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# Consequences of the Dresden-II reactor data for the weak mixing angle and new physics

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ABSTRACT: The Dresden-II reactor experiment has recently reported a suggestive evidence for the observation of coherent elastic neutrino-nucleus scattering, using a germanium detector. Given the low recoil energy threshold, these data are particularly interesting for a low-energy determination of the weak mixing angle and for the study of new physics leading to spectral distortions at low momentum transfer. Using two hypotheses for the quenching factor, we study the impact of the data on: (i) The weak mixing angle at a renormalization scale of  $\sim 10 \,\mathrm{MeV}$ , (ii) neutrino generalized interactions with light mediators, (iii) the sterile neutrino dipole portal. The results for the weak mixing angle show a strong dependence on the quenching factor choice. Although still with large uncertainties, the Dresden-II data provide for the first time a determination of  $\sin^2 \theta_W$  at such scale using coherent elastic neutrino-nucleus scattering data. Tight upper limits are placed on the light vector, scalar and tensor mediator scenarios. Kinematic constraints implied by the reactor anti-neutrino flux and the ionization energy threshold allow the sterile neutrino dipole portal to produce up-scattering events with sterile neutrino masses up to  $\sim 8 \,\mathrm{MeV}$ . In this context, we find that limits are also sensitive to the quenching factor choice, but in both cases competitive with those derived from XENON1T data and more stringent that those derived with COHERENT data, in the same sterile neutrino mass range.

KEYWORDS: Electroweak Precision Physics, Neutrino Interactions, Non-Standard Neutrino Properties, Sterile or Heavy Neutrinos

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#### 1 Introduction

Since first observed by the COHERENT collaboration in 2017 [1] with a CsI detector, and subsequently in 2020 with a liquid argon (LAr) detector [2], coherent elastic neutrinonucleus scattering (CE $\nu$ NS) has been recognized as a powerful tool for Standard Model (SM) measurements and beyond-the-SM (BSM) searches. Examples of the physics cases that can be studied range from the determination of the mean-square radii of neutron distributions and low-energy measurements of the weak mixing angle [3–8], up to searches for new interactions in the neutrino sector covering a whole spectrum of possible mediators (see e.g. [9–29]). Interestingly, the same experimental infrastructures used for CE $\nu$ NS measurements, provide as well environments suitable for searches of new degrees of freedom involving light dark matter (LDM) [30–33] and axion-like particles (ALPs) [34, 35].

Motivated by this wide range of possibilities, plans for further CE $\nu$ NS measurements are underway. They involve experiments using reactor neutrinos (e.g. CONUS [36–38], CONNIE [39], MINER [40], RED-100 [41],  $\nu$ -cleus [42], TEXONO [43], vIOLETA [44], SBC [45] and the Dresden-II reactor experiment [46]), measurements at COHERENT with germanium and NaI detectors [47], the Coherent CAPTAIN-Mills (CCM) experiment [48] as well as at the European Spallation Source (ESS) [49]. Plans to extended measurements/searches with decay-in-flight neutrino beams such as NuMI [50] or LBNF [51] using gaseous targets with the directional  $\nu$ BDX-DRIFT are as well expected [52, 53]. Measurements of CE $\nu$ NS in multi-ton dark matter (DM) detectors and at RES-NOVA, using archaeological lead, are part of the facilities in which CE $\nu$ NS will be looked for [54–63]. Overall, an international program covering the different energy windows where CE $\nu$ NS can be observed is well established. These energy windows offer features that make them particularly suitable for certain types of new physics searches. Pulsed decay-at-rest (DAR) neutrino beams (such as those at the spallation neutron source and the ESS) provide energy and timing spectra, thus making them particularly useful in searches for flavor-dependent new physics. Decay-in-flight (DIF) neutrino beams — instead — are rather suited for testing nuclear physics hypotheses, due to their higher energy. Finally, given the extremely low-energy thresholds of reactor experiments, sensitivity to physics producing spectral distortions at low momentum transfer becomes a main target. Arguably, the prototypical scenario in that case corresponds to neutrino magnetic moments and transitions, for which the differential cross section exhibits a Coloumb divergence [64]. Scenarios with light mediators, although not leading to such pronounced spectral features, can also be tested with reactor data.

In this regard the recent suggestive observation of  $CE\nu NS$  by the Dresden-II reactor experiment [65] offers an opportunity to systematically test the presence of such new light mediators. The Dresden-II reactor experiment consists of a 2.924 kg p-type point contact germanium detector (NCC-1701) operating at 0.2 keV<sub>ee</sub> and located at ~ 10 m from the 2.96 GW Dresden-II nuclear reactor. The data released follow from a 96.4 days exposure with 25 days of reactor operation outages in which no visible  $CE\nu NS$  signal was observed. Analyses relying on these data and investigating the implications of a modified Lindhard quenching factor (QF) as well as limits on light vector mediators have been already presented in ref. [66]. These data have been used also to place limits on a variety of new physics scenarios including neutrino non-standard interactions (NSI), light vector and scalar mediators and neutrino magnetic moments in ref. [67].

In this paper we extend upon these analyses and consider the impact of the Dresden-II reactor data on: (i) Low-energy measurements of the weak mixing angle at a  $\mu \simeq 10 \text{ MeV}$  renormalization scale, (ii) neutrino generalized interactions (NGI) with light mediators, of which light vector and scalar mediators are a subset, (iii) neutrino magnetic transition couplings leading to up-scattering events (the so-called sterile neutrino dipole portal [68, 69],  $\bar{\nu}_e + N \rightarrow F_4 + N$  with  $F_4$  a heavy sterile neutrino).

The remainder of this paper is organized as follows. In section 2 we briefly present the physics scenarios treated in our statistical analysis, including a short discussion on how the weak mixing angle can affect the event rate. In section 3 we discuss differential event rates, total event rates and the details of the statistical analysis we have adopted along with our results. Finally, in section 4 we present our summary and conclusions.

# 2 CE $\nu$ NS differential cross section, weak mixing angle and new physics scenarios

In the SM the  $CE\nu NS$  differential cross section follows from a t-channel neutral current process and reads [70, 71]

$$\left. \frac{d\sigma}{dE_r} \right|_{\rm SM} = \frac{G_F \, m_N}{2\pi} Q_W^2 F^2(q^2) \left( 2 - \frac{m_N E_r}{E_\nu^2} \right) \,, \tag{2.1}$$

where  $G_F$  refers to the Fermi constant,  $m_N$  to the nuclear target mass,  $E_r$  to recoil energy,  $E_{\nu}$  to the incoming neutrino energy and  $Q_W$  to the weak charge coupling, that accounts for the  $Z^0$ -nucleus interaction in the zero momentum transfer limit. Since the scatterer has an internal structure, this coupling is weighted by the nuclear weak form factor  $F^2(q^2)$ .<sup>1</sup> Hence, the "effective" coupling  $Q_W \times F(q^2)$  encapsulates the expected behavior: as the momentum transfer q increases, the weak charge diminishes and so does the strength of the interaction. Neglecting higher-order momentum transfer terms that arise from the nucleon form factors, one explicitly has

$$Q_W = Z g_{V,\text{SM}}^p + (A - Z) g_{V,\text{SM}}^n.$$
(2.2)

Here the proton and neutron vector couplings are dictated by the fundamental  $Z^0 - q$ (q = u, d) couplings, given by  $g_{V,\text{SM}}^p = 1/2 - 2\sin^2\theta_W$  and  $g_{V,\text{SM}}^n = -1/2$ . For the value of the weak mixing angle at  $\mu = m_{Z^0}$ ,  $\sin^2\theta_W|_{\overline{\text{MS}}}(m_{Z^0}) = 0.23122 \pm 0.00003$  [72], one can easily check that the neutron coupling exceeds the proton coupling by about a factor 10, resulting in the  $N^2 = (A - Z)^2$  dependence predicted in the SM for the CE $\nu$ NS cross section. However, a fair amount of events allows for sensitivities to  $\sin^2\theta_W$ . The SM predicted value at q = 0 (obtained by RGE extrapolation in the minimal subtraction ( $\overline{\text{MS}}$ ) renormalization scheme) is

$$\sin^2 \theta_W(q=0) = \kappa(q=0)|_{\overline{\mathrm{MS}}} \sin^2 \theta_W|_{\overline{\mathrm{MS}}}(m_{Z^0}), \qquad (2.3)$$

with  $\kappa(q=0)|_{\overline{\text{MS}}} = 1.03232 \pm 0.00029$  [73]. Variations around this value lead to fluctuations of the predicted cross section and of the event rate (see section 3). Although statistical analyses of the weak mixing angle have been performed in the light of COHERENT data [5, 7] and are expected to follow also from the electron channel at e.g. DUNE [74], the interesting aspect of an analysis using reactor data has to do with the different energy scale of such an indirect measurement (compared with COHERENT or DUNE) and potentially with the amount of data.

#### 2.1 Renormalizable NGI

Effective NGI<sup>2</sup> were first considered by T.D. Lee and Cheng-Ning Yang in ref. [75]. They have been as well considered in the context of neutrino propagation in matter in ref. [76]. More recently they have been considered in the context of  $CE\nu NS$  analyses in ref. [21] and within COHERENT CsI measurements in ref. [19].<sup>3</sup> Although the Dresden-II reactor data can be used to analyze effective NGI, given its rather low recoil energy threshold one could expect beforehand that better sensitivities to NGI induced by light mediators are achievable. Note that an analysis of this scenario in the context of multi-ton DM detectors has been presented recently in ref. [77].

<sup>&</sup>lt;sup>1</sup>For DAR and DIF neutrino beams the form factor plays an important role. For reactor neutrinos, instead, the energy regime is such that to a large degree  $F^2(q^2) \to 1$ .

<sup>&</sup>lt;sup>2</sup>In contrast to the standard effective interaction jargon, here the typical energy scale has to be just above the MeV scale. For reactor experiments this means  $\Lambda > q \simeq 19$  MeV.

<sup>&</sup>lt;sup>3</sup>In this Reference the acronym NGI, and thus the name "neutrino generalized interactions" rather than generalized neutrino interactions, was introduced as to mimic the acronym NSI for "neutrino nonstandard interactions".

Focusing on this case, the most general Lagrangian can be written schematically as follows

$$\mathcal{L}_{\nu-\mathbf{q}} = \sum_{\substack{X=S,P\\V,A,T}} \left[ \overline{\nu} f_X \Gamma_X \nu X + \sum_{\mathbf{q}=u,d} \overline{\mathbf{q}} \Gamma_X \left( g_X^{\mathbf{q}} + i\gamma_5 h_X^{\mathbf{q}} \right) \mathbf{q} X \right], \qquad (2.4)$$

where  $\Gamma_X = \{\mathbb{I}, i\gamma_5, \gamma_\mu, \gamma_\mu\gamma_5, \sigma_{\mu\nu}\}$  with  $\sigma_{\mu\nu} = i[\gamma_\mu, \gamma_\nu]/2$ , the parameters in the quark and neutrino currents  $(f_X, g_X^q)$  and  $h_X^q)$  are taken to be real and the interactions to be lepton flavor universal. Here, X refers to the field responsible for the interaction. Integrating X out leads to an effective Lagrangian that contains, among other terms, NSI as a subset. In the absence of a robust deviation from the SM CE $\nu$ NS prediction, there is no a priori reason for any of these interactions to be preferred over the others. However, those involving nuclear spin (spin-dependent interactions) are expected to produce lower event rates, in particular in heavy nuclei [71]. Dropping those couplings and moving from quark to nuclear operators the resulting Lagrangian reads

$$\mathcal{L}_{\nu-N} = \sum_{X=\text{All}} \overline{\nu} f_X \Gamma_X \nu X + \sum_{\substack{X=S,V,T}} \overline{N} \, \overline{C}_X \Gamma_X N X + \sum_{\substack{(X,Y)=(P,S)\\(A,V)}} \overline{N} \, i \overline{D}_X \Gamma_Y N X \,. \tag{2.5}$$

Expressions for the coupling of the nucleus to the corresponding mediator are given by [19]

$$\overline{C}_{S} = Z \sum_{q} \frac{m_{p}}{m_{q}} f_{T_{q}}^{p} g_{S}^{q} + (A - Z) \sum_{q} \frac{m_{n}}{m_{q}} f_{T_{q}}^{n} g_{S}^{q}, \qquad (2.6)$$

$$\overline{C}_V = Z(2g_V^u + g_V^d) + (A - Z)(g_V^u + 2g_V^d), \qquad (2.7)$$

$$\overline{C}_T = Z(\delta^p_u g^u_T + \delta^p_d g^d_T) + (A - Z)(\delta^n_u g^u_T + \delta^n_d g^d_T), \qquad (2.8)$$

where the different nucleon coefficients are obtained from chiral perturbation theory from measurements of the  $\pi$ -nucleon sigma term and from data of azimuthal asymmetries in semi-inclusive deep-inelastic-scattering and  $e^+e^-$  collisions [78–82]. Expressions for  $D_P$  and  $D_A$  can be obtained by replacing  $g_S^q \to h_P^q$  and  $g_V^q \to h_A^q$  in  $C_S$  and  $C_V$ , respectively.

The differential cross section induced by the simultaneous presence of all the interactions in eq. (2.5) can be adapted to the light mediator case from the result derived in the effective limit in refs. [19, 21]

$$\frac{d\sigma}{dE_r}\Big|_{\text{NGI}} = \frac{G_F^2}{2\pi} m_N F^2(q^2) \left[\xi_S^2 \frac{2E_r}{E_r^{\text{max}}} + \xi_V^2 \left(2 - \frac{2E_r}{E_r^{\text{max}}}\right) + \xi_T^2 \left(2 - \frac{E_r}{E_r^{\text{max}}}\right)\right].$$
(2.9)

Here  $E_r^{\text{max}} \simeq 2E_{\nu}^2/m_N$  and, in contrast to the effective case, the  $\xi_X$  parameters are  $q^2 = 2m_N E_r$  dependent, though they follow the same definitions

$$\xi_S^2 = C_S^2 + D_P^2, \quad \xi_V^2 = C_V^2 + D_A^2, \quad \xi_T^2 = 4C_T^2.$$
 (2.10)

The parameters in the right-hand side are in turn defined as:

$$C_X = \frac{1}{\sqrt{2}G_F} \frac{f_X \overline{C}_X}{2m_N E_r + m_X^2}, \quad D_X = \frac{1}{\sqrt{2}G_F} \frac{f_X \overline{D}_X}{2m_N E_r + m_X^2}, \quad (2.11)$$

with the exception of  $C_V$  which is shifted by the SM contribution,  $C_V \to Q_W + C_V$ , with  $Q_W$  given by eq. (2.2). Two relevant remarks follow from the expressions in eqs. (2.10) and (2.11). First of all, one can notice that in the low momentum transfer limit and with  $m_X \ll q$  the  $\xi_X$  parameters are enhanced. This is at the origin of the spectral distortions that could be expected if any of these interactions sneaks in the signal. Secondly, unlike the effective case, where each  $\xi_X$  can be treated as a free parameter (thus allowing to encapsulate various interactions at the same time, e.g. in  $\xi_S$  a scalar and pseudoscalar interaction), in this case the  $q^2$  dependence does not allow that. Thus, if one considers e.g.  $\xi_S$ , in full generality a four-parameter analysis is required. To assess the impact of the Dresden-II reactor experiment signal, we then proceed by assuming a single mediator at a time:  $\xi_S$  determined only by  $C_S$  and  $\xi_V$  by  $Q_W + C_V$ . Let us finally note that for the case of  $\xi_T$ , such an assumption is not necessary.

#### 2.2 Sterile neutrino dipole portal

In the Dirac case neutrino magnetic and electric dipole moment couplings are dictated by the following Lagrangian [83]

$$\mathcal{L} = \overline{\nu} \,\sigma_{\mu\nu} \,\lambda \,\nu_R \,F^{\mu\nu} + \text{H.c.}\,, \qquad (2.12)$$

where in general  $\lambda$  is a  $3 \times N$  matrix in flavor space. These couplings are chirality flipping and so the scattering process induced by an ingoing active neutrino produces a sterile neutrino in the final state. Thus, Dirac neutrino magnetic moments always induce up-scattering processes ( $\nu_L + N \rightarrow F_4 + N$ ). The mass of the outgoing fermion, being a free parameter, is only constrained by kinematic criteria. Given an ingoing neutrino energy  $E_{\nu}$ , its mass obeys the following relation:

$$m_4^2 \lesssim 2m_N E_r \left(\sqrt{\frac{2}{m_N E_r}} E_\nu - 1\right) \,. \tag{2.13}$$

For the nuclear recoil energies involved at the Dresden-II experiment and for neutrino energies near the kinematic threshold,  $E_{\nu} \sim 9.5 \,\text{MeV}$ , the upper bound  $m_4 \lesssim 8 \,\text{MeV}$  applies.

The interactions in eq. (2.12) contribute to the  $CE\nu NS$  cross section [68] through

$$\frac{d\sigma}{dE_r}\Big|_{\rm DP} = \alpha_{\rm EM} \,\mu_{\nu,\rm Eff}^2 F^2(q^2) Z^2 \left[ \frac{1}{E_r} - \frac{1}{E_\nu} - \frac{m_4^2}{2E_\nu E_r m_N} \left( 1 - \frac{E_r}{2E_\nu} + \frac{m_N}{2E_\nu} \right) + \frac{m_4^4(E_r - m_N)}{8E_\nu^2 E_r^2 m_N^2} \right]. \tag{2.14}$$

Here  $\alpha_{\rm EM}$  refers to the electromagnetic fine structure constant and  $\mu_{\nu,\rm Eff}$  to a dimensionless [normalized to the Bohr magneton,  $\mu_B = e/(2m_e)$ ] parameter space function that involves combinations of the entries of the  $\lambda$  matrix weighted by neutrino mixing angles and possible CP phases (for details see [23, 28]). Note that in the limit  $m_4 \rightarrow 0$ , eq. (2.14) matches the "standard" neutrino magnetic moment cross section [64].

#### 3 The data, the recoil spectrum and the statistical analysis

In this section we present a brief discussion of the data reported by the Dresden-II reactor experiment, provide the technical tools that allow the calculation of the  $CE\nu NS$  signal (within the SM and with new physics) and present our statistical analysis along with our results for the scenarios discussed in section 2.

#### 3.1 Data and recoil spectra

The Dresden-II reactor experiment consists of a p-type point contact (PPC) 2.924 kg ultra-low noise and low energy threshold  $(0.2 \text{ keV}_{ee})$  germanium detector located at ~ 10 m from the 2.96 GW Dresden-II boiling water reactor (BWR): the NCC-1701 detector [46]. The proximity to the detector along with its high power implies a high flux of electron anti-neutrinos. The data accumulated during 96.4 days of effective exposure with the reactor operating at nominal power (Rx-ON), hint to a first ever observation of CE $\nu$ NS using reactor neutrinos, as recently reported in ref. [65]. The residual difference between the full spectrum and the best-fit background components (the suggested CE $\nu$ NS signal) spans over the measured energy range  $E_M \subset [0.2, 0.4] \text{ keV}_{ee}$  and involves 20 data bins equally spaced (0.01 keV<sub>ee</sub>), as shown in figure 1.

The CE $\nu$ NS differential recoil energy spectrum follows from a convolution of the electron anti-neutrino flux and the CE $\nu$ NS cross section, namely

$$\frac{dR}{dE_r} = N_T \int_{E_{\nu}^{\min}}^{E_{\nu}^{\max}} \frac{d\Phi_{\overline{\nu}_e}}{dE_{\nu}} \frac{d\sigma_{\text{CE}\nu\text{NS}}}{dE_r} \, dE_{\nu} \,. \tag{3.1}$$

The number of germanium nuclei in the detector is given by  $N_T = m_{\text{det}} N_A/m_{\text{Ge}}$ , with  $N_A$  the Avogadro number,  $m_{\text{Ge}}$  the germanium molar mass and  $m_{\text{det}} = 2.924 \,\text{kg}$ . The integration limits are given by  $E_{\nu}^{\min} = \sqrt{m_N E_r/2}$ , with  $E_r$  being the recoil energy, and  $E_{\nu}^{\max}$  the kinematic value determined by the electron anti-neutrino flux. We take the values of the atomic number and nuclear mass for <sup>72</sup>Ge. For neutrino energies below 2 MeV we use the anti-neutrino spectral function from Kopeikin [84], while for energies above that value we consider Mueller et al. [85]. For flux normalization we use  $\mathcal{N} = 4.8 \times 10^{13} \overline{\nu}_e/\text{cm}^2/\text{sec}$ , as given in ref. [65]. The differential anti-neutrino flux in eq. (3.1) therefore involves the spectral function and the normalization. The CE $\nu$ NS differential cross section is dictated by eq. (2.1), but can also involve contributions from NGI couplings or the sterile neutrino dipole portal discussed in section 2.1 and 2.2.

For detectors relying on ionization (it applies to scintillation as well), such as the NCC-1701, only a fraction of the nuclear recoil energy is available in a readable format. The characterization of that fraction is given by the QF, Q, defined as the ratio between the nuclear recoil given in ionization  $(E_I)$  over that generated by an electron recoil of the same kinetic energy  $(E_r)$ . Quantitatively, this means that the ionization energy expected from a given recoil energy is given by  $E_I = Q E_r$ . With the aid of the QF, the differential ionization spectrum can then be written according to

$$\frac{dR}{dE_I} = \frac{dR}{dE_r}\frac{dE_r}{dE_I} = \frac{dR}{dE_r}\left(\frac{1}{Q} - \frac{E_I}{Q^2}\frac{dQ}{dE_I}\right).$$
(3.2)



Figure 1. Experimental data from the Dresden-II reactor obtained during 96.4 days exposure time using the NNC-1701 germanium detector. CE $\nu$ NS data follow from residual counts after the subtraction of the best-fit background components [65]. The spectral rates of signal events are also shown, for the SM prediction obtained with the modified Lindhard QF [see eq. (3.3)] (gray curves, solid for q = 0 and dashed for  $q = -20 \times 10^{-5}$ , in both cases k = 0.157) and for various new physics scenarios, with same assumptions on the QF.

For sufficiently high-recoil energy regimes (above  $5 \text{ keV}_{nr}$  or so) the QF is well described by the Lindhard model [86]. However, its validity is questionable in any material for sub-keV energies, as pointed out in ref. [87]. For germanium, recent measurements of its QF using recoils from gamma emission following thermal neutron capture, photo-neutron sources, and a monochromatic filtered neutron beam have shown substantial deviations from the Lindhard model expectations at recoil energies below  $\sim 1.3 \text{ keV}_{nr}$  [88]. In the context of DM direct detection searches, ref. [89] has addressed this issue providing a slight modification of the Lindhard QF

$$Q(E_r) = \frac{k g(\epsilon)}{1 + k g(\epsilon)} - \frac{q}{\epsilon}, \qquad (3.3)$$

where the first term is the standard Lindhard QF with  $g(\epsilon) = 3\epsilon^{0.15} + 0.7\epsilon^{0.6} + \epsilon$  and  $\epsilon = 11.5Z^{-7/3} E_r$ . The second term (the correction) is such that deviations from the standard behavior start to show up at about 0.1 keV. In our analyses we adopt this parametrization, and therefore we include k and q as free parameters. In addition to this QF, we employ as well the "iron-filter" QF reported in the ancillary files provided by ref. [65].

The CE $\nu$ NS ionization differential spectrum in eq. (3.2) has to be smeared by the intrinsic resolution of the detector. Following the information of the README ancillary file [65], we take the resolution to be a Gaussian truncated energy-dependent distribution given by [67]

$$G(E_M, E_I, \sigma) = \frac{2}{1 + \operatorname{erf}\left(E_I/\sqrt{2}/\sigma\right)} \frac{1}{\sqrt{2\pi\sigma}} e^{-\Delta E^2/2/\sigma^2}.$$
 (3.4)

Here, the energy-dependent Gaussian width  $\sigma^2 = \sigma_n^2 + E_I \eta F$  involves the intrinsic electronic noise of the detector  $\sigma_n = 68.5 \text{ eV}$  (for the 96.4 days of Rx-ON data), the average energy of  $e^-$ -hole formation in germanium  $\eta = 2.96 \text{ eV}$ , and the Fano factor whose value we fix to the average value in the range [0.10–0.11], F = 0.105. As stressed in the ancillary file, overall the second term in the Gaussian width measures the dispersion in the number of information carriers ( $e^-$ -hole pairs).

The smearing of the ionization differential spectrum results in the measured differential spectrum

$$\frac{dR}{dE_M} = \int_{\eta}^{\infty} G(E_M, E_I, \sigma) \frac{dR}{dE_I} dE_I , \qquad (3.5)$$

from which the number of events in the *i*th bin is obtained by integration over the measured energy  $E_M$ , in the interval  $[E_M^i - \Delta E_M, E_M^i + \Delta E_M]$  ( $\Delta E_M = 5 \, \text{eV}_{ee}$ ). The integration lower limit is set by the minimum average ionization energy  $\eta \sim 3 \, \text{eV}_{ee}$  required to produce an  $e^-$ -hole pair in germanium.

#### 3.2 Statistical analysis

Our analysis is based on the  $\chi^2$  function

$$\chi^2(\vec{S},\alpha) = \sum_i \left[ \frac{N_{\rm th}^i(\vec{S},\alpha) - N_{\rm meas}^i}{\sigma_i} \right]^2 + \left( \frac{\alpha}{\sigma_\alpha} \right)^2 \,, \tag{3.6}$$

where  $N_{\rm th}^i$  and  $N_{\rm meas}^i$  are the theoretical and measured number of events, respectively, in the *i*th energy bin. Note that in the definition of the  $\chi^2$  function we are assuming the data to follow a Gaussian distribution. Although assuming a Poisson distribution would be a better choice given the dataset, both statistical and systematic errors (which have a bigger impact on the results) can be readily included under the Gaussian assumption. Here,  $\sigma_i$ represents the corresponding uncertainty of the *i*th measurement which includes systematic and statistical uncertainties. Here,  $\vec{S}$  represents a set of new physics parameters, while  $\alpha$  is a nuisance parameter which accounts for the flux normalization uncertainty, for which we consider  $\sigma_{\alpha} = 5\%$ . The theoretical number of events is

$$N_{\rm th}(\vec{S},\alpha) = (1+\alpha)N_{\rm CEvNS}(\vec{S}), \qquad (3.7)$$

which, of course, includes the SM piece in addition to the new physics contribution.

Equipped with the tools discussed in section 3 along with the  $\chi^2$  function in eq. (3.6), we begin our discussion by focusing on the implications for the weak mixing angle. Figure 2 shows the  $\Delta \chi^2$  distributions in terms of  $\sin^2 \theta_W$  for the two QFs considered in the analysis. In the case of the modified Lindhard QF, our result is obtained by marginalizing over the



**Figure 2**.  $\Delta \chi^2$  profiles for  $\sin^2 \theta_W$  for the two QFs considered [modified Lindhard QF, eq. (3.3), and iron-filter QF as given in the ancillary files in ref. [65]]. For the modified Lindhard QF the result follows after marginalization over k and q [see eq. (3.3)].

parameters k and q [see eq. (3.3)]. Notice that the  $\Delta \chi^2$  profile for the case of Lindhard QF is rather flat at the bottom, thus making its best fit value not very statistically meaningful. Specifically, the Lindhard parameters are allowed to float in the ranges<sup>4</sup> 0.14  $\leq k \leq 0.27$ and  $-40 \leq q/10^{-5} \leq 0$ . As expected, a strong dependence on the QF is observed. The best-fit values differ by about ~ 6.5%, with the iron-filter QF favoring a larger  $\sin^2 \theta_W$ value. The 1 $\sigma$  ranges read

Modified Lindhard QF: 
$$\sin^2 \theta_W = 0.178^{+0.280}_{-0.090}$$
  
Iron filter QF:  $\sin^2 \theta_W = 0.190^{+0.039}_{-0.046}$ , (3.8)

thus showing the disparity of the values obtained as a consequence of a different QF model. One can notice as well that both values differ substantially from the SM RGE expectation. In particular, the best fit result from the iron-filter QF analysis is compatible with the SM RGE prediction at 80.7% C.L., whereas the result from the modified Lindhard QF is in agreement at  $1\sigma$ , given the spread of its  $\Delta\chi^2$  distribution. From these results one can conclude that with the current data set and the lack of a better knowledge of the germanium QF, a robust determination of the weak mixing angle seems not possible.

Although featuring a moderate disparity, these results can be understood as a first determination of the weak mixing angle at low energies using  $CE\nu NS$  data from reactor antineutrinos. They can be compared with the values obtained from COHERENT CsI and LAr data [1, 2] and other dedicated experiments that include atomic parity violation (APV) [90, 91], proton weak charge from Cs transitions ( $Q_{\text{weak}}$ ) [92], Møller scattering (E158) [93], parity violation in deep inelastic scattering (PVDIS) [94] and neutrino-nucleus scattering (NuTeV) [95]. A summary of these results is displayed in figure 3, which shows as well

<sup>&</sup>lt;sup>4</sup>Note, that k = 0.27 corresponds to the limit set by CONUS [37].



Figure 3. Weak mixing angle RGE running in the SM, calculated in the  $\overline{\text{MS}}$  renormalization scheme as obtained in ref. [96], along with measurements at different renormalization scales [1, 2, 90–95]. The  $1\sigma$  result obtained using the Dresden-II reactor data is shown assuming the modified Lindhard and iron-filter QF (see text for further details).

the RGE running calculated in the MS renormalization scheme [96]. The value for the weak mixing angle at the  $1\sigma$  level extracted from the best fit in figure 2 is shown. For the renormalization scale at which the measurement applies, we have adopted a rather simple procedure. We have translated the ionization energy range into recoil energy with the aid of the QF. With the values obtained for  $E_r^{\min}$  and  $E_r^{\max}$  we have then calculated the momentum transfer by using the kinematic relation  $q^2 = 2m_N E_r$ . This result corresponds to the first CE $\nu$ NS-based determination of  $\sin^2 \theta_W$  with reactor data at  $\mu \sim 10$  MeV. With further data, and more importantly a better understanding of the germanium QF, this result is expected to highly improve in the future.

We now move on to the case of NGI. For this analysis we assume universal quark couplings and switch off the pseudovector couplings in the vector case (those controlled by  $\xi_V$ ) as well as the pseudoscalar couplings in the scalar case (those controlled by  $\xi_S$ ). These simplifications reduce the analysis to pure vector and pure scalar interactions, controlled by the couplings  $g_V^2 = g_V^q f_V$  and  $g_S^2 = g_S^q f_S$  (and the mediators masses), as investigated in refs. [66, 67]. For the tensor case no assumption on different contributions is required. The cross section is determined by  $\xi_T$  and, under the assumption of universal quark couplings, it is eventually controlled by  $g_T^2 = g_T^q f_T$ . Again, for the statistical analysis using the modified Lindhard QF we vary as well q and k. The analysis in this case is therefore a four parameter problem, while for the iron-filter QF only two parameters matter, i.e. the new mediator mass and coupling.

Our extracted sensitivities are illustrated in figure 4 at 1, 2, 3  $\sigma$  (assuming two d.o.f., i.e.  $\Delta \chi^2 = 2.3, 6.18, 11.83$  respectively). The upper row stands for the vector case, the middle row for the scalar and the bottom row for the tensor, while left (right) panels are obtained using the modified Lindhard (iron-filter) QF. As can be seen, at the 1 $\sigma$  level and above,



Figure 4. Constraints on vector NGI (upper row), scalar NGI (central row) and tensor NGI (lower row) in the coupling-mass plane, obtained using the modified Lindhard QF (left column) and the iron-filter QF (right column). In all panels, purple regions indicate exclusion limits. Where present, dark blue stars specify the best fit solutions. Moreover, constraints from CONUS [38], CONNIE [97] and COHERENT CsI+LAr [98] are shown for comparison. Additionally, the black dashed vertical line marks the transition from the light to the effective regime.

large portions of parameter space are ruled out, disfavoring couplings as low as  $7.5 \times 10^{-6}$  for  $m_V \lesssim 100$  keV. At the 1 and  $2\sigma$  level, two "islands" in the region of noneffective interactions  $(m_V \gtrsim 10 \text{ MeV})$  are open as well. At the  $3\sigma$  level these spots are gone and the constraint becomes a little less stringent. Turning to the analysis done assuming the iron-filter QF, we find that about the same regions in parameter space are excluded, though the most stringent limit is a little more pronounced in this case  $(4 \times 10^{-6} \text{ for } m_V \lesssim 100 \text{ keV})$ . The parameter space "islands" found with the modified Lindhard QF are present in this case as well, but cover a somewhat wider area. At the 90% C.L. constraints on the vector NGI scenario amount to  $g_V \lesssim 8 \times 10^{-6}$  (Lindhard QF) and  $g_V \lesssim 4.5 \times 10^{-6}$  (iron-filter QF) for vector masses up to 100 keV. This limit should be compared with results from COHERENT CsI and LAr, for which refs. [7, 13] found  $g_V \lesssim 6 \times 10^{-5}$  at the 90% C.L. We can then conclude that the Dresden-II data largely improve limits for vector interactions in the low vector mass window. This result can be attributed to the sub-keV recoil energy threshold the experiment operates with.

In the scalar case the situation is as follows. The modified Lindhard QF and the scalar hypothesis tend to produce smaller deviations from the data. This can be readily understood from the left graph in the bottom row of figure 1. At low scintillation energy the event rate tends to increase, but slightly less than in the vector case, a behavior somehow expected, see e.g. ref. [18]. While the scalar coupling contributes to the CE $\nu$ NS cross section quadratically, the vector does it linearly because of its interference with the SM contribution. As a consequence, at  $1\sigma$  level and above, limits are slightly less stringent than in the vector case. In contrast to that case as well, the parameter space "islands" are gone. Their disappearance can be traced back to the fact that these interactions do not sizably interfere with the SM term. Limits for scalar masses below ~ 1 MeV at the 90% C.L. amount to  $g_S \lesssim 3.0 \times 10^{-5}$  at the 90% C.L., implying a slight improvement on the limit. For the iron-filter QF one finds about the same trend, with limits at different statistical significances spreading uniformly. The 90% C.L. limit at low scalar mass amounts to  $g_S \lesssim 1.8 \times 10^{-6}$ , for scalar masses up to 100 keV.

Results for the light tensor case resemble those found in the NGI light scalar scenario, though limits are a little weaker. At the  $1\sigma$  level and above, we find  $g_T \leq 1.0 \times 10^{-5}$ (Lindhard QF) and  $g_T \leq 6.0 \times 10^{-6}$  (iron-filter QF) for tensor masses below ~ 100 keV. Although with small differences, among the NGI we have considered, the tensor couplings are the less constrained by the Dresden-II data set. This result is inline with that found when analyzing tensor NGI using CSI COHERENT data [19].

To our knowledge, limits on light tensor interactions using COHERENT CsI and LAr data have been discussed only in [99]. On the other hand, there are some forecasts for searches for this type of interactions at multi-ton DM detectors [77]. Searches relying on the CE $\nu$ NS nuclear recoil channel are expected to be sensitive up to  $g_T \sim 2.0 \times 10^{-5}$  for tensor masses up to  $\sim 1$  MeV at the 90% C.L. These numbers lead to the same conclusion than in the scalar case: in the light mediator regime, constraints obtained using Dresden-II data seem to improve upon available sensitivities.

As we have already pointed out, given the kinematic threshold of the electron antineutrino flux and the small ionization energy of the Dresden-II data set, up-scattering



**Figure 5.** Results of the analysis for the sterile neutrino dipole portal based on two QF hypotheses: modified Lindhard QF (left graph) and iron-filter QF (right graph). For the former case results follow after marginalization over q and k. Shaded areas indicate the excluded regions at different statistical significance levels:  $1\sigma$ ,  $2\sigma$  and  $3\sigma$  as shown in the graphs. Constraints from CENNS10, TEXONO, COHERENT CsI and XENON1T (see ref. [28]) are also shown for comparison.

via dipole portal interactions can produce sterile neutrinos with masses up to  $\sim 8 \,\mathrm{MeV}$ . In full generality, one can expect constraints on the effective magnetic dipole moment coupling to be less severe as the mass of the up-scattered fermion increases. The kinematic suppression increases, reaching zero when the sterile neutrino mass hits the kinematic production threshold limit given by eq. (2.13). The 1, 2, 3  $\sigma$  (assuming two d.o.f., i.e.  $\Delta \chi^2 = 2.3, 6.18, 11.83$  respectively) results of our analysis for this case are shown figure 5, left (right) graph obtained with the modified Lindhard (iron-filter) QF. Limits from the different exclusion regions tend to be a little more uniform in terms of the up-scattered sterile state mass, in comparison to the NGI scenarios previously considered in terms of the mediator mass. For the modified Lindhard QF analysis, values of the order of  $\mu_{\nu_e} \lesssim 4 \times 10^{-10} \,\mu_B$  are excluded for sterile neutrino masses below 100 keV, at the  $1\sigma$ level. Assuming instead the modified iron-filter QF the constraints are slightly tighter,  $\mu_{\nu_e} \lesssim 1 \times 10^{-10} \,\mu_B$  for sterile neutrino masses below 100 keV, at the 1 $\sigma$  level. A comparison of these values with those obtained using CsI and LAr COHERENT data sets (shown in the graphs),  $\mu_{\nu_e} \lesssim (3-4) \times 10^{-9} \mu_B$  at the 90% C.L. [28], demonstrates that the Dresden-II experimental data improve upon these results (the 90% C.L. upper limits are  $(2-8) \times 10^{-10} \mu_B$  for  $m_4 \lesssim 100 \,\text{keV}$ ). They are competitive with the constraints implied by XENON1T data (indeed more constraining if one focuses only on the nuclear recoil channel) [69], are stronger than those derived from CENNS10 [28] and comparable (or even tighter) than those following from TEXONO depending on the QF model used for the analysis, as can be read directly from the graphs. If compared with explanations of the XENON1T electron excess using electron neutrinos [28], one can see that our results are consistent with that possibility,<sup>5</sup> regardless of the QF choice. Note that the sterile neutrino dipole portal and NGI results, in contrast to those found for the weak mixing angle, are to

 $<sup>^{5}</sup>$ Explanations of the excess using tau neutrinos are not affected by this result either [100].

a large extent rather insensitive to the QF model. Thus, from that point of view they are more robust.

#### 4 Conclusions

We have studied the implications of the recently released Dresden-II reactor data on the weak mixing angle and on new physics scenarios sensitive to the low-energy threshold of the experiment, namely NGI generated by light vector, scalar and tensor mediators and the sterile neutrino dipole portal. In order to check for the dependences on the QF, we have performed the analyses considering: (i) A modified Lindhard model, (ii) a QF provided by the collaboration (iron-filter QF).

The low scintillation energy threshold provides a determination of the weak mixing angle at a renormalization scale of order 10 MeV, a scale for which up to now no determination was yet available. Our result shows a rather pronounced dependence on the QF model, with differences between the best-fit values of about 6%. The precision of the determination of  $\sin^2 \theta_W$  has also a strong dependence on that choice, leading to best fit values that are compatible with the SM RGE prediction at 80.7% C.L. and 1 $\sigma$ , respectively. A better understanding of the germanium QF is thus required to improve upon the determination of this parameter. However, regardless of these disparities, the Dresden-II data provides the first hint ever of the value of  $\sin^2 \theta_W$  at  $\mu \sim 10$  MeV.

Regarding our analysis of NGI with light mediators, also in this case our findings show that at the  $1\sigma$  level results depend on the QF model. For vector interactions, results derived using the modified Lindhard QF tend to produce slightly less stringent bounds. In both cases, though, at large vector mediator masses (above 10 MeV or so) the  $1\sigma$  and  $2\sigma$  limits produce two nonoverlapping exclusion regions. At the  $3\sigma$  level these regions are gone and constraints are restricted to a single area, where for vector boson masses of the order of 100 keV the coupling is constrained to be below  $\sim 10^{-5}$ .

The same trend is found for scalar and tensor interactions through light mediators. Regardless of the QF choice, results lead to constraints that amount to about  $g_S \leq 1.0 \times 10^{-6}$ and  $g_T \leq 1.0 \times 10^{-5}$ , respectively, for mediator masses below ~ 100 keV at the 1 $\sigma$  level. In all scenarios, the derived constraints turn out to improve upon other existing bounds from CE $\nu$ NS experiments (COHERENT CsI+LAr, CONUS and CONNIE) and even upon predictions made for multi-ton DM detector measurements.

Finally, concerning the sterile neutrino dipole portal we find that the Dresden-II results rule out larger regions of parameter space, not excluded by COHERENT and CONUS and are rather competitive with limits from XENON1T data. Actually, they are more stringent if one compares only with XENON1T nuclear recoil data. Compared with those regions where the sterile neutrino dipole portal can account for the XENON1T electron excess, the Dresden-II data is not able to test them yet. However, with more statistics and better understanding of the germanium QF the situation might improve in the future.

To conclude, the recent evidence for  $CE\nu NS$  from the Dresden-II reactor experiment provides unique opportunities to investigate physics scenarios sensitive to low-energy thresholds, complementary to other  $CE\nu NS$  measurements with spallation sources. However, current results show a dependence on the QF model at low recoil energies thus calling for a deeper understanding of the germanium QF along with more data.

**Note added in proof.** After completion of the manuscript results from the first science run of the XENONnT collaboration [101] have ruled out the electron excess previously reported by the XENON1T collaboration [102].

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