration), Direct measurement of the Be-7 solar neutrino flux with 192 days of borexino data, *Phys. Rev. Lett.* **101**, 091302 (2008).
 - [9] M. G. Aartsen *et al.* (IceCube Collaboration), Evidence for high-energy extraterrestrial neutrinos at the IceCube detector, *Science* **342**, 1242856 (2013).
 - [10] K. Eguchi *et al.* (KamLAND Collaboration), First results from KamLAND: Evidence for reactor anti-neutrino disappearance, *Phys. Rev. Lett.* **90**, 021802 (2003).
 - [11] J. Álvarez-Muñiz *et al.* (GRAND Collaboration), The giant radio array for neutrino detection (GRAND): Science and design, *Sci. China Phys. Mech. Astron.* **63**, 219501 (2020).
 - [12] L. Baudis, A. Ferella, A. Kish, A. Manalaysay, T. Marrodan Undagoitia, and M. Schumann, Neutrino physics with multi-ton scale liquid xenon detectors, *J. Cosmol. Astropart. Phys.* **01** (2014) 044.
 - [13] L. Pattavina, N. Ferreiro Iachellini, and I. Tamborra, Neutrino observatory based on archaeological lead, *Phys. Rev. D* **102**, 063001 (2020).
 - [14] G. Agnolet *et al.* (MINER Collaboration), Background studies for the MINER coherent neutrino scattering reactor experiment, *Nucl. Instrum. Methods Phys. Res., Sect. A* **853**, 53 (2017).
 - [15] R. Strauss *et al.* (NUCLEUS Collaboration), Gram-scale cryogenic calorimeters for rare-event searches, *Phys. Rev. D* **96**, 022009 (2017).
 - [16] A. Aguilar-Arevalo *et al.* (CONNIE Collaboration), Exploring low-energy neutrino physics with the Coherent Neutrino Nucleus Interaction Experiment, *Phys. Rev. D* **100**, 092005 (2019).
 - [17] C. Augier *et al.* (Ricochet Collaboration), Ricochet progress and status, *J. Low Temp. Phys.* **212**, 127 (2023).
 - [18] S. Adrian-Martinez *et al.* (KM3Net Collaboration), Letter of intent for KM3NeT 2.0, *J. Phys. G* **43**, 084001 (2016).
 - [19] K. Abe *et al.* (T2K Collaboration), Measurements of neutrino oscillation parameters from the T2K experiment using 3.6×10^{21} protons on target, *Eur. Phys. J. C* **83**, 782 (2023).
 - [20] P. Abratenko *et al.* (MicroBooNE Collaboration), First double-differential measurement of kinematic imbalance

- in neutrino interactions with the MicroBooNE detector, *Phys. Rev. Lett.* **131**, 101802 (2023).
- [21] A. D. Avrorin *et al.* (Baikal-GVD Collaboration), Neutrino telescope in lake baikal: Present and future, *Proc. Sci., ICRC2019* (2020) 1011 [arXiv:1908.05427].
- [22] E. Vitagliano, I. Tamborra, and G. Raffelt, Grand unified neutrino spectrum at earth: Sources and spectral components, *Rev. Mod. Phys.* **92**, 045006 (2020).
- [23] E. Aprile *et al.* (XENON Collaboration), First indication of solar ^8B neutrinos via coherent elastic neutrino-nucleus scattering with XENONnT, *Phys. Rev. Lett.* **133**, 191002 (2024).
- [24] Z. Bo *et al.* (PandaX Collaboration), First indication of Solar ^8B neutrinos through coherent elastic neutrino-nucleus scattering in PandaX-4T, *Phys. Rev. Lett.* **133**, 191001 (2024).
- [25] D. Z. Freedman, Coherent neutrino nucleus scattering as a probe of the weak neutral current, *Phys. Rev. D* **9**, 1389 (1974).
- [26] M. Abdullah *et al.*, Coherent elastic neutrino-nucleus scattering: Terrestrial and astrophysical applications, arXiv:2203.07361.
- [27] M. Schumann, Direct detection of WIMP dark matter: Concepts and status, *J. Phys. G* **46**, 103003 (2019).
- [28] D. S. Akerib *et al.*, Snowmass2021 cosmic frontier dark matter direct detection to the neutrino fog, in *Snowmass 2021*, arXiv:2203.08084.
- [29] J. Cooley *et al.*, Report of the topical group on particle dark matter for snowmass 2021, arXiv:2209.07426.
- [30] D. Akimov *et al.* (COHERENT Collaboration), Observation of coherent elastic neutrino-nucleus scattering, *Science* **357**, 1123 (2017).
- [31] D. Akimov *et al.* (COHERENT Collaboration), First measurement of coherent elastic neutrino-nucleus scattering on argon, *Phys. Rev. Lett.* **126**, 012002 (2021).
- [32] D. Akimov *et al.* (COHERENT Collaboration), COHERENT Collaboration data release from the first detection of coherent elastic neutrino-nucleus scattering on argon, arXiv:2006.12659.
- [33] N. Ackermann *et al.*, First observation of reactor anti-neutrinos by coherent scattering, arXiv:2501.05206.
- [34] D. G. Cerdeño, M. Fairbairn, T. Jubb, P. A. N. Machado, A. C. Vincent, and C. Boehm, Physics from solar neutrinos in dark matter direct detection experiments, *J. High Energy Phys.* **05** (2016) 118; **09** (2016) 48.
- [35] B. Dutta, S. Liao, L. E. Strigari, and J. W. Walker, Non-standard interactions of solar neutrinos in dark matter experiments, *Phys. Lett. B* **773**, 242 (2017).
- [36] M. Abdullah, J. B. Dent, B. Dutta, G. L. Kane, S. Liao, and L. E. Strigari, Coherent elastic neutrino nucleus scattering as a probe of a Z' through kinetic and mass mixing effects, *Phys. Rev. D* **98**, 015005 (2018).
- [37] P. B. Denton, Y. Farzan, and I. M. Shoemaker, Testing large non-standard neutrino interactions with arbitrary mediator mass after COHERENT data, *J. High Energy Phys.* **07** (2018) 037.
- [38] D. W. P. Amaral, D. G. Cerdeno, P. Foldenauer, and E. Reid, Solar neutrino probes of the muon anomalous magnetic moment in the gauged $U(1)_{L_\mu-L_\tau}$, *J. High Energy Phys.* **12** (2020) 155.
- [39] O. G. Miranda, D. K. Papoulias, M. Tórtola, and J. W. F. Valle, Probing new neutral gauge bosons with CE ν NS and neutrino-electron scattering, *Phys. Rev. D* **101**, 073005 (2020).
- [40] L. M. G. de la Vega, L. J. Flores, N. Nath, and E. Peinado, Complementarity between dark matter direct searches and CE ν NS experiments in $U(1)'$ models, *J. High Energy Phys.* **09** (2021) 146.
- [41] A. Majumdar, D. K. Papoulias, and R. Srivastava, Dark matter detectors as a novel probe for light new physics, *Phys. Rev. D* **106**, 013001 (2022).
- [42] Y.-F. Li and S.-y. Xia, Constraining light mediators via detection of coherent elastic solar neutrino nucleus scattering, *Nucl. Phys.* **B977**, 115737 (2022).
- [43] T. Schwemberger and T.-T. Yu, Detecting beyond the standard model interactions of solar neutrinos in low-threshold dark matter detectors, *Phys. Rev. D* **106**, 015002 (2022).
- [44] M. Atzori Corona, M. Cadeddu, N. Cargioli, F. Dordei, C. Giunti, Y. F. Li, E. Picciau, C. A. Ternes, and Y. Y. Zhang, Probing light mediators and $(g-2)_\mu$ through detection of coherent elastic neutrino nucleus scattering at COHERENT, *J. High Energy Phys.* **05** (2022) 109.
- [45] M. Atzori Corona, W. M. Bonivento, M. Cadeddu, N. Cargioli, and F. Dordei, New constraint on neutrino magnetic moment and neutrino millicharge from LUX-ZEPLIN dark matter search results, *Phys. Rev. D* **107**, 053001 (2023).
- [46] D. W. P. Amaral, D. Cerdeno, A. Cheek, and P. Foldenauer, A direct detection view of the neutrino NSI landscape, *J. High Energy Phys.* **07** (2023) 071.
- [47] C. Giunti and C. A. Ternes, Testing neutrino electromagnetic properties at current and future dark matter experiments, *Phys. Rev. D* **108**, 095044 (2023).
- [48] D. Aristizabal Sierra, N. Mishra, and L. Strigari, Implications of first neutrino-induced nuclear recoil measurements in direct detection experiments: Probing nonstandard interaction via CE ν NS, *Phys. Rev. D* **111**, 055007 (2025).
- [49] V. De Romeri, D. K. Papoulias, and C. A. Ternes, Bounds on new neutrino interactions from the first CE ν NS data at direct detection experiments, *J. Cosmol. Astropart. Phys.* **05** (2025) 012.
- [50] V. De Romeri, D. K. Papoulias, G. Sanchez Garcia, C. A. Ternes, and M. Tórtola, Neutrino electromagnetic properties and sterile dipole portal in light of the first solar CE ν NS data, arXiv:2412.14991.
- [51] V. De Romeri, D. K. Papoulias, and G. Sanchez Garcia, Implications of the first CONUS+ measurement of coherent elastic neutrino-nucleus scattering, *Phys. Rev. D* **111**, 075025 (2025).
- [52] T. N. Maity and C. Boehm, First measurement of the weak mixing angle in direct detection experiments, arXiv:2409.04385.
- [53] P. Blanco-Mas, P. Coloma, G. Herrera, P. Huber, J. Kopp, I. M. Shoemaker, and Z. Tabrizi, Clarity through the neutrino fog: Constraining new forces in dark matter detectors, arXiv:2411.14206.
- [54] S. Xia, Measuring solar neutrino fluxes in direct detection experiments in the presence of light mediators, *Nucl. Phys.* **B1009**, 116738 (2024).

- [55] J. Aalbers *et al.* (DARWIN Collaboration), Solar neutrino detection sensitivity in DARWIN via electron scattering, *Eur. Phys. J. C* **80**, 1133 (2020).
- [56] A. de Gouvêa, E. McGinness, I. Martinez-Soler, and Y. F. Perez-Gonzalez, pp solar neutrinos at DARWIN, *Phys. Rev. D* **106**, 096017 (2022).
- [57] J. Aalbers *et al.*, A next-generation liquid xenon observatory for dark matter and neutrino physics, *J. Phys. G* **50**, 013001 (2023).
- [58] J. Aalbers *et al.* (XLZD Collaboration), The XLZD design book: Towards the next-generation liquid xenon observatory for dark matter and neutrino physics, [arXiv:2410.17137](https://arxiv.org/abs/2410.17137).
- [59] D. Aristizabal Sierra, V. De Romeri, L. J. Flores, and D. K. Papoulias, Light vector mediators facing XENON1T data, *Phys. Lett. B* **809**, 135681 (2020).
- [60] K. A. ShivaSankar, A. Majumdar, D. K. Papoulias, H. Prajapati, and R. Srivastava, Implications of first LZ and XENONnT results: A comparative study of neutrino properties and light mediators, *Phys. Lett. B* **839**, 137742 (2023).
- [61] A. N. Khan, Constraints on general light mediators from PandaX-II electron recoil data, *Phys. Lett. B* **819**, 136415 (2021).
- [62] V. De Romeri, D. K. Papoulias, and C. A. Ternes, Light vector mediators at direct detection experiments, *J. High Energy Phys.* **05** (2024) 165.
- [63] G. Bellini *et al.* (Borexino Collaboration), Measurement of the solar ^8B neutrino rate with a liquid scintillator target and 3 MeV energy threshold in the Borexino detector, *Phys. Rev. D* **82**, 033006 (2010).
- [64] G. Bellini *et al.* (Borexino Collaboration), Final results of Borexino Phase-I on low energy solar neutrino spectroscopy, *Phys. Rev. D* **89**, 112007 (2014).
- [65] G. Bellini *et al.* (BOREXINO Collaboration), Neutrinos from the primary proton–proton fusion process in the Sun, *Nature (London)* **512**, 383 (2014).
- [66] M. Agostini *et al.* (Borexino Collaboration), Improved measurement of ^8B solar neutrinos with 1.5 kt·y of Borexino exposure, *Phys. Rev. D* **101**, 062001 (2020).
- [67] M. Agostini *et al.* (Borexino Collaboration), First simultaneous precision spectroscopy of pp , ^7Be , and pep solar neutrinos with Borexino Phase-II, *Phys. Rev. D* **100**, 082004 (2019).
- [68] M. Agostini *et al.* (BOREXINO Collaboration), Comprehensive measurement of pp -chain solar neutrinos, *Nature (London)* **562**, 505 (2018).
- [69] M. Agostini *et al.* (BOREXINO Collaboration), Seasonal modulation of the ^7Be solar neutrino rate in Borexino, *Astropart. Phys.* **92**, 21 (2017).
- [70] M. Agostini *et al.* (BOREXINO Collaboration), Experimental evidence of neutrinos produced in the CNO fusion cycle in the Sun, *Nature (London)* **587**, 577 (2020).
- [71] A. N. Khan, W. Rodejohann, and X.-J. Xu, Borexino and general neutrino interactions, *Phys. Rev. D* **101**, 055047 (2020).
- [72] P. Coloma, M. C. Gonzalez-Garcia, M. Maltoni, J. P. Pinheiro, and S. Urrea, Constraining new physics with Borexino Phase-II spectral data, *J. High Energy Phys.* **07** (2022) 138; **11** (2022) 138(E).
- [73] M. Bauer, P. Foldenauer, and J. Jaeckel, Hunting all the hidden photons, *J. High Energy Phys.* **07** (2018) 094.
- [74] P. Ilten, Y. Soreq, M. Williams, and W. Xue, Serendipity in dark photon searches, *J. High Energy Phys.* **06** (2018) 004.
- [75] T. Han, J. Liao, H. Liu, and D. Marfatia, Nonstandard neutrino interactions at COHERENT, DUNE, T2HK and LHC, *J. High Energy Phys.* **11** (2019) 028.
- [76] B. Dutta, R. F. Lang, S. Liao, S. Sinha, L. Strigari, and A. Thompson, A global analysis strategy to resolve neutrino NSI degeneracies with scattering and oscillation data, *J. High Energy Phys.* **09** (2020) 106.
- [77] C. Antel *et al.*, Feebly-interacting particles: FIPs 2022 workshop report, *Eur. Phys. J. C* **83**, 1122 (2023).
- [78] J. Seguinot, T. Ypsilantis, and A. Zichichi, A High rate solar neutrino detector with energy determination, *Conf. Proc. C* **920310**, 289 (1992).
- [79] F. Arzarello, F. Ciralli, F. Frascioni, G. Bonvicini, G. Laurenti, A. Zichichi, J. Seguinot, T. Ypsilantis, and S. Tzamarias, HELLAZ: A high rate solar neutrino detector with neutrino energy determination, in *International School on Cosmological Dark Matter, Valencia, Spain* (1994), pp. 205–228.
- [80] C. Arpesella, C. Broggin, and C. Cattadori, A possible gas for solar neutrino spectroscopy, *Astropart. Phys.* **4**, 333 (1996).
- [81] M. Leyton, S. Dye, and J. Monroe, Exploring the hidden interior of the Earth with directional neutrino measurements, *Nat. Commun.* **8**, 15989 (2017).
- [82] Z. Yang *et al.*, First attempt of directionality reconstruction for atmospheric neutrinos in a large homogeneous liquid scintillator detector, *Phys. Rev. D* **109**, 052005 (2024).
- [83] M. Askins *et al.* (Theia Collaboration), THEIA: An advanced optical neutrino detector, *Eur. Phys. J. C* **80**, 416 (2020).
- [84] M. Agostini *et al.* (BOREXINO Collaboration), First directional measurement of sub-MeV solar neutrinos with Borexino, *Phys. Rev. Lett.* **128**, 091803 (2022).
- [85] M. Agostini *et al.* (BOREXINO Collaboration), Correlated and integrated directionality for sub-MeV solar neutrinos in Borexino, *Phys. Rev. D* **105**, 052002 (2022).
- [86] D. Basilico *et al.* (BOREXINO Collaboration), Final results of Borexino on CNO solar neutrinos, *Phys. Rev. D* **108**, 102005 (2023).
- [87] J. B. R. Battat *et al.*, Readout technologies for directional WIMP dark matter detection, *Phys. Rep.* **662**, 1 (2016).
- [88] S. E. Vahsen, C. A. J. O'Hare, and D. Loomba, Directional recoil detection, *Annu. Rev. Nucl. Part. Sci.* **71**, 189 (2021).
- [89] P. Grothaus, M. Fairbairn, and J. Monroe, Directional dark matter detection beyond the neutrino bound, *Phys. Rev. D* **90**, 055018 (2014).
- [90] C. A. J. O'Hare, A. M. Green, J. Billard, E. Figueroa-Feliciano, and L. E. Strigari, Readout strategies for directional dark matter detection beyond the neutrino background, *Phys. Rev. D* **92**, 063518 (2015).
- [91] F. Mayet *et al.*, A review of the discovery reach of directional dark matter detection, *Phys. Rep.* **627**, 1 (2016).
- [92] C. A. J. O'Hare, B. J. Kavanagh, and A. M. Green, Time-integrated directional detection of dark matter, *Phys. Rev. D* **96**, 083011 (2017).

- [93] C. A. J. O'Hare, Can we overcome the neutrino floor at high masses?, *Phys. Rev. D* **102**, 063024 (2020).
- [94] J. Billard, L. Strigari, and E. Figueroa-Feliciano, Implication of neutrino backgrounds on the reach of next generation dark matter direct detection experiments, *Phys. Rev. D* **89**, 023524 (2014).
- [95] F. Ruppin, J. Billard, E. Figueroa-Feliciano, and L. Strigari, Complementarity of dark matter detectors in light of the neutrino background, *Phys. Rev. D* **90**, 083510 (2014).
- [96] C. A. J. O'Hare, Dark matter astrophysical uncertainties and the neutrino floor, *Phys. Rev. D* **94**, 063527 (2016).
- [97] J. B. Dent, B. Dutta, J. L. Newstead, and L. E. Strigari, Effective field theory treatment of the neutrino background in direct dark matter detection experiments, *Phys. Rev. D* **93**, 075018 (2016).
- [98] J. B. Dent, B. Dutta, J. L. Newstead, and L. E. Strigari, Dark matter, light mediators, and the neutrino floor, *Phys. Rev. D* **95**, 051701 (2017).
- [99] G. B. Gelmini, V. Takhistov, and S. J. Witte, Casting a wide signal net with future direct dark matter detection experiments, *J. Cosmol. Astropart. Phys.* **07** (2018) 009.
- [100] C. A. J. O'Hare, New definition of the neutrino floor for direct dark matter searches, *Phys. Rev. Lett.* **127**, 251802 (2021).
- [101] B. Carew, A. R. Caddell, T. N. Maity, and C. A. J. O'Hare, Neutrino fog for dark matter-electron scattering experiments, *Phys. Rev. D* **109**, 083016 (2024).
- [102] T. N. Maity, Neutrinos as background and signal in Migdal search, [arXiv:2412.17649](https://arxiv.org/abs/2412.17649).
- [103] S. Vahsen *et al.*, CYGNUS: Feasibility of a nuclear recoil observatory with directional sensitivity to dark matter and neutrinos, [arXiv:2008.12587](https://arxiv.org/abs/2008.12587).
- [104] C. A. J. O'Hare *et al.*, Recoil imaging for dark matter, neutrinos, and physics beyond the Standard Model, in *Snowmass 2021* (2022), [arXiv:2203.05914](https://arxiv.org/abs/2203.05914).
- [105] B. Surrow *et al.*, Micro-pattern gaseous detectors, in *2022 Snowmass Summer Study* (2022), [arXiv:2209.05202](https://arxiv.org/abs/2209.05202).
- [106] N. Agafonova *et al.* (NEWSdm Collaboration), Discovery potential for directional dark matter detection with nuclear emulsions, *Eur. Phys. J. C* **78**, 578 (2018).
- [107] M. C. Marshall, M. J. Turner, M. J. H. Ku, D. F. Phillips, and R. L. Walsworth, Directional detection of dark matter with diamond, *Quantum Sci. Technol.* **6**, 024011 (2021).
- [108] R. Bernabei, P. Belli, F. Nozzoli, and A. Incicchitti, Anisotropic scintillators for WIMP direct detection: Revisited, *Eur. Phys. J. C* **28**, 203 (2003).
- [109] P. Belli *et al.*, Measurements of ZnWO_4 anisotropic response to nuclear recoils for the ADAMO project, *Eur. Phys. J. A* **56**, 83 (2020).
- [110] A. Drukier, K. Freese, A. Lopez, D. Spergel, C. Cantor, G. Church, and T. Sano, New dark matter detectors using DNA or RNA for nanometer tracking, [arXiv:1206.6809](https://arxiv.org/abs/1206.6809).
- [111] L. M. Capparelli, G. Cavoto, D. Mazzilli, and A. D. Polosa, Directional dark matter searches with carbon nanotubes, *Phys. Dark Universe* **9–10**, 24 (2015); **11**, 79(E) (2016).
- [112] C. A. J. O'Hare *et al.*, Particle detection and tracking with DNA, *Eur. Phys. J. C* **82**, 306 (2022).
- [113] J. Schueler, M. Ghrear, S. E. Vahsen, P. Sadowski, and C. Deaconu, Deep learning for improved keV-scale recoil identification in high resolution gas time projection chambers, [arXiv:2206.10822](https://arxiv.org/abs/2206.10822).
- [114] M. Ghrear, P. Sadowski, and S. E. Vahsen, Deep probabilistic direction prediction in 3D with applications to directional dark matter detectors, [arXiv:2403.15949](https://arxiv.org/abs/2403.15949).
- [115] J. B. R. Battat *et al.* (DRIFT Collaboration), Low threshold results and limits from the DRIFT directional dark matter detector, *Astropart. Phys.* **91**, 65 (2017).
- [116] T. Ikeda *et al.*, Direction-sensitive dark matter search with the low-background gaseous detector NEWAGE-0.3b", *Prog. Theor. Exp. Phys.* **2021**, 063F01 (2021).
- [117] E. Baracchini *et al.*, CYGNO: A gaseous TPC with optical readout for dark matter directional search, *J. Instrum.* **15**, C07036 (2020).
- [118] F. D. Amaro *et al.*, The CYGNO experiment, a directional detector for direct dark matter searches, *Nucl. Instrum. Methods Phys. Res., Sect. A* **1054**, 168325 (2023).
- [119] B. D. Almeida *et al.*, Noise assessment of CMOS active pixel sensors for the CYGNO experiment, *Meas. Sci. Technol.* **34**, 125145 (2023).
- [120] F. D. Amaro *et al.* (CYGNO Collaboration), The CYGNO experiment: A directional dark matter detector with optical readout, *J. Instrum.* **18**, C09010 (2023).
- [121] F. D. Amaro *et al.*, Modeling the light response of an optically readout GEM based TPC for the CYGNO experiment, [arXiv:2505.06362](https://arxiv.org/abs/2505.06362).
- [122] C. Deaconu *et al.*, Measurement of the directional sensitivity of dark matter time projection chamber detectors, *Phys. Rev. D* **95**, 122002 (2017).
- [123] M. Ghrear, S. E. Vahsen, and C. Deaconu, Observables for recoil identification in high-definition gas time projection chambers, *J. Cosmol. Astropart. Phys.* **10** (2021) 005.
- [124] D. Santos *et al.*, MIMAC: A micro-tpc matrix for directional detection of dark matter, *EAS Publ. Ser.* **53**, 25 (2012).
- [125] R. Yakabe *et al.*, First limits from a 3D-vector directional dark matter search with the NEWAGE-0.3b' detector, *Prog. Theor. Exp. Phys.* **2020**, 113F01 (2020).
- [126] S. Vahsen, H. Feng, M. Garcia-Sciveres, I. Jaegle, J. Kadyk *et al.*, The directional dark matter detector (D^3), *EAS Publ. Ser.* **53**, 43 (2012).
- [127] I. Jaegle, P. Lewis, M. Garcia-Sciveres, M. Hedges, T. Hemperek, J. Janssen, Q. Ji, D.-L. Pohl, S. Ross, J. Schueler *et al.*, Compact, directional neutron detectors capable of high-resolution nuclear recoil imaging, *Nucl. Instrum. Methods Phys. Res., Sect. A* **945**, 162296 (2019).
- [128] Y. Tao *et al.*, Track length measurement of $^{19}\text{F}^+$ ions with the MIMAC directional dark matter detector prototype, *Nucl. Instrum. Methods Phys. Res., Sect. A* **985**, 164569 (2021).
- [129] C. Lisotti *et al.*, CYGUS: Detecting solar neutrinos with directional gas time projection chambers, *Eur. Phys. J. C* **84**, 1021 (2024).
- [130] M. Abdullah, D. Aristizabal Sierra, B. Dutta, and L. E. Strigari, Coherent elastic neutrino-nucleus scattering with directional detectors, *Phys. Rev. D* **102**, 015009 (2020).

- [131] W. C. Haxton, R. G. Hamish Robertson, and A. M. Serenelli, Solar neutrinos: Status and prospects, *Annu. Rev. Astron. Astrophys.* **51**, 21 (2013).
- [132] M. Blennow, T. Ohlsson, and H. Snellman, Day-night effect in solar neutrino oscillations with three flavors, *Phys. Rev. D* **69**, 073006 (2004).
- [133] N. Mishra and L. E. Strigari, Solar neutrinos with CE ν NS and flavor-dependent radiative corrections, *Phys. Rev. D* **108**, 063023 (2023).
- [134] K. J. Kelly, N. Mishra, M. Rai, and L. E. Strigari, $\nu\mu$ and $\nu\tau$ elastic scattering in Borexino, *Phys. Rev. D* **110**, 113004 (2024).
- [135] S. Navas *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
- [136] T. Ohlsson, Status of non-standard neutrino interactions, *Rep. Prog. Phys.* **76**, 044201 (2013).
- [137] B. Dev *et al.*, Neutrino non-standard interactions: A status report, *SciPost Physics Proceedings*, 10.21468/scipostphysproc.2.001 (2019).
- [138] S. Davidson, C. Pena-Garay, N. Rius, and A. Santamaria, Present and future bounds on nonstandard neutrino interactions, *J. High Energy Phys.* **03** (2003) 011.
- [139] M. Lindner, W. Rodejohann, and X.-J. Xu, Coherent neutrino-nucleus scattering and new neutrino interactions, *J. High Energy Phys.* **03** (2017) 097.
- [140] O. G. Miranda and H. Nunokawa, Non standard neutrino interactions: Current status and future prospects, *New J. Phys.* **17**, 095002 (2015).
- [141] Y. Farzan and I. M. Shoemaker, Lepton flavor violating non-standard interactions via light mediators, *J. High Energy Phys.* **07** (2016) 033.
- [142] J. Gehrlein, P. A. N. Machado, and J. P. Pinheiro, Constraining non-standard neutrino interactions with neutral current events at long-baseline oscillation experiments, *J. High Energy Phys.* **05** (2025) 065.
- [143] I. Esteban, M. C. Gonzalez-Garcia, M. Maltoni, I. Martinez-Soler, J. P. Pinheiro, and T. Schwetz, Nufit-6.0: Updated global analysis of three-flavor neutrino oscillations, *J. High Energy Phys.* **12** (2024) 216.
- [144] S. Torelli, Feasibility of a directional solar neutrino measurement with the CYGNO/INITIUM experiment, [arXiv:2408.03760](https://arxiv.org/abs/2408.03760).
- [145] M. Ghreear and S. E. Vahsen, Angular resolution of electrons in gaseous targets, [arXiv:2502.12387](https://arxiv.org/abs/2502.12387).
- [146] I. J. Arnquist, C. Beck, M. L. di Vacri, K. Harouaka, and R. Saldanha, Ultra-low radioactivity Kapton and copper-Kapton laminates, *Nucl. Instrum. Methods Phys. Res., Sect. A* **959**, 163573 (2020).
- [147] G. Cowan, K. Cranmer, E. Gross, and O. Vitells, Asymptotic formulae for likelihood-based tests of new physics, *Eur. Phys. J. C* **71**, 1554 (2011); **73**, 2501(E) (2013).
- [148] I. Esteban, M. C. Gonzalez-Garcia, M. Maltoni, I. Martinez-Soler, and J. Salvado, Updated constraints on non-standard interactions from global analysis of oscillation data, *J. High Energy Phys.* **08** (2018) 180.
- [149] N. Vinyoles, A. M. Serenelli, F. L. Villante, S. Basu, J. Bergström, M. C. Gonzalez-Garcia, M. Maltoni, C. Peña-Garay, and N. Song, A new generation of standard solar models, *Astrophys. J.* **835**, 202 (2017).
- [150] J. D. Lewin and P. F. Smith, Review of mathematics, numerical factors, and corrections for dark matter experiments based on elastic nuclear recoil, *Astropart. Phys.* **6**, 87 (1996).
- [151] R. Abdel Khaleq, J. L. Newstead, C. Simenel, and A. E. Stuchbery, Detailed nuclear structure calculations for coherent elastic neutrino-nucleus scattering, *Phys. Rev. D* **111**, 033003 (2025).