

**KARADENİZ TECHNICAL UNIVERSITY
GRADUATE SCHOOL OF NATURAL AND APPLIED SCIENCES**

DEPARTMENT OF PHYSICS

**NEUTRINO INTERACTIONS BEYOND THE STANDARD MODEL AT DARK MATTER
EXPERIMENTS WITH SOLAR NEUTRINOS**

DOCTORAL DISSERTATION

Muhammad Fauzi MUSTAMIN

**OCTOBER 2025
TRABZON**



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**This dissertation is accepted to give the degree of
DOCTOR OF PHILOSOPHY
by Karadeniz Technical University**

Date of Submission : 02 / 09 / 2025

Date of Examination : 01 / 10 / 2025

Supervisor : Prof. Dr. Mehmet DEMİRCİ

Trabzon 2025

ACKNOWLEDGMENT

This dissertation entitled “Neutrino Interactions Beyond the Standard Model at Dark Matter Experiments with Solar Neutrinos” was prepared in the doctoral program of the Department of Physics, Institute of Science, Karadeniz Technical University. During the completion, this research was supported by TUBITAK 1002-A under project number 123F186 and TUBITAK-COST 2515 under project number 124F416.

I would like to express my heartfelt respect to my supervisor, Prof. Mehmet Demirci. It is a great privilege being able to be advised by him. I have gained a lot from his knowledge and wide horizon on fundamental physics. Thank you for the support and guidance throughout the completion of this work. You have guided me as a physicist and shown me what doing proper research is.

My gratitude is also extended toward the Presidency of Turks Abroad and Related Communities (YTB) for providing me the opportunity to pursue doctoral studies at Karadeniz Technical University. Moreover, I would like to extend this to my Indonesian companions and also the Trabzon International Student Association for the great experiences in the process.

Last but not least, I would like to thank my parents and family, my beloved mother Bahjah, my father Mustamin, my brother Mukram and my sisters Nadhrah and Nihlah. Thank you for supporting me from afar during my academic journey. Ultimately, I would express my gratefulness to the Almighty, who only with His benevolence, brings all particles into harmony in the observable universe, as well as in the yet unobserved one.

Trabzon, October 2025

Muhammad Fauzi MUSTAMIN

STATEMENT OF ETHICS

I declare this work is an original research. It has been conducted by me, under the supervision of Prof. Dr. Mehmet Demirci. It has not been submitted to any previous university or published in any institution. I confirm that this is my own work. The data and information from other researches have been cited in the reference section. Publication ethics and rules have been obeyed during the completion of this work and I am fully responsible for every possible outcome of this work. 01/10/2025

Muhammad Fauzi MUSTAMIN

CONTENTS

	<u>Page No.</u>
ACKNOWLEDGMENT	III
STATEMENT OF ETHICS	IV
CONTENTS	VI
SUMMARY	VII
ÖZET	VIII
LIST OF FIGURES	XII
LIST OF TABLES	XIII
ABBREVIATIONS	XIV
1. INTRODUCTION	1
1.1. Neutrinos in the Standard Model	1
1.2. Neutrino Oscillation Phenomena	4
1.3. Neutrino Sources	6
1.3.1. Artificial sources	9
1.3.2. Natural sources	10
1.4. Scope and Outline of Dissertation	12
2. SOLAR NEUTRINO INTERACTIONS AT DIRECT DETECTION EXPERIMENTS	14
2.1. The Solar Neutrino Spectrum	15
2.2. Elastic Neutrino-Electron Scattering ($E\nu$ ES)	20
2.3. Coherent Elastic Neutrino-Nucleus Scattering ($CE\nu$ NS)	24
2.4. Dark Matter Direct Detection Experiments	31
2.4.1. CDEX-10	33
2.4.2. PandaX-4T	34
2.4.3. XENONnT	36

3. BEYOND THE STANDARD MODEL PHYSICS WITH NEUTRINO INTERACTIONS	38
3.1. New Light Mediator Models	38
3.1.1. Universal Scalar Mediator	39
3.1.2. Universal Vector Mediator	41
3.1.3. Universal Tensor Mediator	43
3.1.4. Anomaly-free $U(1)'$ Models	46
3.2. Sterile Neutrino Dipole Portal	52
3.3. Triplet Higgs Model	55
4. ANALYSIS DETAILS	59
4.1. Expected Event Rates	59
4.2. Likelihood Analysis	62
4.3. Projected Sensitivity	64
5. RESULTS AND DISCUSSIONS	67
5.1. Predicted Event Rates	67
5.2. Previous Limits	71
5.3. New Light Mediator Models	72
5.3.1. Results from Electron Recoil Data	72
5.3.2. Results from Nuclear Recoil Data	78
5.4. Sterile Neutrino Dipole Portal	84
5.4.1. Results from Electron Recoil Data	84
5.4.2. Results from Nuclear Recoil Data	88
5.5. Triplet Higgs Model	92
5.6. Conclusion	94
5.6.1. New light mediator models	94
5.6.2. Active-sterile neutrino transition magnetic moment	96
5.6.3. Charged Higgs boson	97
6. SUGGESTIONS	99
7. REFERENCES	100
8. APPENDICES	109
CURRICULUM VITAE	

Doctoral Thesis
Summary

NEUTRINO INTERACTIONS BEYOND THE STANDARD MODEL AT DARK
MATTER EXPERIMENTS WITH SOLAR NEUTRINOS

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Supervisor: Prof. Dr. Mehmet DEMİRCİ
2025, 108 pages, 7 pages appendix

In this work, we investigate the effects of new physics scenarios on solar neutrino-induced interactions in dark matter direct detection experiments. We consider new light mediator models, where a new light boson mediates additional neutrino interactions; the sterile neutrino dipole portal, which introduces transition magnetic moments between active and sterile neutrino states; and the Higgs triplet model, where singly charged Higgs bosons contribute to neutrino scattering processes. Our analysis focuses on two fundamental processes: elastic neutrino-electron scattering ($E\nu$ ES) and coherent elastic neutrino-nucleus scattering ($CE\nu$ NS). We use datasets from the CDEX-10, PandaX-4T, and XENONnT dark matter experiments. By performing a χ^2 -analysis, we derive new and competitive constraints on the considered models from recent low-energy recoil data of these experiments. The analysis incorporates solar neutrino survival probabilities to account for flavor-dependent effects. For comparison, we also include limits from previous studies based on other experimental data. Our results demonstrate significantly improved sensitivity relative to existing bounds, highlighting the potential of low-energy solar neutrino measurements in advancing the search for physics beyond the Standard Model.

Keywords: Solar neutrino, $CE\nu$ NS, $E\nu$ ES, new light mediator models, sterile neutrino dipole portal, charged Higgs boson.

Doktora Tezi
ÖZET

GÜNEŞ NÖTRİNOLARIYLA YAPILAN KARANLIK MADDE DENEYLERİNDE
STANDART MODEL ÖTESİNDE NÖTRİNO ETKİLEŞMELERİ

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2025, 108 sayfa, 7 sayfa ek

Nötrinolar, yalnızca zayıf etkileşme yoluyla maddeyle etkileşen ve kütlelerinin çok küçük olması nedeniyle doğrudan tespiti son derece güç olan temel parçacıklardır. Standart Model (SM) çerçevesinde başlangıçta kütsüz olarak kabul edilen nötrinoların deneysel olarak gözlemlenen salınımları, bu parçacıkların sıfır olmayan kütlelere sahip olduğunu ortaya koymakta ve dolayısıyla SM'in eksikliğine işaret eden önemli bir gösterge sunmaktadır. Nötrino kütlelerinin kökeni ve doğası, parçacık fiziğinde açık bir problem olarak varlığını sürdürmekte ve yeni fizik senaryolarının geliştirilmesini gerektirmektedir.

Bu tez çalışması, SM ötesi çeşitli etkileşme senaryolarını düşük enerjili güneş nötrinoların maddeyle olan etkileşmeleri üzerinden, özellikle karanlık maddenin doğrudan tespiti amacıyla inşa edilmiş deneylerin güncel verileri kullanılarak araştırmayı amaçlamaktadır. Güneş nötrinoları, bu deneylerde yalnızca önemli bir arkaplan kaynağı olmakla kalmayıp, aynı zamanda yeni fizik etkileşmelerini test etmek için doğal ve sürekli bir nötrino kaynağı olarak da işlev görmektedir. Dolayısıyla, güneş nötrino kaynaklı olayların detaylı biçimde modellenmesi, karanlık madde-nükleon saçılmalarına dair potansiyel sinyallerin ayrıştırılmasında kritik rol oynamaktadır.

Bu bağlamda üç temel SM ötesi etkileşme mekanizması ele alınmıştır: (i) hafif aracı bozonlar aracılığıyla gerçekleşen yeni etkileşmeler, (ii) aktif-sterile nötrino geçiş manyetik momenti üzerinden elektromanyetik saçılma süreçleri ve (iii) yüklü Higgs bozonlarının katkı sağladığı yeni skaler etkileşmeler. Hafif aracı bozon modelleri, minimal yeni alan içeren genişletilmiş senaryoların düşük enerjili limitlerini temsil eder ve burada, özellikle anomali içermeyen $U(1)'$ uzantıları dikkate alınmıştır. Aktif-sterile nötrino geçiş manyetik momenti senaryosu kapsamında, aktif nötrinolar ile sterile nötrino arasındaki dipol etkileşmeler ele

alınmıştır. Bu tür etkileşmeler, nötrinoların madde içerisindeki elektromanyetik alanlarla saçılmasına ve bunun sonucunda sterile nötrinoların üretimine olanak tanımaktadır. Bu mekanizma, özellikle kısa-tabanlı nötrino salınım deneylerinde gözlemlenen anomalilerin açıklanmasında öne çıkan alternatif yaklaşımlardan biridir. Son senaryo olarak üçlü-Higgs modelinde ortaya çıkan ve Majorana nötrino kütlelerini ağaç seviyesinde üretebilen yüklü Higgs bozonu dikkate alınarak nötrino saçılma sürecine katkıları incelenmiştir.

SM ötesi bu senaryolar, güneş nötrinolarıyla gerçekleşen elastik nötrino-elektron saçılması ve koherent elastik nötrino-çekirdek saçılması süreçleri çerçevesinde incelenmiştir. Güneş nötrinolarının çeşni dönüşümleri hesaba katılarak, güncel düşük enerjili karanlık madde deneylerinden (CDEX-10, PandaX-4T ve XENONnT) elde edilen veri setleriyle istatistiksel analizler gerçekleştirilmiş ve ilgili model parametreleri üzerine üst sınırlar türetilmiştir. Elde edilen sınırlar, literatürde mevcut olan sonuçlarla karşılaştırmalı olarak değerlendirilmiş ve birçok senaryo için daha sıkı kısıtlamalar sağlandığı gösterilmiştir. Bu bulgular, özellikle düşük enerjili hassas ölçümlerin, yeni etkileşmeleri araştırmak adına yüksek potansiyele sahip olduğunu ortaya koymaktadır.

Anahtar Kelimeler: Güneş nötrinoları, elastik nötrino-elektron saçılması, koherent elastik nötrino-çekirdek saçılması, SM ötesi etkileşmeler, hafif aracı bozonlar, geçiş manyetik momenti, steril nötrino, yüklü Higgs bozonu.

LIST OF FIGURES

	<u>Page No.</u>
Figure 1. The Grand Unified Neutrino Spectrum at Earth.	7
Figure 2. Representative of neutrino cross sections from various sources.	8
Figure 3. The pp-chains (a) and the CNO cycle (b) processes.	16
Figure 4. Spectrum of solar neutrino fluxes observed on Earth.	18
Figure 5. Solar neutrinos survival probabilities.	20
Figure 6. Feynman diagram representations at tree level for $E\nu$ ES.	21
Figure 7. Feynmann diagram of $CE\nu$ NS in the SM.	24
Figure 8. Various processes involving neutrino.	25
Figure 9. Spin dependent structure factors with (a) ^{72}Ge and (b) ^{131}Xe target for $CE\nu$ NS.	28
Figure 10. Quenching factor measurement results for germanium.	30
Figure 11. A general scheme for detecting a DM particle.	31
Figure 12. The detector scheme of the CDEX-10.	34
Figure 13. The detector scheme of the PandaX-4T.	35
Figure 14. The detector scheme of the XENON program.	36
Figure 15. Feynman diagrams of (a) $E\nu$ ES and (b) $CE\nu$ NS in the presence of light mediators.	39
Figure 16. Feynman diagram representation for the up-scattering for (a) $E\nu$ ES and (b) $CE\nu$ NS.	53
Figure 17. Feynman diagram representation of the charged Higgs contribution to $E\nu$ ES.	57
Figure 18. Efficiency of CDEX-10.	61
Figure 19. Efficiency for PandaX-4T.	61
Figure 20. Efficiency of XENONnT.	62
Figure 21. Differential cross-sections for (a) $E\nu$ ES and (b) $CE\nu$ NS with germanium and xenon targets.	68
Figure 22. Predicted event rates of $E\nu$ ES using (a) germanium and (b) xenon target.	68
Figure 23. Predicted event rates of $CE\nu$ NS using (a) germanium and (b) xenon target.	69
Figure 24. Event rate predictions from CDEX-10 results for (a) $E\nu$ ES and (b) $CE\nu$ NS with contribution from new physics.	69

Figure 25. Event rate predictions from PandaX-4T (a) Run0 and (b) Run1 with contribution from light mediator and transition magnetic moment.	70
Figure 26. Event rate predictions from XENONnT electronic recoil signals with contribution from light mediator and transition magnetic moment are shown.	71
Figure 27. (a) Exclusion limits on scalar mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	73
Figure 28. (a) Exclusion limits on vector mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	73
Figure 29. (a) Exclusion limits on tensor mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	75
Figure 30. (a) Exclusion limits on $U(1)_{L_e-L_\mu}$ mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	75
Figure 31. (a) Exclusion limits on $U(1)_{L_e-L_\tau}$ mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	76
Figure 32. (a) Exclusion limits on $U(1)_{L_\mu-L_\tau}$ mediator using PandaX-4T Run0 and Run1 data with $E\nu$ ES and (b) comparison with previous limits.	76
Figure 33. (a) Exclusion limits on universal light scalar mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	78
Figure 34. (a) Exclusion limits on universal light vector mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	79
Figure 35. (a) Exclusion limits on universal light tensor mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	80
Figure 36. (a) Exclusion limits on $U(1)_{B-L}$ mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	81
Figure 37. (a) Exclusion limits on $U(1)_{B-3L_e}$ mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	83
Figure 38. (a) Exclusion limit on $U(1)_{B-3L_\mu}$ mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	83
Figure 39. (a) Exclusion limits on $U(1)_{B-3L_\tau}$ mediator from CDEX-10 with $CE\nu$ NS and (b) comparison with previous limits.	84
Figure 40. (a) Exclusion limit on $\mu_{\nu_{e4}}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.	85
Figure 41. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.	86
Figure 42. (a) Exclusion limit on $\mu_{\nu_{\tau 4}}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.	87
Figure 43. (a) Exclusion limit on $\mu_{\nu_{\tau 4}}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.	87

Figure 44. (a) Exclusion limit on $\mu_{\nu_{e4}}$ from CDEX-10 with CE ν NS and the projected future results and (b) comparison of the projected results with existing limits.	88
Figure 45. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from CDEX-10 with CE ν NS and the projected future results and (b) comparison of the projected results with existing limits.	89
Figure 46. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from CDEX-10 with CE ν NS and the projected future results and (b) comparison of the projected results with existing limits.	90
Figure 47. (a) Exclusion limit on $\mu_{\nu_{e4}}$ from CDEX-10 with CE ν NS and the projected future results and (b) comparison of the projected results with existing limits.	91
Figure 48. Exclusion limit on the considered parameters from PandaX-4T and XENONnT for (a) $f_{ee}/m_{H^\pm}$, (b) $f_{e\mu}/m_{H^\pm}$, and (c) $f_{e\tau}/m_{H^\pm}$	92
Figure 49. The 90% C.L. allowed regions with 2 d.o.f. in the (a) $f_{e\mu}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ and (b) $f_{e\tau}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ planes.	93
Figure 50. Kinematics for $\nu(p_1) + X(p_2) \rightarrow \nu(p_3) + X(p_4)$ process.	110

LIST OF TABLES

	<u>Page No.</u>
Table 1. The SM matter particle irreducible representations with its quantum numbers.....	2
Table 2. Neutrino sources and their properties.	8
Table 3. Neutrino oscillation parameters from NuFit-5.3.	20
Table 4. Number of the ionizable effective electron charges for xenon and germanium.....	23
Table 5. The universal and flavor-dependent $U(1)'$ vector charges.....	51
Table 6. Significance value and its corresponding degree of freedom.....	63
Table 7. Upper limits with 90% C.L. on g_X for light mediator models from PandaX-4T.....	95
Table 8. Upper limits with 90% C.L. on g_X for light mediator models from CDEX-10.....	95
Table 9. Upper limits with 90% C.L. on the transition magnetic moments from PandaX-4T Run0 and Run1 and previous limits.....	96
Table 10. Upper limits with 90% C.L. on the transition magnetic moments from CDEX-10 and projected limits.....	97
Table 11. Upper limits with 1 d.o.f. on the charged Higgs parameters obtained from analyzing data of XENONnT, PandaX-4T, and their combination. .	98

ABBREVIATIONS

BBN	Big-Bang Nucleosynthesis
BSM	Beyond the Standard Model
C.L.	Confidence Level
CC	Charged Current
CDEX	China Dark Matter Experiment
d.o.f.	Degree of Freedom
$\text{CE}\nu\text{NS}$	Coherent elastic neutrino-nucleus scattering
CJPL	China Jinping Underground Laboratory
CNB	Cosmic neutrino background
CNO	Carbon-Nitrogen-Oxygen
DD	Direct Detection
DM	Dark Matter
$\text{E}\nu\text{ES}$	Elastic neutrino electron scattering
EW	Electroweak
FC	Flavor-conserving
FV	Flavor-violating
GUNS	Grand Unified Neutrino Spectrum
HTM	Higgs Triplet Model
NC	Neutral Current
PandaX	Particle and Astrophysical Xenon Detector
ROI	Region of Interest
SM	Standard Model
SN	Supernova
SSM	Standard Solar Model
TPC	Time Projection Chamber
WIMP	Weakly Interacting Massive Particle

1. INTRODUCTION

The Standard Model (SM) of particle physics has remarkably explained how the universe's basic building blocks interact. It successfully predicts phenomena related to electromagnetic, weak, and strong forces that have been confirmed by experiments at advanced particle physics facilities. However, experimental evidence clearly indicates that the SM is incomplete. Several outstanding puzzles point to the need for physics beyond our current understanding. Among these unresolved issues are the hierarchy problem, the nature of dark matter, the mechanism behind cosmic inflation, and the incorporation of gravity into a quantum framework. Particularly intriguing ones are the nature of neutrinos, their masses, and their interactions, which span from the smallest particle scales to macroscopic phenomena. Consequently, a deeper understanding of neutrino interactions is essential for advancing future research in this area. This chapter provides an introduction to the properties of neutrinos within the SM and discusses their various sources.

1.1. Neutrinos in the Standard Model

Neutrinos are elusive and mysterious particles that are ubiquitous in the universe. While several of their properties are well explained within the framework of the SM, others remain unexplained and point towards physics beyond the SM (BSM). These intriguing properties have been consistently confirmed by experimental observations, leading to robust and widely accepted conclusions. Continued advancements in neutrino physics have driven numerous scientific breakthroughs and remain at the forefront of particle physics research.

The origin of the idea of neutrino was proposed by Pauli (1930) to address the continuous spectrum of β -decays. The discovery of neutrinos necessitated the concept of the weak force. This force was initially described by Fermi (1934) as an interaction involving a weakly interacting fermion with a mass comparable to that of an electron. He was also the one who proposed the name *neutrino* for the elusive particle and considered it to be massless and neutrally charged. Since then, experimental efforts to detect the particle were started until its first

detection in the Savannah River detector from the proton inverse β -decay process (Cowan *et al.*, 1956). It matched the theoretical prediction of the cross-section of the neutrino was predicted to be $\sigma_{\bar{\nu}} \leq 10^{-44} \text{ cm}^2$ with energy $E_{\bar{\nu}} \approx 2 \text{ MeV}$ (Bethe & Peierls, 1934). These achievements were the inspiration for the development of neutrino detectors until now.

The SM tried to incorporate the neutrino properties, yet it has not covered all of its mysteries. A brief review of the paradigm is important to be mentioned. The SM is a product of a series of theoretical developments (Glashow, 1961; Weinberg, 1967; Salam, 1968) in the same era as the development of neutrino physics. It is constructed as a gauge theory based on the local symmetry group $SU(3)_C \times SU(2)_L \times U(1)_Y$, where C stands for color, L for left-handed chirality, and Y for hypercharge. The $SU(3)_C$ describes color charges from Quantum Chromodynamics (QCD), a model that explains the dynamics of quarks together with their binding gluons. The $SU(2)_L \times U(1)_Y$ is the group transformation that is obeyed by all the particles, quarks and leptons, correspond to the electroweak (EW) interactions. This symmetry is spontaneously broken via the Higgs mechanism to $U(1)_Q$, where Q denotes the electric charge. This breaking gives mass to the weak gauge bosons while preserving the gauge invariance associated with electromagnetism. The fermion components of the model is categorized by the corresponding quantum number, as can be seen in the Table 1. Notice

Table 1. The SM particle irreducible representation with its quantum numbers.

Particles	$SU(3)$	$SU(2)_L$	I^3	Y	Q
Leptons					
$\begin{pmatrix} \nu_{eL} \\ e_L \end{pmatrix}, \begin{pmatrix} \nu_{\mu L} \\ \mu_L \end{pmatrix}, \begin{pmatrix} \nu_{\tau L} \\ \tau_L \end{pmatrix}$	1	2	$\begin{pmatrix} +1/2 \\ -1/2 \end{pmatrix}$	-1/2	$\begin{pmatrix} 0 \\ -1 \end{pmatrix}$
e_R, μ_R, τ_R	1	1	0	-1	-1
Quarks					
$\begin{pmatrix} u_L \\ d_L \end{pmatrix}, \begin{pmatrix} c_L \\ s_L \end{pmatrix}, \begin{pmatrix} t_L \\ b_L \end{pmatrix}$	3	2	$\begin{pmatrix} +1/2 \\ -1/2 \end{pmatrix}$	+1/6	$\begin{pmatrix} +2/3 \\ -1/3 \end{pmatrix}$
u_R, c_R, t_R	3	1	0	+2/3	+2/3
d_R, s_R, b_R	3	1	0	-1/3	-1/3

that in the SM, neutrinos are left-handed since they were assumed to be massless particles in the early development. Both quarks and leptons carry hypercharge Y , where they are related with electric charge Q and weak isospin I^3 by $Q = I^3 + Y$.

Neutrinos have three flavors. They are correspond to each lepton types, namely elec-

tron (e^-), muon (μ^-), and tau (τ^-). The ν_e type occurs in the β -decay process. The ν_μ can be found in the pion decay, while the ν_τ from a leptonic-tau decay process. Examples of these two are $\pi^+ \rightarrow \mu^+ + \nu_\mu$ and $\tau^- \rightarrow \nu_\tau + e^- + \bar{\nu}_e$. These processes obey lepton number conservation, namely for lepton $L = 1$, for antilepton $L = -1$, and for non-leptonic $L = 0$. Moreover, there are also number conservation for each generations denoted by L_e, L_μ , and L_τ which can have value of 1 for particle and -1 for antiparticle. Both reactions above and the β -decay for example obey this lepton conservation. The restriction of three generations in the EW theory is determined from Z boson decay. The decay width consists of the summation of visible hadrons, e^-e^+ , $\mu^-\mu^+$, and $\tau^-\tau^+$, and invisible process widths. In the weak interaction, the flavor eigenstates are ν_e, ν_μ, ν_τ , in accordance with charged leptons in their production process by the charge current (CC). Meanwhile, the neutral current (NC) does not play a role in this case. The branching ratio of the invisible final states for SM consideration would be due to the number of the light neutrino species N_ν . Therefore, the branching ratio of the invisible decay would be $\Gamma_{\text{inv}} = N_\nu \Gamma_{\nu\bar{\nu}}$ and from combined LEP observations (Schael *et al.*, 2006), the number of light neutrino species, excluding non-standard possibilities, is $N_\nu = 2.984 \pm 0.0082$. Hence, the SM should contain three neutrino flavors.

Chirality and helicity play important role in neutrino physics. The concept of chirality comes from the fifth gamma-matrix γ^5 in the $V - A$ structure of weak interactions. It gives rise to two possible chiral operators, $P_R = (1 + \gamma_5)/2$ for right projection and $P_L = (1 - \gamma_5)/2$ for left projection. In the weak interactions of the SM, the force carriers interact exclusively with left-handed fermion, $\psi_L = P_L \psi$, and right-handed anti-fermion $\bar{\psi}_R = (P_R \psi)^\dagger \gamma^0 = \bar{\psi} P_L$, with ψ denotes a fermion. In fact, it is a direct implication of the parity violation, which explains the preferred directionality in the spin of emitted particles in the β -decay experiment. Helicity, on the other hand, is the projection of particle spin to its corresponding momentum, i.e. direction of particle's motion. For massless particles, such as neutrinos in the SM, helicity is equivalent to chirality. We can recognize the difference between chirality and helicity from the pion decay process to $\mu^+ \nu_\mu$ or to $e^+ \nu_e$. Spinless pion at initial state sets the decay products to have opposite spins. Since ν_μ and ν_e are left-handed in helicity with the same chirality, μ^+ and e^+ also need to have the left-handed helicity to preserve angular momentum conservation. But this requirement is forbidden as anti-particle must be in right-handed helicity. The process then contains a mixture of left-handed helicity for a right-handed

chirality state (Naumov, 2011). In short, the structure $V - A$ in the weak interaction of the SM only couples left-handed neutrino and right-handed antineutrino as a consequence of the vertex structure correction from the initial proposed model that violates parity.

1.2. Neutrino Oscillation Phenomena

In the SM of particle physics, neutrinos are considered massless and exist in three distinct flavors. However, results from solar and atmospheric neutrino experiments – most notably by Fukuda *et al.* (1998) and (Ahmad *et al.*, 2002) – demonstrated that neutrinos do, in fact, have mass. These experiments revealed that neutrinos can change from one flavor to another as they travel over large distances, a phenomenon known as neutrino oscillation. This idea was first proposed by Pontecorvo (1957), and later developed further by Maki, Nakagawa, & Sakata (1962).

To explain neutrino oscillations, it is necessary to distinguish between flavor eigenstates (which participate in weak interactions) and mass eigenstates (which propagate in space). Neutrino oscillation arises due to a quantum mechanical mixing between these eigenstates, implying that the flavor states are superpositions of the mass eigenstates. The general representation are ν_1 , ν_2 , and ν_3 , which correspond to eigenvalues m_1 , m_2 , and m_3 . To relate it with flavor eigenstate, a unitary 3×3 unitary matrix U is introduced

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = \begin{pmatrix} U_{e1} & U_{e2} & U_{e3} \\ U_{\mu1} & U_{\mu2} & U_{\mu3} \\ U_{\tau1} & U_{\tau2} & U_{\tau3} \end{pmatrix} \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}. \quad (1.1)$$

The matrix is generally shown as U_{PMNS} , where PMNS stands for Pontecorvo-Maki-Nakagawa-Sakata. This expression can be expressed as $|\nu_\alpha\rangle = \sum_i U_{\alpha i} |\nu_i\rangle$, where $\alpha = e, \mu, \tau$ and $i = 1, 2, 3$. Explicitly, the matrix is expressed as (Thompson *et al.*, 2009)

$$U_{\text{PMNS}} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & s_{13}e^{-i\delta} \\ 0 & 1 & 0 \\ -s_{13}e^{-i\delta} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (1.2)$$

In this expression, $c_{ab} = \cos \theta_{ab}$ and $s_{ab} = \sin \theta_{ab}$, where $\theta_{12}, \theta_{13}, \theta_{23}$ denote the mixing

angles between flavors, and δ represents the complex phase of charge-parity (CP)-violation. Using this matrix, the generic formula to account for the probability of neutrino oscillation from flavor α to flavor β is

$$P(\alpha \rightarrow \beta) = \sum_{j,k} U_{\alpha j}^* U_{\beta j} U_{\alpha k} U_{\beta k}^* e^{-i \frac{\Delta m_{jk}^2 L}{2|\vec{p}|}}, \quad (1.3)$$

where $\Delta m_{jk}^2 \equiv m_k^2 - m_j^2$ is the squared mass difference between the state ν_j and ν_k , $U_{\alpha j}$ is the α, j element of the matrix, L represents the distance traveled by neutrino from its origin to a detector or called baseline, while $|\vec{p}|$ denotes the neutrino's momentum. In its observations, neutrinos originating from a source (sun, μ -decay, nuclear reaction, etc.) are produced and then detected after traveling a distance L . During propagation, neutrinos experience a weak interaction process in combination with flavor states and randomly distributed in space-time. It implies a slightly different velocity phase, hence the probability of obtaining a different state between the detected and produced neutrinos is non-zero.

The observation of neutrino oscillations provides strong evidence for a mismatch between the neutrino flavor and mass eigenstates. This phenomenon can only occur if neutrinos have non-zero masses, a feature not accounted for in the formulation of the SM. As such, the existence of massive neutrinos indicates that the SM is incomplete and points toward the need for physics beyond the SM to explain the origin of neutrino masses. A significant amount of model building has been undertaken to incorporate the experimental evidence for neutrino masses and oscillations. Introducing a mass term for neutrinos within the theoretical framework depends on the fundamental nature of the neutrino specifically, whether it is a Dirac or Majorana particle. However, the inclusion of such mass terms is not possible within the SM. This is because the SM lacks right-handed neutrino fields and does not allow lepton number violation, both of which are necessary for constructing viable neutrino mass terms in Majorana case. As a result, extending the SM is essential to accommodate massive neutrinos.

The simplest extension to the SM can be made by adding ν_R to form a Dirac mass term. In the literature it is also associated with a sterile neutrino. The corresponding Yukawa coupling can be expressed as (Granelli, 2024)

$$\mathcal{L}_Y = -\lambda \bar{L} \tilde{H} \nu_R + h.c. \quad (1.4)$$

with the lepton doublet $L = \begin{pmatrix} \nu_L & e_L \end{pmatrix}^T$ and $\tilde{H} \equiv i\sigma_2 H^*$, where H denotes the Higgs field and λ its coupling. This form generates a Dirac mass term $\sim -\frac{v_H}{\sqrt{2}} \bar{\nu}_R \lambda \nu_L + \text{h.c.}$ as the neutral Higgs field acquires a vacuum expectation value ν_H . It leads to a Dirac mass that conserve lepton number. Taking the neutrino mass to be sub-eV scale implies the coupling is approximately $\sim 10^{-12}$. This result is too small and there is no description about the strong mass hierarchy between neutrinos and the charged leptons. Assuming a naive expectation to the hierarchy in the quark sector, this is in contradiction with the experimental results.

Another possible way to describe a massive neutrino is by proposing a Majorana mass term. The interaction term can be written as

$$\mathcal{L}_Y = \frac{\lambda}{\Lambda} \bar{L} \tilde{H} C \tilde{H}^T \bar{L}^T + \text{h.c.}, \quad (1.5)$$

which is widely known as introducing the effective Weinberg operator (Weinberg, 1979). It has dimension five and hence needs a mass scale Λ . It is the only five-dimensional operator that can be built up by the SM content. Adding this term into the SM leads to a non-renormalizable theory. It means that the theory is not valid at any scale, but limited by the scale of Λ . This operator breaks lepton number by two units. As the Higgs boson gets a vacuum expectation value, we obtain a Majorana mass term $\sim \frac{\lambda v_H^2}{2\Lambda} \bar{\nu}_L C \bar{\nu}_L^T + \text{h.c.}$, which is similar to the Fermi theory. Note that the Fermi theory is the low energy model of the weak interactions mediated by the W -boson. A lot of the research in theoretical neutrino physics today focuses on searching for new particles and interactions, focusing on understanding the nature of neutrinos either as a Dirac or a Majorana particle.

1.3. Neutrino Sources

We are constantly surrounded by an immense number of neutrinos, as they are produced by a wide variety of natural and artificial sources. Artificially, neutrinos are generated in nuclear reactors and particle accelerators during controlled interactions. Naturally, neutrino sources include geoneutrinos (emitted from radioactive decay within the Earth), atmospheric neutrinos (produced by cosmic ray interactions in the atmosphere), solar neutrinos (from nuclear fusion in the Sun), supernova neutrinos, and neutrinos from other cosmological

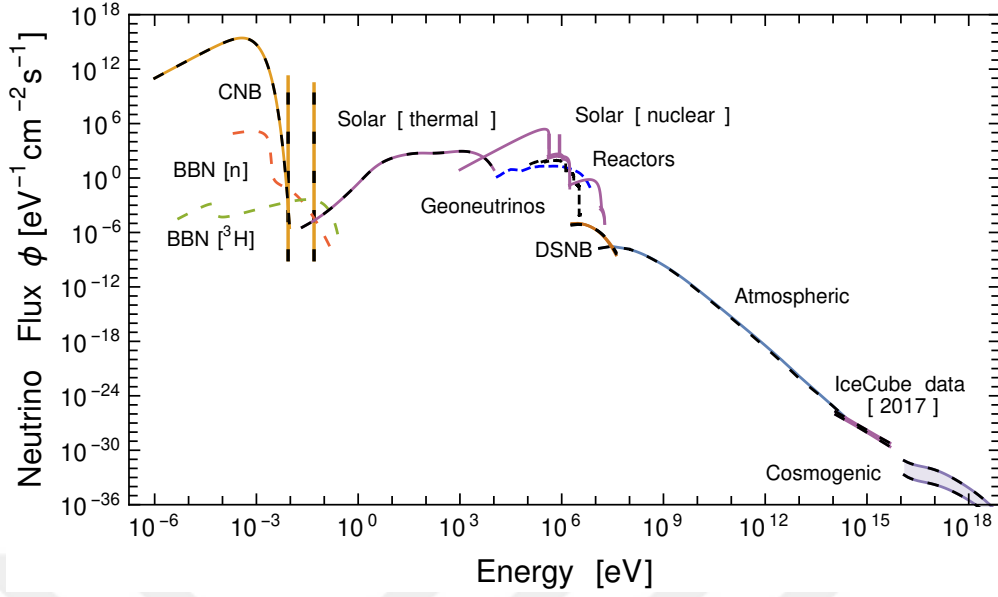


Figure 1. The Grand Unified Neutrino Spectrum at Earth. The plot is reproduced following explanation in (Vitagliano *et al.*, 2020).

and astrophysical events such as gamma-ray bursts and active galactic nuclei. Keeping up with advancements in neutrino astronomy is demanding because of the wide energy range of neutrinos and the diverse sources available, and the detectors involved. It constitutes a fundamental component of the current paradigm of multi-messenger astronomy, in which the investigation of the Universe is no longer restricted to electromagnetic observations but is complemented by cosmic rays, gravitational waves, and neutrinos, which serve as essential astrophysical messengers. The broad span of energy scales, together with the diversity of sources and detector technologies, makes it challenging to remain fully informed about all developments across the field of neutrino astronomy. An approach to addressing this issue is the introduction of the Grand Unified Neutrino Spectrum (GUNS) (Vitagliano *et al.*, 2020). It provides a summary of the neutrino and antineutrino background observed at Earth in a single plot. We present the plot of GUNS in Figure 1, where neutrinos are represented by lines, and antineutrinos by dashed or dotted lines.

Figure 2 shows neutrino cross sections across all currently explored energy scales. The interactions of neutrinos with matter depicted in this figure are well described by the SM of particle physics. It is important to emphasize that neutrinos are electrically neutral and possess extremely small masses. As a result, they interact only weakly with matter and are, for all practical purposes, invisible to our senses. We have to deduce the presence of a

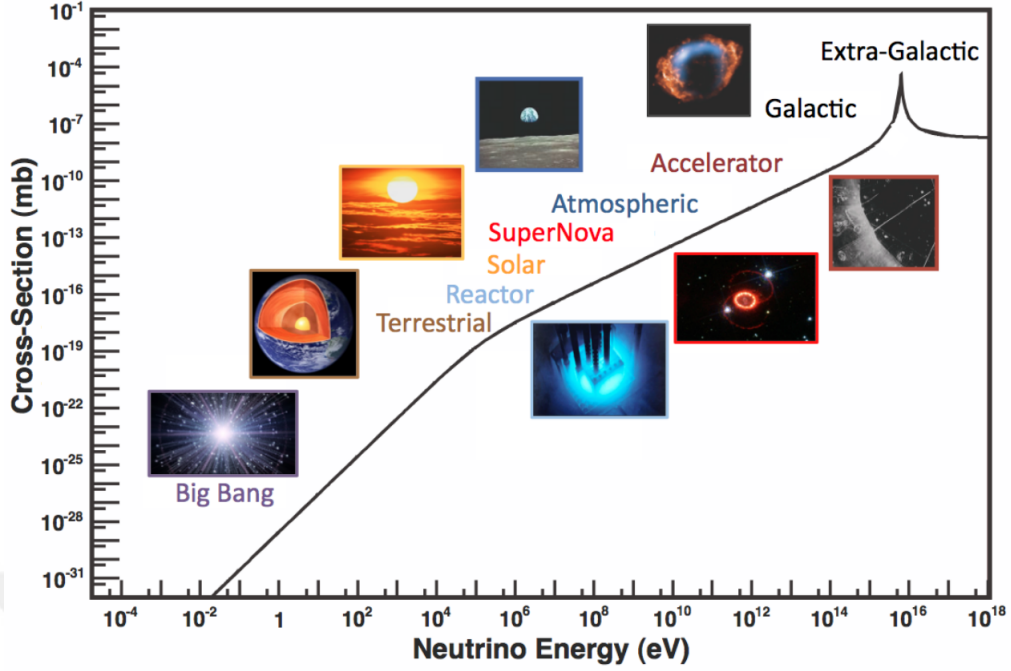


Figure 2. Representative of neutrino cross sections from various sources as a function of energy, adapted from (Formaggio & Zeller, 2012).

neutrino by observing the particles it leaves behind after an interaction. By reconstructing the track of a certain process in an experiment, we can recognize the interaction type by analyzing the kinematics of the final state. In general, due to the small cross section, we need to build a detector with a large dimension and high mass to obtain a large enough event rate to perform useful studies. For completeness, we provide in Table 2 a list of various features of

Table 2. Neutrino sources and their properties. The cross section σ is in $[\text{cm}^{-2}]$.

Note that $\star$ indicates that the value is still not known precisely.

Source	Time	Distance	σ	Events	Ref.
Earth	4 yr	$\star$	$\sim 10^{-44}$	10	(Bellini <i>et al.</i> , 2013)
Reactor	4 yr	10^3 km	$\sim 10^{-42}$	30	(Bellini <i>et al.</i> , 2013)
Accelerator (ν_μ)	1 yr	500 km	$\sim 10^{-37}$	3000	(Aliu <i>et al.</i> , 2005)
Atmosphere	1 yr	500 km	$\sim 10^{-38}$	3000	(Fukuda <i>et al.</i> , 1998)
Sun(pp)	1 day	10^8 km	$\sim 10^{-45}$	100	(Bellini <i>et al.</i> , 2014)
Supernova	1 yr	300 kpc	$\sim 10^{-42}$	0.6	(Vissani & Pagliaroli, 2011)
Galactic	1 yr	1 kpc	$\sim 10^{-35}$	$1\star$	(Halzen <i>et al.</i> , 2008)
High Energy	3yr	$\star$	$\sim 10^{-34}$	10	(Aartsen <i>et al.</i> , 2014)

the neutrino source, including average energies, distances, cross sections, and the predicted events. Brief explanations about the artificial and natural neutrino sources are as follows.

1.3.1. Artificial sources

Reactors

Nuclear reactors produce anti-neutrino from the process of β -decay, arising from unstable fission of ^{238}U or ^{239}Pu . The process produces only electron-antineutrinos. The flux prediction is relatively easy, since it falls off as the inverse square distance of the detector from the reactor. The peak of energy is at about 3 MeV and extend up to around 8 MeV. The process

$$\bar{\nu}_e + p \rightarrow e^+ + n, \quad (1.6)$$

which is the inverse beta decay, with about 1.8 MeV energy threshold is the predominant process to be observed. Generally, cosmic rays interactions in the surrounding material and natural radioactivity are the main backgrounds at nuclear reactors.

Neutrino beamlines

Neutrino beams can be produced in accelerators. Conventional neutrino beams are made following the same way as neutrinos produced from cosmic rays. At first, a synchrotron sends bunches of protons with high energy toward a fixed target. The pion and kaon beams are created by proton collisions with the targets. These produced mesons are focused by aiming them into a magnetic focusing channel. There are secondary decay processes

$$\pi^+ \rightarrow \mu^+ + \nu_\mu, \quad K^+ \rightarrow \mu^+ + \nu_\mu. \quad (1.7)$$

The branching ratio for pion is almost 100% while for kaon is 63.5%. The produced muon can further decay via

$$\mu^+ \rightarrow e^+ + \bar{\nu}_\mu + \nu_e, \quad (1.8)$$

generating muon-antineutrinos and electron-antineutrinos. A long shield at the end of the tunnel absorbs mesons which have not decayed and stopped the charged muons in the beam. Only neutrino beams pass through it. After the shield, the experiment of detecting the neutrinos are placed.

Off-axis beams

The off-axis type of neutrino experiment places the detector at an angle to the axis of the beam. This configuration is developed so that a monochromatic neutrino spectrum can be produced, namely neutrino flux with a single energy. The monochromatic neutrino spectrum is ideally required to study neutrino oscillations. This type of neutrino experiment is currently either being constructed or in the planning stage. It is built on the foundation of intense proton beams.

1.3.2. Natural sources

Relic neutrinos

In the early universe, it is predicted that about 10^{-6} seconds following the big bang all fundamental particles composed a thermal plasma. In this stage, the reactions rates of various particles were considered in thermal equilibrium. When the universe expands, the equilibrium state can not be maintained as temperature drops. The particle component of the plasma decouple or *freeze-out* from the evolution of the universe. At this stage, the scenario has photons, electrons, positrons, and neutrinos in thermal equilibrium. As the temperature falls further, there is not enough energy to create the neutral Z -boson, implying the neutrinos then decouple from the photons and electrons and develop independently. This relic neutrinos that decouple after about 1 second are still around. The current model suggests that their energy is about 0.1 meV and there are around 330 per cubic centimeter. They only interact via the neutral current. Their low energy falls below the electron mass and to detect them is a big challenge. If we could detect them, relic neutrinos would provide extremely valuable information of the early universe, much closer to the Big Bang than our current best probes.

Supernova neutrinos

A massive star explosion is known as supernova. Neutrinos are produced in huge numbers during this violent astrophysical event. For a massive star, the electrons inside its core can have high energy. It is then possible for the electrons to be captured by the protons $e^- + p \rightarrow n + \nu_e$ and heavy nuclei $e + (Z + A) \rightarrow (Z, A - 1) + \nu_e$ within the core. The core begins to contract at supersonic speeds as the number of electron decrease. A significant

number of neutrinos are left behind, which then begin to propagate outward. As the core continues to collapse, reaching the point densities of an atomic nucleus, the pressure from neutron degeneracy and the strong force prevents the collapse. The shrinking matter rebounds and produces a shock wave. The wave then propagates out of the star. It is predicted that about 99% of the supernova's released energy is carried away by neutrinos.

Solar neutrinos

Neutrinos are continuously produced by the nuclear fusion inside the Sun's core. This process fuel the sun and generating electron neutrinos. There are two main processes that responsible during the process, namely the *pp*-chain and the CNO-cycle. More details about the topic are provided in the next section.

Atmospheric neutrinos

Cosmic rays constantly bombard the outer layer of Earth's atmosphere. The major component are protons, with minor components of alpha particles and heavier nuclei and electrons. In the atmosphere, a cascade of hadrons are produced as the primary cosmic rays collide with nuclei. Atmospheric neutrinos are produced as the hadrons decay. The following are some examples of such processes

$$\pi^+ \rightarrow \mu^+ + \nu_\mu, \quad \pi^- \rightarrow \mu^- + \bar{\nu}_\mu \quad (1.9)$$

$$\mu^+ \rightarrow e^+ + \nu_e + \bar{\nu}_\mu, \quad \mu^- \rightarrow e^- + \bar{\nu}_e + \nu_\mu. \quad (1.10)$$

Geoneutrinos

The production of geoneutrinos originates from the unstable radioactive elements that decay inside the Earth. The elements can be uranium, thorium, and potassium. The decay processes also generate heat of the geothermal heat flow. Beyond 10 km from the surface, the precise amount of radioactive material and where it is located remain unclear. Furthermore, recent observations provide the upper limit on the heat generation, indicating the evidence of the significant radioactivity inside the core of the Earth. Measurements of geoneutrinos can help us address these issue.

1.4. Scope and Outline of Dissertation

The primary objective of this dissertation is to explore BSM physics by investigating various models using solar neutrinos at dark matter (DM) – direct detection (DD) experiments. In particular, we investigate the potential signatures and implications of the following BSM scenarios:

- **New Light Mediator Models:** These models are motivated by universal considerations and include $U(1)'$ gauge extensions, where a new light gauge boson mediates neutrino interactions.
- **Sterile Neutrino Dipole Portal:** The scenario involves the transition magnetic moments between active and sterile neutrino states, providing a possible portal to a hidden sector.
- **Triplet Higgs Model:** In this model, contribution from a charged Higgs boson is considered, offering an extended scalar sector that can influence neutrino interactions and mass generation.

Neutrinos from the Sun are detected in Earth-based experiments through their interactions with electrons and nuclei in the detector's target material. These interactions primarily occur via two processes: the elastic neutrino-electron scattering ($E\nu ES$) and the coherent elastic neutrino-nucleus scattering ($CE\nu NS$). Both processes contribute to the neutrino background observed in DM–DD experiments, many of which were originally designed for searching DM candidates and are now being utilized or developed for precision neutrino measurements. We are particularly interested in the recent results from CDEX, PandaX-4T, and XENONnT experiments. As these experiments continue to collect high-quality data, they offer a promising opportunity to probe BSM physics. In particular, they may provide valuable insights into the mechanisms responsible for generating neutrino mass.

In Chapter 2, we describe solar neutrino interactions relevant to DM-DD experiments. We first discuss the solar neutrino spectrum and its origin. We then move on to the SM formulation of the $E\nu ES$ and the $CE\nu NS$. Finally, we provide a brief overview of the experimental setups considered in this study. In Chapter 3, we elaborate on the BSM proposals explored in this thesis, including new light mediator models, the sterile neutrino dipole portal, and the

triplet Higgs model. For each model, we briefly discuss theoretical motivations and connections to solar neutrino interactions in DM-DD experiments. In Chapter 4, we describe the analysis methods employed. Our results and discussions, including predicted event rates and parameter constraints for each model, are presented in Chapter 5. We finally give suggestions in Chapter 6.



2. SOLAR NEUTRINO INTERACTIONS AT DIRECT DETECTION EXPERIMENTS

The search for DM remains one of the most compelling and actively pursued challenges in fundamental physics. This pursuit is motivated by robust evidence for gravitational interactions between ordinary baryonic matter and an elusive non-luminous component, observed across astronomical and cosmological scales. Despite extensive research, the fundamental nature of DM remains unresolved. It is widely hypothesized that DM is composed of yet-undiscovered elementary particles, whose detection would profoundly advance our understanding of the universe. If these hypothetical particles are stable, electrically neutral, and possess masses ranging from sub-GeV/ c^2 up to a few TeV/ c^2 , they could be eligible as weakly interacting massive particles (WIMPs) (Arcadi *et al.*, 2018). Such particles, potentially responsible for DM, are expected in numerous proposals of the SM extensions.

Over the past decades, numerous experimental efforts have been undertaken to search for these massive particles through their scattering in detector media. The experiments to conduct this observation are generally located underground. These studies aim to determine the mass of DM particles and their interaction cross-section with ordinary matter. Unveiling the essence of DM would resolve the one profound open question in particle physics and significantly advance our understanding of the Universe and its evolution.

DM-DD experiments typically aim to observe nuclear recoils resulting from DM-nucleus scattering. Modern experiments have achieved sufficiently low energy thresholds and large exposures, making them sensitive to solar neutrinos. In fact, solar neutrinos constitute a significant background when probing very small DM-nucleon scattering cross sections, a limit commonly cited as the “neutrino floor” (O’Hare, 2021). The feasible impact of solar neutrinos on DD experiments was first recognized over 40 years ago (Cabrera *et al.*, 1984), and has been extensively studied since.

In this chapter, we first explain the solar neutrino spectrum. It includes the production process of neutrinos and the propagation processes to the Earth. Next, we explain $E\nu$ ES and $CE\nu$ NS, which can be triggered by solar neutrino interactions at DM-DD experiments. We

then describe the concept of such experiments and discuss recent developments accordingly. We close this chapter by discussing CDEX-10, PandaX-4T, and XENONnT collaborations, from which we utilize their data in our study.

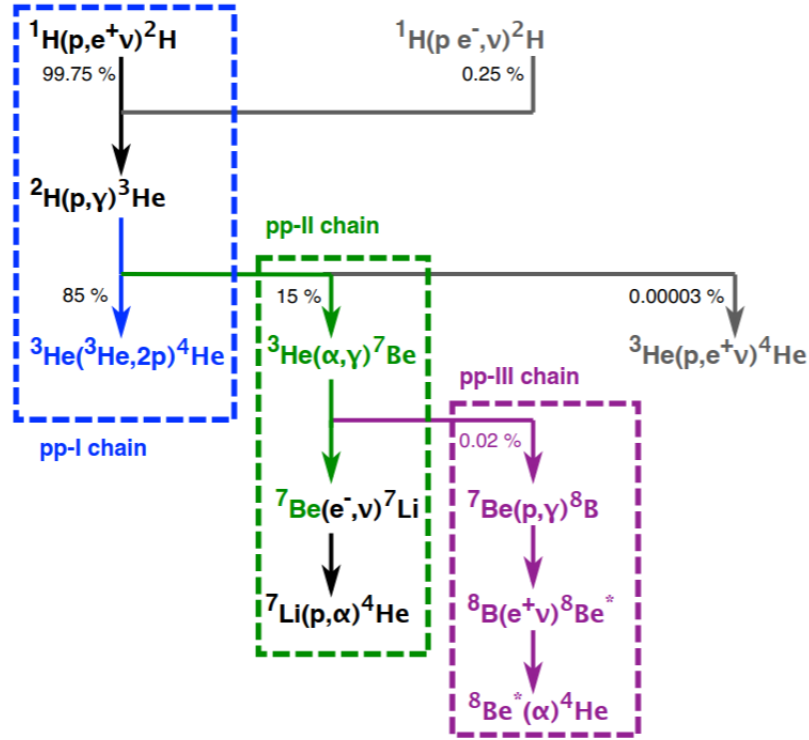
2.1. The Solar Neutrino Spectrum

Stellar nucleosynthesis explains that the Sun produces an immense amount of neutrinos from the process of nuclear fusion inside the core. The first solar neutrino experiment at Homestake (Davis *et al.*, 1968) obtained a precise measurement of $2.56 \pm 0.16 \pm 0.16$ SNU (Cleveland *et al.*, 1998) where $1 \text{ SNU} = 10^{-36} \text{ events/atom/sec}$. This value indicates a deficit of solar neutrinos by one-third compared to theoretical predictions. Similar findings were confirmed by other experiments at GALLEX (Anselmann, P. *et al.*), SAGE (Abdurashitov *et al.*, 2002), Kamiokande (Hirata *et al.*, 1989), and SNO (Ahmad *et al.*, 2002). This discrepancy between theory and experimental results became the so-called “solar neutrino problem.” Many theoretical scenarios were proposed to explain the discrepancy. At last, the standard solution of this solar neutrino problem is neutrino oscillation with the Mikheyev-Smirnov-Wolfenstein (MSW) (Mikheev & Smirnov, 1986; Wolfenstein, 1978) effect and a large mixing angle (LMA).

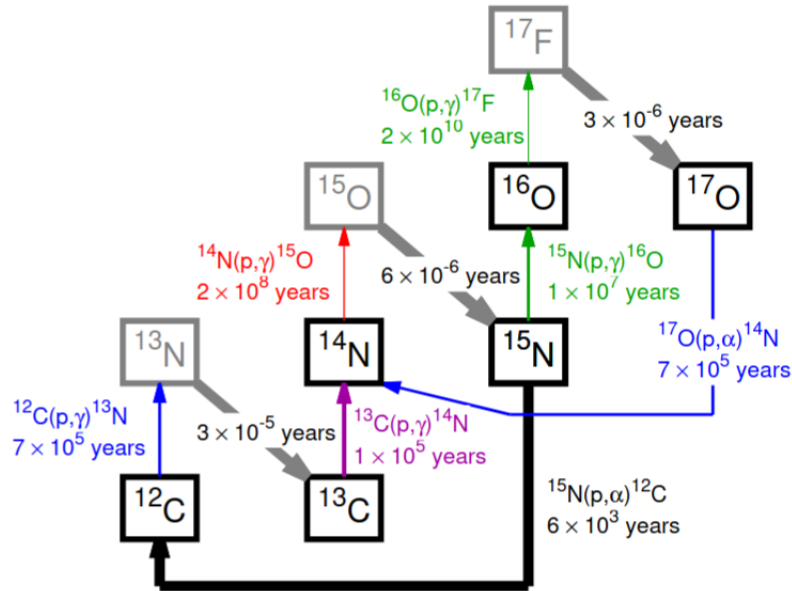
Solar neutrino production is described by the Standard Solar Model (SSM) (Vinyoles *et al.*, 2017), a comprehensive study essential for understanding the Sun’s structure and behavior. The SSM is based on some fundamental principles: energy transport, hydrostatic equilibrium, and the nuclear reactions that generate the energy of the Sun. By combining these principles with observational constraints – like the solar radius, luminosity, age, and radiative opacity – a detailed forecast of the Sun’s internal structure, such as temperature, density, pressure, as well as neutrino fluxes, can be explored. Solar neutrinos, in particular, are produced in the process of nuclear fusion, where protons turn into helium nuclei. They represent the primary natural source of neutrino interaction events in DM-DD experiments, contributing to both electron and nuclear recoil signals.

Hydrogen burning is relevant at the temperature of the solar core, and the proton-proton (pp) chains dominate. The energy production in this process is dominated by three pp chains. These are termed pp-I, pp-II, and pp-III. The second process is the carbon-nitrogen-oxygen

(CNO) cycle. The two main processes are the CN- and NO-cycle. These cycles produce the CNO neutrino fluxes. The diagram of both processes can be seen in Fig. 3, adapted from (Gaan, 2021). For a simple representation of the nuclear reactions, the following notation is



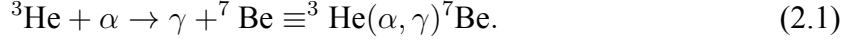
(a)



(b)

Figure 3. The pp-chains (a) and the CNO cycle (b) processes.

commonly used:



Other processes with this notation are α decay (α), electron capture (e^- , ν_e), and positron capture (e^+ , ν_e). This form helps us understand the nuclear process of producing neutrinos.

We can see the process of pp-chains in Figure 3.(a). The initial process ${}^1\text{H}(p, e^+\nu_e){}^2\text{H}$ controls the rate of all three pp-chains. The cross-section is too low to be accessible experimentally. Next, the ${}^2\text{H}(p, \gamma){}^3\text{He}$ process occurs much faster. In between the pp-I phase and pp-II phase, there are two processes: ${}^3\text{He}({}^3\text{He}, 2p){}^4\text{He}$ (83%) and ${}^3\text{He}(\alpha, \gamma){}^7\text{Be}$ (17%). In the pp-III chain, there is a competition between the electron-capture decay of ${}^7\text{Be}$ and the ${}^7\text{Be}(p, \gamma){}^8\text{B}$ reaction. Additionally, two secondary nuclear reactions branch off from the three primary pp chains: ${}^1\text{H}(pe^-, \nu_e){}^2\text{H}$ (pep) and ${}^3\text{He}(p, e^+\nu_e){}^4\text{He}$ (hep).

Next, we can see the process of CNO-cycles in Figure 3 (b). The process begins from the preexisting ${}^{12}\text{C}$ inside the core. The overall timescale of the cycle is governed by ${}^{14}\text{N}(p, \gamma){}^{15}\text{O}$, the slowest reaction. Two additional reactions, ${}^{12}\text{C}(p, \gamma){}^{13}\text{N}$ and ${}^{16}\text{O}(p, \gamma){}^{17}\text{F}$, play significant roles in CNO burning, as they influence the initiation of the process. The CN-cycle and NO-cycle intersect at ${}^{15}\text{N}$, where the reactions ${}^{15}\text{N}(p, \alpha){}^{12}\text{C}$ and ${}^{15}\text{N}(p, \gamma){}^{16}\text{O}$ take place. The preexisting ${}^{16}\text{O}$ participates in the ${}^{16}\text{O}(p, \gamma){}^{17}\text{F}$ reaction, and the resulting ${}^{17}\text{F}$ is rapidly converted back to ${}^{14}\text{N}$ via ${}^{17}\text{O}(p, \alpha){}^{14}\text{N}$ reaction. These reactions continue to be investigated in order to enhance our knowledge about nuclear processes within the Sun.

We provide the typical solar neutrino flux in Figure 4, normalized to B16-GS98 SSM (Vinyoles *et al.*, 2017). We have discussed above the two primary nuclear fusion processes relevant for producing solar neutrinos. The pp chain is the dominant source of neutrino flux that can induce interactions in direct detection experiments. In contrast, the CNO cycle contributes minimally to the overall energy output from the Sun, but is crucial for probing its composition. This cycle generates rather small solar neutrino fluxes.

The current technology of DD experiments is sensitive to the process of $\text{E}\nu\text{ES}$ and $\text{CE}\nu\text{NS}$. The $\text{E}\nu\text{ES}$ process is induced by pp and ${}^7\text{Be}$ solar neutrinos, while the $\text{CE}\nu\text{NS}$ process is mostly induced by the ${}^8\text{B}$ solar neutrino. The pp solar neutrinos are produced during the fusion reaction of two protons into helium through deuterium, resulting in a substantial

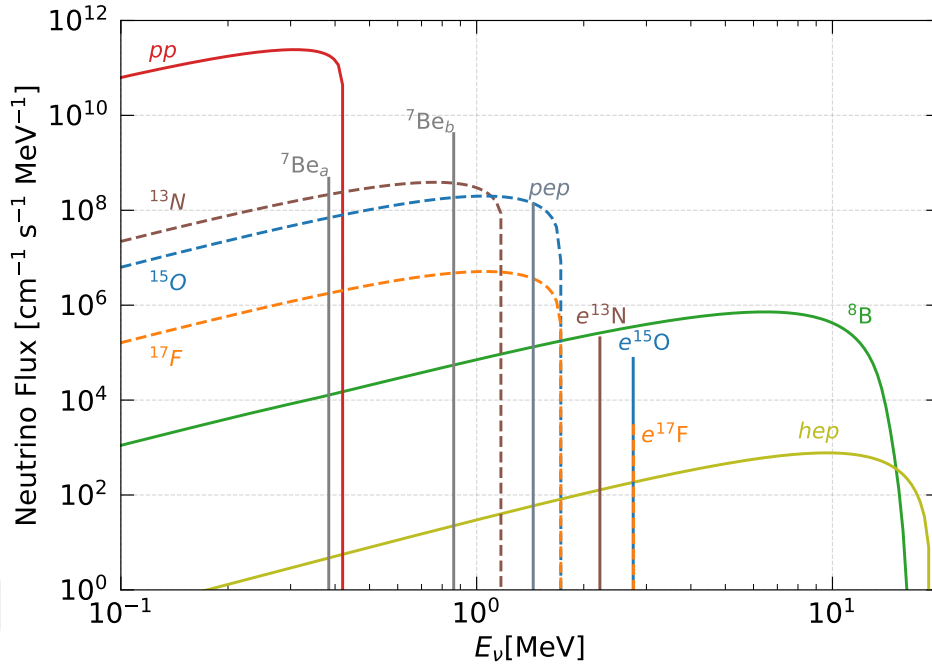


Figure 4. Spectrum of solar neutrino fluxes observed on Earth.

flux of pp solar neutrinos with energies less than 0.3 MeV. The solar neutrinos from ${}^7\text{Be}$ arise from an intermediate step in helium fusion, involving electron capture to form lithium. These neutrinos appear as two discrete lines at 0.862 MeV and 0.384 MeV. The branching ratio of the two channels is 90:100. Neither produce a continuous spectrum because the insufficient nuclear energy to allow β^+ decay. The ${}^8\text{B}$ solar neutrino arises from the decay ${}^8\text{B} \rightarrow {}^7\text{Be}^\square + e^+ + \nu_e$. It occurs in the final step of the pp chain with sizable flux at higher energies. The final channel of the pp cycle generates the hep solar neutrino flux, which has the highest energy but an extremely low intensity, rendering it unobserved. Since the $E\nu\text{ES}$ cross section depends on neutrino flavor, the flavor composition of the flux reaching Earth must be taken into account. In contrast, $\text{CE}\nu\text{NS}$ is flavor-independent because it occurs via neutral currents.

Solar neutrinos experience oscillations as they travel to Earth, arriving at detectors as an admixture of electron-, muon-, and tau-flavor. To incorporate these oscillations, the differential cross section, in terms of recoil energy T_r , must be multiplied by the corresponding

neutrino survival probabilities

$$\frac{d\sigma_{\text{total}}}{dT_r} = P_{ee} \frac{d\sigma_{\nu_e}}{dT_r} + \sum_{f=\mu,\tau} P_{ef} \frac{d\sigma_{\nu_f}}{dT_r}, \quad (2.2)$$

where $P_{e\mu} = (1 - P_{ee}) \cos^2 \vartheta_{23}$ and $P_{e\tau} = (1 - P_{ee}) \sin^2 \vartheta_{23}$. The electron-neutrino case is represented by P_{ee} . We consider (Maltoni & Smirnov, 2016)

$$P_{ee} = (c_{13}^2 c_{13}^{m^2}) \left(\frac{1}{2} - \frac{1}{2} \cos(2\vartheta_{12}^m) \cos(2\vartheta_{12}) \right) + (s_{13}^2 s_{13}^{m^2}). \quad (2.3)$$

In this equation, $c_{13} = \cos \vartheta_{13}$, $s_{13} = \sin \vartheta_{13}$ and the matter effect is represented by the index m . The $\cos(2\vartheta_{12}^m)$ represents the matter angle, with the explicit form

$$\cos(2\vartheta_{12}^m) \approx \frac{\cos(2\vartheta_{12}) - \beta_{12}}{\sqrt{(\cos 2\vartheta_{12} - \beta_{12})^2 + \sin^2 2\vartheta_{12}}}. \quad (2.4)$$

Here, $(s_{13}^m)^2 \approx s_{13}^2(1 + 2\beta_{13})$, where

$$\beta_{12} = \frac{2c_{13}^2 V_e^0 E_\nu}{\Delta m_{21}^2}, \quad \beta_{13} = \frac{2V_e^0 E_\nu}{\Delta m_{31}^2}. \quad (2.5)$$

In these relations, V_e^0 is the value of V_e at the core where electron-neutrino is produced, while the index m represents the matter effect. Notice that $V_e = \sqrt{2}G_F n_e$ is the MSW effective potential induced by coherent forward scattering of neutrinos with electrons in matter, with n_e denoting the electron number density. In calculating the survival probabilities, we account for the day-night asymmetry arising from the Earth's matter effects. After averaging out oscillating parts, the difference of P_{ee} due to this effect can be expressed as

$$\Delta_P = P_{ee}^{\text{day}} - P_{ee}^{\text{night}} \approx \frac{1}{2} c_{13}^6 \frac{\cos 2\theta_{12}^m \sin^2 2\theta_{12} K V_\oplus}{K^2 - 2c_{13}^2 \cos 2\theta_{12} V_\oplus K + V_\oplus^2}. \quad (2.6)$$

In this relation, $K = \delta m_{21}^2 / (2E_\nu)$ and $V_\oplus$ represents the average value of V_e in the Earth. For neutrino oscillation parameters, we consider the normal-ordering without the Super-Kamiokande atmospheric data from the latest results of 3- ν oscillation from NuFit-5.3 (Esteban *et al.*, 2020), summarized in Table 3. Accordingly, we plot the survival probabilities in Figure 5. The day and night asymmetry effect can be inferred for high neutrino energy.

Table 3. Neutrino oscillation parameters from NuFit-5.3. We consider the parameters without SK atmospheric data, obtained from (Esteban *et al.*, 2020).

Parameters	Best fit $\pm 1 \sigma$
$\sin^2 \theta_{12}$	$0.304^{+0.013}_{-0.012}$
$\theta_{12}/^\circ$	$33.44^{+0.78}_{-0.75}$
$\sin^2 \theta_{23}$	$0.570^{+0.018}_{-0.024}$
$\theta_{23}/^\circ$	$49.00^{+1.1}_{-1.4}$
$\sin^2 \theta_{13}$	$0.02221^{+0.00068}_{-0.00062}$
$\theta_{13}/^\circ$	$8.57^{+0.13}_{-0.12}$
$\delta_{CP}/^\circ$	195^{+51}_{-25}
$\frac{\Delta m_{21}^2}{10^{-5} \text{eV}^2}$	$7.42^{+0.21}_{-0.20}$
$\frac{\Delta m_{3\ell}^2}{10^{-5} \text{eV}^2}$	$2.514^{+0.028}_{-0.027}$

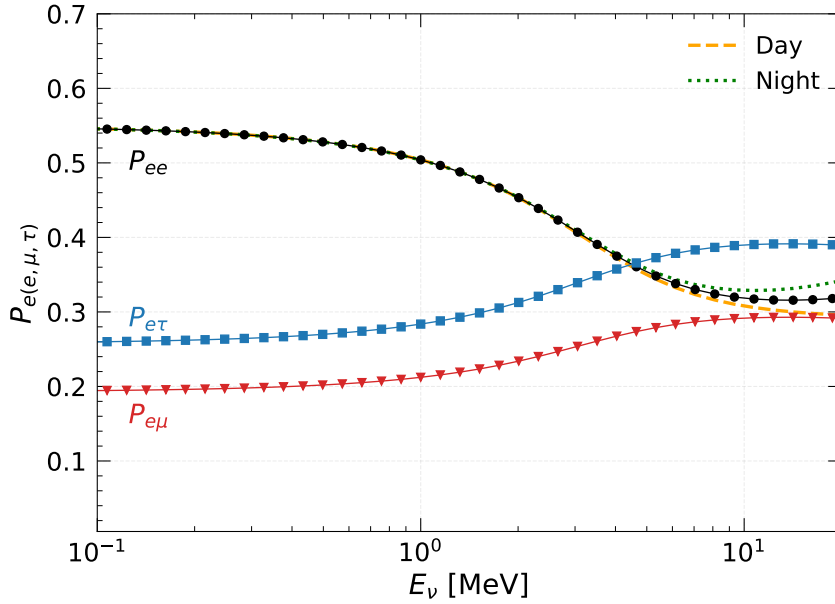


Figure 5. Solar neutrinos survival probabilities.

2.2. Elastic Neutrino-Electron Scattering ($\text{E}\nu\text{ES}$)

The process of $\text{E}\nu\text{ES}$ occurs when a neutrino interacts with an electron via the exchange of either a charged or a neutral gauge boson ('t Hooft, 1971). It is a purely leptonic SM process. At the tree level, we present the Feynman diagrams of the process in Fig. 6. In the diagrams, Z denotes the neutral-gauge boson and W the charged-gauge boson of the SM. The amplitudes of $\nu(p_1) + e^-(p_2) \rightarrow \nu(p_3) + e^-(p_4)$ process for the CC and NC are

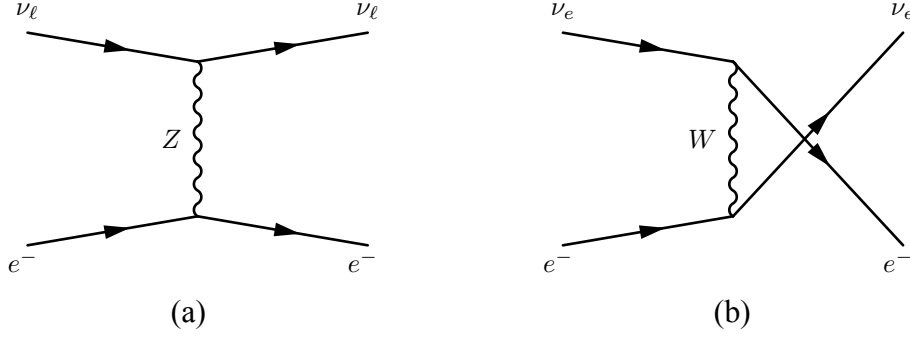


Figure 6. Feynman diagrams representations at tree level for the SM $E\nu$ ES via (a) neutral-current (NC) and (b) charge-current (CC) channels.

$$\mathcal{M}_W = \frac{g^2}{2m_W^2} [\bar{\nu}(p_3)\gamma^\mu P_L \nu(p_1)] [\bar{e}(p_4)\gamma_\mu P_L e(p_2)], \quad (2.7)$$

$$\mathcal{M}_Z = \frac{g^2}{4c_W^2 m_Z^2} [\bar{\nu}(p_3)\gamma^\mu P_L \nu(p_1)] [\bar{e}(p_4)\gamma_\mu (g_V^e - g_A^e \gamma^5) e(p_2)]. \quad (2.8)$$

Here, m_W is the charged W -mass and m_Z the neutral Z -mass, with $m_W = m_Z c_W$. For simplicity, we have used $c_W \equiv \cos \theta_W$, and similarly $s_W \equiv \sin \theta_W$. These vector and axial-vector couplings are

$$g_V^e = -1/2 + 2s_W^2 + \delta_{\ell e} \quad g_A^e = -1/2 + \delta_{\ell e}. \quad (2.9)$$

We note that the Kronecker delta symbol $\delta_{\ell e}$ in the above relations stands for the NC and CC nature of the process. The two currents are relevant for the case of ν_e , while only the NC contributes for the ν_μ and ν_τ cases. In this work, we use $s_W^2 = 0.23857$ in the $\overline{\text{MS}}$ scheme (Navas *et al.*, 2024). The total amplitude of the process is given by adding the two contributions $\mathcal{M}_Z - \mathcal{M}_W$, where the minus sign is necessary to take into account the crossing symmetry. We find

$$\mathcal{M}_{\nu e} = \frac{g^2}{8m_W^2} [\bar{\nu}(p_3)\gamma^\mu (1 - \gamma_5) \nu(p_1)] [\bar{e}(p_4)\gamma_\mu ((g_V^e + 1) - (g_A^e + 1)\gamma_5) e(p_2)]. \quad (2.10)$$

By averaging the absolute square of this amplitude, we find

$$\begin{aligned} \langle |\mathcal{M}_{\nu e}|^2 \rangle = & 16G_F^2 \left[((g_V^e + 1) + (g_A^e + 1))^2 (p_1 \cdot p_2)(p_3 \cdot p_4) \right. \\ & + ((g_V^e + 1) - (g_A^e + 1))^2 (p_1 \cdot p_4)(p_2 \cdot p_3) \\ & \left. - m_e^2 ((g_V^e + 1)^2 - (g_A^e + 1)^2) (p_1 \cdot p_3) \right], \end{aligned} \quad (2.11)$$

where $G_F/\sqrt{2} = g^2/(8m_W^2)$ is the Fermi coupling constant. The four-momentum of the incoming and outgoing particles is denoted by p_i for $i = 1, 2, 3, 4$. The four-vector multiplication can be calculated using the kinematics of the laboratory frame in Equation (A.8), as we discuss in the Appendix. The differential cross section as a function of the electron recoil energy T_{er} can be expressed as

$$\left[\frac{d\sigma_{\nu e}}{dT_{er}} \right]_{\text{SM}}^{\nu e} = Z_{\text{eff}}(T_{er}) \frac{G_F^2 m_e}{2\pi} \left[(g_V^e + g_A^e)^2 + (g_V^e - g_A^e)^2 \left(1 - \frac{T_{er}}{E_\nu} \right)^2 - (g_V^e{}^2 - g_A^e{}^2) \frac{m_e T_{er}}{E_\nu^2} \right]. \quad (2.12)$$

In the above formula, m_e represents the electron mass and E_ν denotes the initial neutrino energy. Note that for the antineutrino-electron scattering, the relative sign between g_V^e and g_A^e at the first and second terms is reversed in the above equation. These two couplings depend on the incoming neutrino's flavor. Meanwhile, the factor $Z_{\text{eff}}(T_{er})$ accounts for the number of the effective ionizable electron charges. This multiplication factor is required, since when the electron recoil energy is comparable to the electron binding energy, about the order of a few keV, the approximation of a free electron typically breaks down because of the atomic ionization effects (Chen *et al.*, 2017). Note that the free-electron approximation assumes electrons to be free and stationary. This factor of the number of charge electrons satisfies

$$Z_{\text{eff}}(T_{er}) = \sum_{\alpha} n_{\alpha} \theta(T_{er} - B_{\alpha}) \quad (2.13)$$

where the electron number is denoted by n_{α} and the binding energy in the corresponding atomic shell is denoted by B_{α} . The term $\theta(T_e - B_{\alpha})$ is the Heaviside function. This factor is tabulated in Table 4 for xenon and germanium. The correction is required for adjusting the cross section obtained using the free-electron approximation.

A notable characteristic of neutrino-electron scattering is its strong directional depen-

Table 4. Number of the ionizable effective electron charges for xenon and germanium (Thompson *et al.*, 2009).

Xenon		Germanium	
$Z_{\text{eff}}(T_{er})$	T_{er} in keV	$Z_{\text{eff}}(T_{er})$	T_{er} in keV
54	$T_{er} > 34.561$	32	$T_{er} > 11.103$
52	$34.561 \geq T_{er} > 5.4528$	30	$11.103 \geq T_{er} > 1.4146$
50	$5.4528 \geq T_{er} > 5.1037$	28	$1.4146 \geq T_{er} > 1.2481$
48	$5.1037 \geq T_{er} > 4.7822$	26	$1.2481 \geq T_{er} > 1.217$
44	$4.7822 \geq T_{er} > 1.1487$	22	$1.217 \geq T_{er} > 1.801$
42	$1.1487 \geq T_{er} > 1.0021$	20	$1.801 \geq T_{er} > 1.249$
40	$1.0021 \geq T_{er} > 0.9406$	18	$1.249 \geq T_{er} > 1.208$
36	$0.9406 \geq T_{er} > 0.6890$	14	$1.208 \geq T_{er} > 0.0298$
32	$0.6890 \geq T_{er} > 0.6764$	10	$0.0298 \geq T_{er} > 0.0292$
26	$0.6764 \geq T_{er} > 0.2132$	4	$T_{er} \leq 0.0292$
24	$0.2132 \geq T_{er} > 0.1467$		
22	$0.1467 \geq T_{er} > 0.1455$		
18	$0.1455 \geq T_{er} > 0.0695$		
14	$0.0695 \geq T_{er} > 0.0675$		
10	$0.0675 \geq T_{er} > 0.0233$		
4	$0.0233 \geq T_{er} > 0.0134$		
2	$0.0134 \geq T_{er} > 0.0121$		
0	$0.0121 \leq T_{er}$		

dence, meaning the outgoing electron is emitted at a very small angle relative to the incoming neutrino's direction (Formaggio & Zeller, 2012). Using the laboratory frame, the angle of the outgoing electron with energy E'_e relative to the neutrino beam is

$$\cos \theta_e = \frac{E'_\nu E'_e - m_e^2 - m_e(E'_\nu - E'_e)}{E'_\nu \sqrt{E_e'^2 - m_e^2}}. \quad (2.14)$$

With small angle approximation and $m_e \ll E'_e$, we find

$$E'_e \theta^2 = 2m_e \left(1 - \frac{E'_e}{E'_\nu} \right). \quad (2.15)$$

It gives us an important constraint $E'_e \theta^2 < 2m_e$, indicating that the scattered electron is always into the forward direction. For the process triggered by solar neutrino, the electron scattered in the forward hemisphere. This distinctive property has been exploited in numerous neutrino experiments, particularly for the detection of solar neutrinos, such as Super-Kamiokande (Hirata *et al.*, 1987; Fukuda *et al.*, 1998), SNO (Ahmad *et al.*, 2002; Ahmed *et al.*, 2004) and BOREXINO (Bellini *et al.*, 2010) experiments. As mentioned earlier, Solar

neutrinos can produce elastic neutrino-electron scattering events in DM-DD experiments by interacting with the electrons in the target material. $E\nu$ ES, along with the $CE\nu$ NS discussed next, constitutes a significant source of neutrino background in these experiments. Principally, the signal can be reduced. However, it is difficult in practice to eliminate entirely in experiments. Consequently, interactions involving neutrinos exhibit a strong degeneracy with the corresponding DM signals. making it a valuable framework for probing BSM signatures.

2.3. Coherent Elastic Neutrino-Nucleus Scattering ($CE\nu$ NS)

The $CE\nu$ NS process can be depicted as neutrinos interact with a nucleus coherently via neutral boson exchange in low energy (Freedman, 1974). Coherent here means that neutrinos are capable to reach the nucleus structure without breaking it. It indicates that nuclear form factors can effectively capture the complexity of the strong interaction. Through this procedure, properties of neutrinos when they mingle with matters in low momentum transfer can be investigated. The corresponding Feynman diagram is shown in Figure 7. It shows a

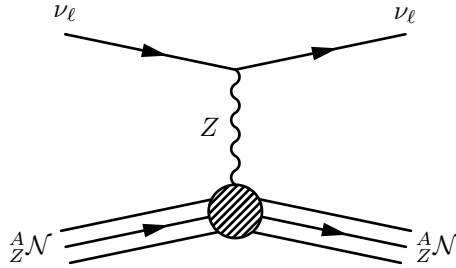


Figure 7. Feynmann diagram of $CE\nu$ NS in the SM.

neutrino interacting with the nucleus as a whole by exchanging neutral boson. So far, $CE\nu$ NS is a process with the largest cross-section involving neutrinos as shown in Figure 8. A relatively large cross-section is possessed by $CE\nu$ NS comparing to other low-energy processes involving neutrino. The wide range of neutrino sources suitable for detection – including terrestrial, solar, atmospheric, and artificial sources such as accelerators and reactors – is likely to stimulate significant scientific endeavors in the years to come. The challenge in detecting this process arises from the nuclear recoil energies, which are typically on the scale of a few keV. This difficult goal has ultimately become achievable due to advancements in

accelerator-based neutrino technology in recent years by COHERENT experiment (Akimov *et al.*, 2017) after its first theory was proposed. The collaboration detected the signal with neutrino energies less than 10 MeV. They utilized heavy target such as mercury which is bombarded by high-energy proton. It then produces a pion-stopped source for fabricating neutrinos (π DAR process) with three types: electron neutrino, muon neutrino, and muon antineutrino. The full coherency was predicted to be achieved from reactors or solar neutrino experiments (Kerman *et al.*, 2016). During the preparation of this thesis, CONUS Collaboration has announced their result of CE ν NS observation from neutrino detector (Bonet *et al.*, 2020). Detection of CE ν NS offer a new perspective of understanding other phenomena in nature.

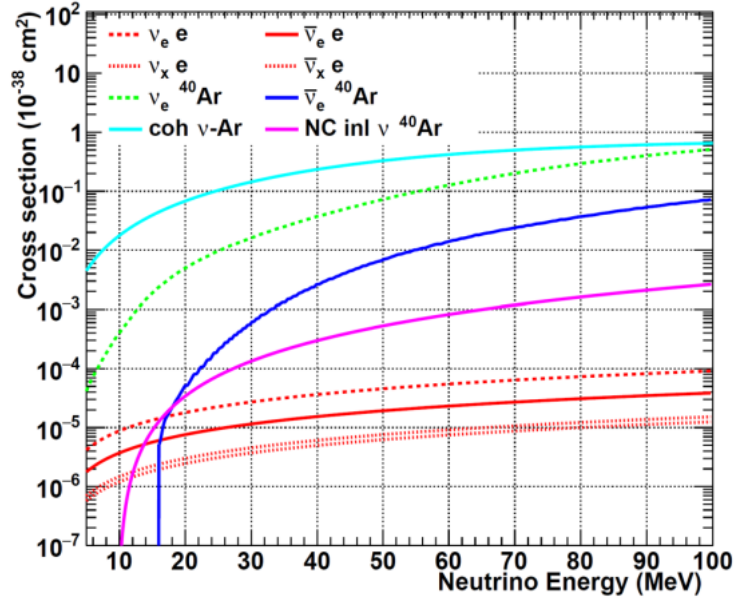


Figure 8. Various processes involving neutrino, adapted from (Sinev, 2020).

The scattering amplitude for the process of $\nu(p_1) + N(p_2) \rightarrow \nu(p_3) + N(p_4)$, where assuming the nucleus as a Dirac particle

$$\langle N(p_4) | J_{EM}^\mu | N(p_2) \rangle = [\mathcal{F} g_V^p + \mathcal{N} g_V^n] \bar{u}(p_4) \gamma^\mu u(p_2) F(q^2), \quad (2.16)$$

can be written as

$$\mathcal{M}_{\nu N} = \frac{G_F}{\sqrt{2}} Q_{SM} F(q^2) g_L^\nu \bar{\nu}(p_3) \gamma^\mu (1 - \gamma^5) \nu(p_1) \bar{u}(p_4) \gamma_\mu u(p_2), \quad (2.17)$$

with the weak charge defined by (Scholberg, 2005)

$$Q_{SM} = \mathcal{N} - \mathcal{Z}(1 - 4s\mathbf{w}^2). \quad (2.18)$$

The average square of this amplitude can be written as

$$\langle |\mathcal{M}_{\nu\mathcal{N}}|^2 \rangle = \sum_{ss'} |\mathcal{M}|^2 = \frac{1}{2} G_F^2 Q_{SM}^2 |F(q^2)|^2 (g_L^\nu)^2 L_\nu^{\mu\nu} W_{\mu\nu}^N \quad (2.19)$$

where we note that neutrino appears only in the left-handed state so that no need for spin averaging factor. We have defined the neutrino tensor $L_\nu^{\mu\nu}$ and nucleus tensor $W_{\mu\nu}^N$. Using the lab frame, the four-vector momentum contraction gives

$$\begin{aligned} L_\nu^{\mu\nu} W_{\mu\nu}^N &= 64 [m_N^2 (E_\nu - T_{nr})^2 + E_\nu^2 m_N^2 - m_N^3 T_{nr}] \\ &= 128 m_N^2 E_\nu^2 \left[1 - \frac{T_{nr}}{E_\nu} - \frac{m_N T_{nr}}{2E_\nu^2} + \frac{T_{nr}^2}{2E_\nu^2} \right]. \end{aligned} \quad (2.20)$$

We notice that only the last term distinguishes this result from the spinless nucleus. The cross section then becomes

$$\left[\frac{d\sigma}{dT_{nr}} \right]_{SM}^{\nu\mathcal{N}} = \frac{G_F^2 Q_{SM}^2 m_N |F(q^2)|^2}{4\pi} \left(1 - \frac{m_N T_{nr}}{2E_\nu^2} \right), \quad (2.21)$$

where terms of order T_{nr}/E_ν have been neglected due to typical initial neutrino energies $E_\nu \lesssim 50$ MeV.

In the differential cross-section of Equation (2.21), the form factor determines inner structure of nucleus target. There are many possibilities for choosing it. We consider in this study the Helm parameterization (Helm, 1956)

$$F_{\text{Helm}}(q^2) = 3 \frac{j_1(qR)}{qR} e^{-\frac{1}{2}q^2 s^2}. \quad (2.22)$$

Here, $j_1(x) = \sin x/x^2 - \cos x/x$ denotes the first order Spherical Bessel function. Parameter $R = \sqrt{c^2 + \frac{7}{3}\pi^2 - 5s^2}$ denotes the nuclear radius, with $c = (1.23\mathcal{A}^{1/3} - 6)$ fm and $s = 0.9$ fm (Duda *et al.*, 2007). Notice that $\mathcal{A}$ is the nucleus mass number. The form factor for proton and neutron are considered equal. Note that relation of the momentum transfer obeys $Q^2 = -q^2 = 2MT_{nr}$.

In its development, interference terms can arise between vector and axial-vector interactions. Such contributions become significant in precision studies of nuclei with spin-dependent interactions. This additional contribution reads

$$\left[\frac{d\sigma}{dT_{nr}} \right]_A^{\nu\mathcal{N}} = \frac{G_F^2 m_N F_A(q^2)}{4\pi} \left(1 + \frac{m_N T_{nr}}{2E_\nu^2} \right), \quad (2.23)$$

where the axial-vector form-factor is defined by

$$F_A(q^2) = \frac{8\pi}{2J+1} \left((g_A^{s,N})^2 S_{00}^{\mathcal{T}}(q^2) - g_A g_A^{s,N} S_{01}^{\mathcal{T}}(q^2) + (g_A)^2 S_{11}^{\mathcal{T}}(q^2) \right). \quad (2.24)$$

The nuclear spin is denoted by J . The structure factor is represented by $S_{ij}^{\mathcal{T}}(q^2)$ with i, j can be 0 or 1, representing isoscalar/isovector combinations, and $\mathcal{T}$ denotes the transversal component. This follows the same with DM calculations. The couplings are related by

$$g_A^p = \frac{G_F}{\sqrt{2}}(g_A - g_A^{s,N}), \quad g_A^n = -\frac{G_F}{\sqrt{2}}(g_A + g_A^{s,N}) \quad (2.25)$$

$$g_A^0 = -\frac{G_F}{\sqrt{2}}g_A^{s,N}, \quad g_A^1 = \frac{G_F}{\sqrt{2}}g_A, \quad (2.26)$$

where it is found from strong interaction studies that $g_A = 1.27641$ and $g_A^{s,N} = -0.035$ (Hoferichter *et al.*, 2020).

The structure factor needs more explanation. In the case of CE ν NS, only the transversal component is considered. It is given by

$$S_{00}^{\mathcal{T}} = \sum_L \left[\mathcal{F}_+^{\Sigma'_L}(q^2) \right]^2, \quad (2.27)$$

$$S_{11}^{\mathcal{T}} = \sum_L \left[[1 + \delta'(q^2)] \mathcal{F}_-^{\Sigma'_L}(q^2) \right]^2, \quad (2.28)$$

$$S_{01}^{\mathcal{T}} = \sum_L 2[1 + \delta'(q^2)] \mathcal{F}_+^{\Sigma'_L}(q^2) \mathcal{F}_-^{\Sigma'_L}(q^2), \quad (2.29)$$

where these can be further decomposed as proton and neutron basis

$$\mathcal{F}_\pm^{\Sigma'_L}(q^2) = \mathcal{F}_p^{\Sigma'_L}(q^2) \pm \mathcal{F}_n^{\Sigma'_L}(q^2). \quad (2.30)$$

Correction beyond the leading spin-dependent coupling is encoded within $\delta'(q^2)$ that satisfies

$$\delta'(\mathbf{q}^2) = -\frac{\mathbf{q}^2 \langle r_A^2 \rangle}{6}. \quad (2.31)$$

with $\langle r_A^2 \rangle = 0.46 \text{ fm}^2$. For completeness, we also provide the longitudinal component as follows

$$S_{00}^{\mathcal{L}} = \sum_L \left[\mathcal{F}_+^{\Sigma_L''}(q^2) \right]^2, \quad (2.32)$$

$$S_{11}^{\mathcal{L}} = \sum_L \left[[1 + \delta''(q^2)] \mathcal{F}_-^{\Sigma_L''}(q^2) \right]^2, \quad (2.33)$$

$$S_{01}^{\mathcal{L}} = \sum_L 2[1 + \delta''(q^2)] \mathcal{F}_+^{\Sigma_L''}(q^2) \mathcal{F}_-^{\Sigma_L''}(q^2). \quad (2.34)$$

These contribute to the tensor operator. Note that we have neglected two-body current contributions. We show the spin dependent structure factors with ^{72}Ge and ^{131}Xe required for

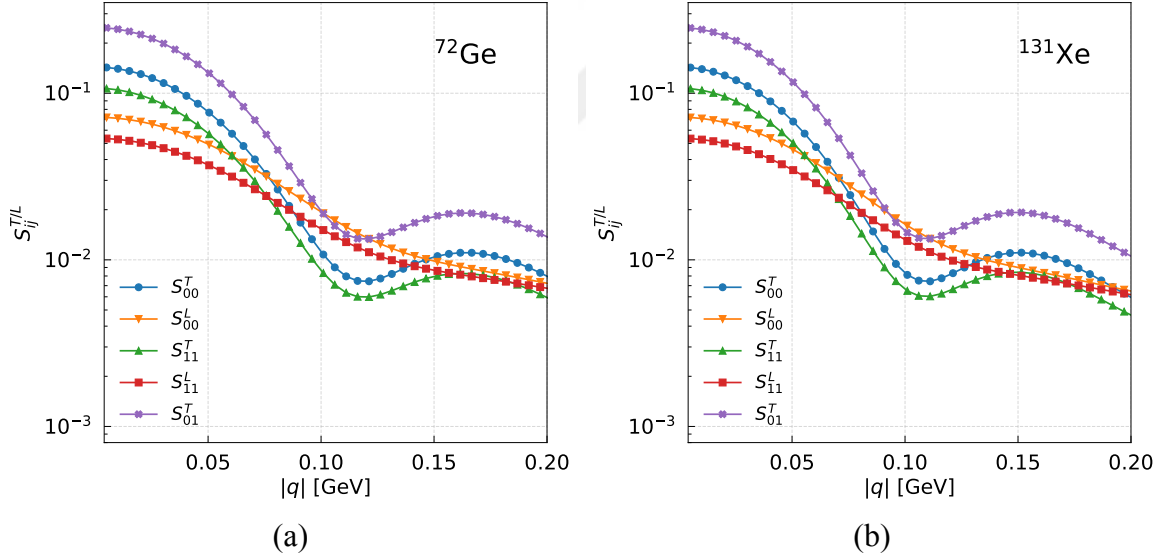


Figure 9. Spin dependent structure factors with (a) ^{72}Ge and (b) ^{131}Xe target for CE ν NS.

CE ν NS. In plotting this, we considered $b = 2.1035 \text{ fm}$ for ^{72}Ge and $b = 2.2925 \text{ fm}$ for ^{131}Xe . We followed the prescription in (Hoferichter *et al.*, 2020), where the details of nuclear response functions of proton and neutron can be elaborated for interested reader.

In this first measurement, the COHERENT collaboration measures the number of photoelectrons (PE), n_{PE} of CE ν NS events. The differential rate that depends on the nuclear

recoil energy is related to the electron recoil energy with quenching factor $Q(T_{nr})$. This quantity can then be converted to PE using (Akimov *et al.*, 2016)

$$n_{PE} = Q(T_{nr})L_Y T_{nr}, \quad (2.35)$$

with L_Y is the light yield quantity. The quenching factor connects the energy of nuclear recoils to that of electron recoils by

$$T_{er} = Q(T_{nr})T_{nr}, \quad (2.36)$$

relating the ratio between emitted scintillation light in nuclear and electron recoils.

There is a wide range of study regarding the quenching factor. The mostly used form in the literature is the Lindhard parametrization (Lindhard *et al.*, 1963)

$$Q(T_{nr}) = \frac{kg(\epsilon)}{1 + kg(\epsilon)}, \quad (2.37)$$

where

$$g(\epsilon) = 3\epsilon^{0.15} + 0.7\epsilon^{0.6} + \epsilon, \quad \epsilon = 11.5\mathcal{Z}^{7/3}T_{nr}. \quad (2.38)$$

The term $\mathcal{Z}$ denotes the number of protons in the nucleus target. Meanwhile, the coefficient of $g(\epsilon)$ has the general form

$$k = 0.133\mathcal{Z}^{2/3}\mathcal{A}^{-1/2}. \quad (2.39)$$

In practice, the parameter k is typically treated as a free parameter. This is due to the various ranges of this parameter from experimental view point. For example, result from germanium target in low-energy regime gives $k = 0.162$ (Bonhomme *et al.*, 2022). Notice that the Lindhard parametrization of the quenching factor is appropriate for high recoil energy. It is good approximation for nuclear recoil $T_{nr} > 0.254$ keV. Below this energy, there are plenty of works that still trying to model the phenomena. For example the model of high ionization-

efficiency, for germanium target the quenching factor is (Bonhomme *et al.*, 2022)

$$Q(T_{nr}) = 0.18 \left[1 - \exp \frac{0.15 - T_{nr}}{71.03} \right]. \quad (2.40)$$

This form is accepted in the nuclear recoil range $0.015 \text{ keV} < T_{nr} < 0.254 \text{ keV}$. Some

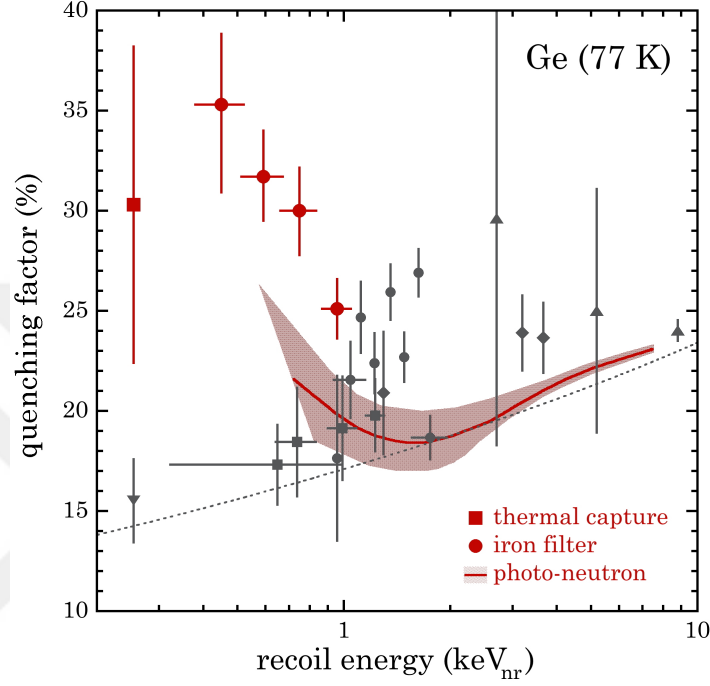


Figure 10. Quenching factor measurements for germanium from (Collar *et al.*, 2021).

quenching factor measurements can be seen in Figure 10, adapted from (Collar *et al.*, 2021). Prediction of the Lindhard quenching factor with $k = 0.157$ is shown by the dotted line. We can clearly see the deviation of the data from thermal capture and iron filter experiments for nuclear recoil energy less than 1 keV. Previous measurements are shown in gray data points.

In this study, we examine the $\text{CE}\nu\text{NS}$ process driven by solar neutrinos to present event predictions for direct-detection experiments and explore potential effects of BSM physics. Related to this purpose, we are interested in studying the light mediator models, transition magnetic moments from active into sterile neutrino, and charged Higgs contributions in the framework of $\text{E}\nu\text{ES}$ and $\text{CE}\nu\text{NS}$. Any distortion from the SM prediction could indicate the presence of new physics and may provide insights into BSM scenarios. The process is also predicted to be dominant in the core-collapse of a supernova, as well as other astrophysical and terrestrial events. The subject of direct detection of WIMP as one of the most probable

dark matter candidate has also been affected since their similar nuclear recoil energy behavior and also from the irreducible neutrino background.

2.4. Dark Matter Direct Detection Experiments

The direct detection process is one method from the three methods to detect the particle DM hypothesis. The other two are production at a particle accelerator and indirect detection from annihilation products. A schematic diagram describing these methods is shown in Figure 11. In this diagram, χ denotes the possible DM which couples to an ordinary particle P .

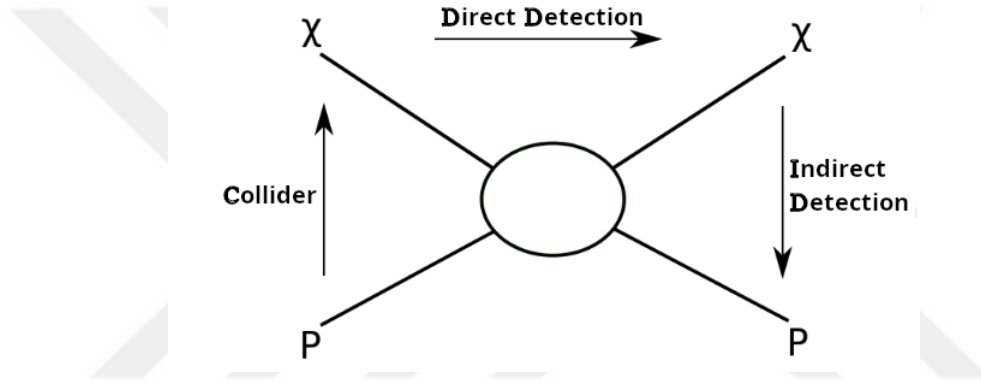


Figure 11. A general scheme for detecting a DM particle.

The goal of DD experiments is to detect WIMPs via the nuclear recoils, caused by WIMP elastic scattering, in dedicated low background detectors (Goodman & Witten, 1985).

When a DM particle elastically scatters off a target material, the process transfers energy to the nuclei. The energy can be detected in DD experiments. There are three different signals, depending on the technology utilized on the detector, and which strategies used to detect signals, focusing either on one of these or on the combination of two of these signals. They are the production of heat as phonon in a crystal, direct ionization of atoms of the target material, and a target nucleus excitation which releases photons (Undagoitia & Rauch, 2016). These technologies utilize phonon, charge, or light signals – or a combination thereof. Such methods allow for the discrimination between nuclear and electronic recoils, which is crucial for reducing experimental backgrounds. Some of the main difficulties in DM-DD lie in lowering the energy threshold sufficiently to detect very small recoils, minimizing background to enhance signal detection, and employing a detector with a large mass to increase

the likelihood of interactions within the target.

We focus this dissertation on DM-DD experiments utilizing technologies of a germanium detector and liquid noble-gas xenon detectors. Germanium detectors offer a threshold with very low energy and target material with high radio-purity. This technology has been further refined in recent years, especially to lower the energy threshold and reduce background levels. Some experiments with this technology are CoGeNT (Aalseth *et al.*, 2013), MAJORANA (Giovanetti *et al.*, 2014), and CDEX (Kang *et al.*, 2013). Meanwhile, detectors with noble-gas liquid provide the benefits of large, homogeneous targets, possible to obtain high outcome of scintillation and ionization. At present, liquid argon (LAr) and liquid xenon (LXe) are the most commonly employed detector targets, with liquid neon also being considered. To distinguish the primary background from electronic recoils versus nuclear recoils, this technology employs the method of pulse-shape discrimination and the method of the ratio of charge-to-light signal. Collaborations of XMASS (Abe *et al.*, 2013), DarkSide (Agnes *et al.*, 2015), ArDM (Badertscher *et al.*, 2013), XENON (Aprile *et al.*, 2011), PandaX (Xiao *et al.*, 2014), LUX (Akerib *et al.*, 2013), and ZEPLIN (Akimov *et al.*, 2007) are examples that use this technology.

With sensitivity down to keV energies, DM-DD experiments begin to observe neutrino interactions as target masses increase to hundreds of kilograms or even tons. Neutrinos thus constitute an important background, particularly for both electron recoil and nuclear recoil. Note that solar neutrinos can elastically scatter off a cloud of electrons in the material target through NC and CC interactions for electron-neutrino, and via NC only for other neutrino flavors. Because of the higher fluxes, solar neutrinos from pp and ${}^7\text{Be}$ are expected to be the first detectable. A recoiling electron is the resulting signal, in contrast to WIMP interactions, which have the nuclear recoil signal. Consequently, $E\nu\text{ES}$ represents a key component, particularly for experiments unable to discriminate between electron recoil and nuclear recoil signal. Solar neutrinos can also induce $\text{CE}\nu\text{NS}$, resulting in a few keV of energy from nuclear recoils. DM detectors have shown their sensitivity from recent results of XENONnT (Aprile *et al.*, 2024) and Pandax-4T (Bo *et al.*, 2024). The CDEX-10 (Geng *et al.*, 2023), furthermore, have also provided recent data that can be interpreted as neutrino-electron and coherent neutrino-nucleus even rate. Coherent scattering of solar neutrinos imposes a sensitivity limit on dark matter experiments for low-mass WIMPs, a few GeV, with cross sections

of about $\sim 10^{-45} \text{ cm}^2$. For higher-mass WIMPs, coherent scattering of atmospheric neutrinos constrains dark matter searches at approximately $\sim 10^{-49} \text{ cm}^2$. The rate of solar neutrinos should peak around January 3rd, while the WIMP rate should peak around June 2nd. It is because of the larger solid angle during the perihelion. This is also the case for atmospheric neutrinos since the changes in atmospheric density resulting from seasonal temperature variations in this period. Brief information about our considered experiments is summarized in the following.

2.4.1. CDEX-10

The CDEX experiment operates within the China Jinping Underground Laboratory (CJPL), a facility devoted to directly detect possible DM using detectors with a p -type point contact germanium (p PCGe), enabling access of ultra-low energy thresholds. The first generation of the experiment, the CDEX-1A (1B), used a 1-kg-scale single-element p PCGe cooled by a cold finger since 2010. Since its inception, CDEX has carried out searches for various exotic physics phenomena, including WIMPs, axions, and dark photons. The second generation of the program is termed CDEX-10. The total detector mass is about 10 kg and is equipped with three triple-element p PCGe, directly submerged in liquid nitrogen. It is categorized as low- Z material shielding, which provides better control of the radiation background compared to finger cooling. The scheme is depicted in Figure 12, where for the detailed analysis we refer to the collaboration works (Jiang *et. al.*, 2019).

Both $E\nu$ ES and $CE\nu$ NS processes can improve the detection sensitivity of BSMs within the region of interest (ROI) at DM-DD experiments. Accordingly, we apply the CDEX-10 data (there are 20 data points) related to neutrino-nucleus scattering. The published data are in terms of electron-equivalent recoil energy. It then needs to be converted into nuclear recoil energy by applying quenching factors. With their ultra-low energy threshold and minimal background, these data are well-suited for investigating exotic processes, such as searches for BSM physics. The CDEX-10 has a threshold of around 100 eVee (the extra "ee" represents electron equivalent) and approximately 1 count/kg/keVee/day.

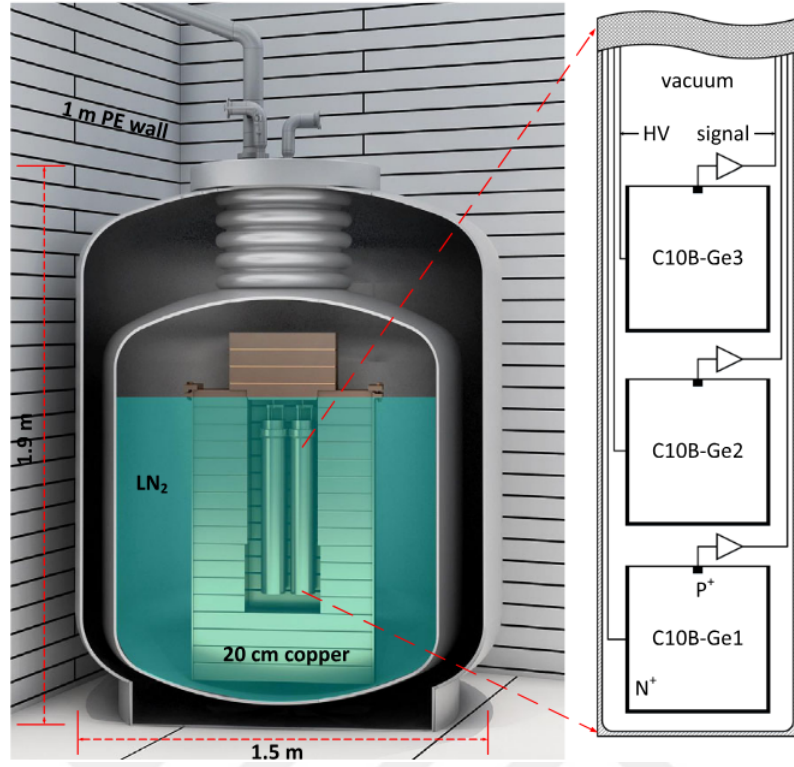


Figure 12. The detector scheme of the CDEX-10, taken from (Jiang *et. al.*, 2019).

2.4.2. PandaX-4T

The collaboration is situated at CJPL, like CDEX. It is a dual-phase xenon DM-DD facility using time projection chamber (TPC) technology. PandaX-I, the first phase of the experiment, performed observation of WIMP in 2014 with exposure of 54×80.1 kg-day produced a strong limit on the WIMP-nucleon cross-section for a mass $\lesssim 10$ GeV/c². The next phase, PandaX-II, announced in 2017, is a world-leading WIMP search with a total of 54 ton-day exposure for a WIMP mass 40 GeV/c². The experiment entered its third phase, called the PandaX-4T, with total exposure of 1.54 ton-year. We focus our work using the reported results of this latest phase of the experiment.

The PandaX-4T is part of the latest facility advancement at the CJPL. The experiment uses a liquid xenon TPC inside a dual-vessel stainless steel cryostat. The structure of the TPC forms a cylindrical shape with a 1.2 m diameter, consisting of two thick copper plates and 24 pieces of interlocking polytetrafluoroethylene panels (Zhang *et. al.*, 2019). The scheme of this TPC is shown in Figure 13. Energy deposition in xenon generates prompt scintillation

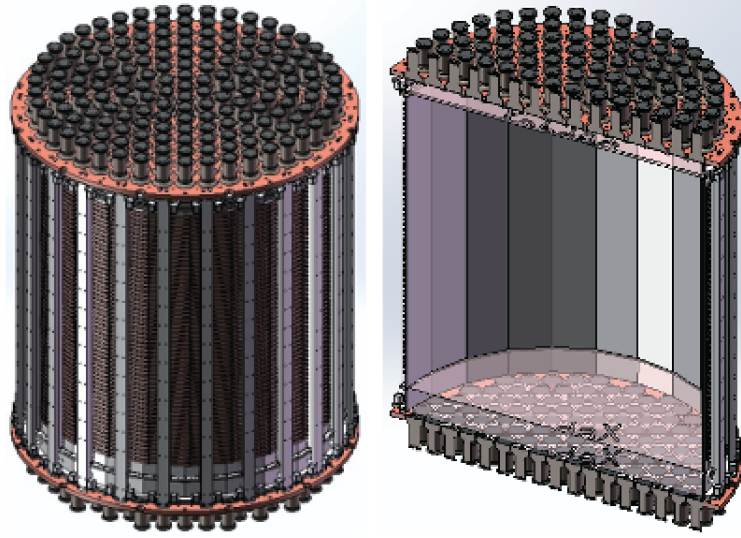


Figure 13. The detector scheme of the PandaX-4T, taken from (Zhang *et. al.*, 2019).

photons, referred to as S1, along with ionized electrons. These electrons drift upward when an electric field is applied in the sensitive volume and produce electroluminescence delayed signals, labeled S2, in the container with gaseous xenon. Photons from both S1 and S2 signals are detected by photomultiplier tubes (PMTs) arrays of Hamamatsu R11410-23 at the top and bottom layer of the TPC. An event position at three-dimension can be determined by combining the photon distribution on the PMT arrays with the time delay between S1 and S2. Since beginning operations, the experiment has pursued exploration of new physics, for example, axions, the electromagnetic properties of neutrinos, and dark photons.

The PandaX-4T collaboration uses a TPC, with 3.7 tons of LXe in the sensitive volume, for measuring the precise position, timing, and energy of each event. Electron recoil signals have been reported by the collaboration recently, with datasets termed Run0 and Run1, both showing low electron recoil energies below 30 keV. As a background in DM-DD experiments, the contribution from SM neutrino-electron scattering is well understood and approximately flat across the spectrum. Therefore, it is typically subtracted in analyzing standard DM data. Nevertheless, specific BSM scenarios have the potential to strongly increase the observed signal. Motivated by this, we explore the potential impact of BSM physics on the neutrino-electron scattering background.

2.4.3. XENONnT

The XENON program is a DM search project whose main goal is to detect WIMP dark matter using dual-phase TPCs with a liquid xenon (LXe) target. LXe is an ideal target for WIMP since it has high stopping power for gamma and beta radiation, providing self-shielding from external backgrounds. A schematic diagram of the detector can be seen in Figure 14. The first phase of the program was conducted with XENON10 (Aprile *et al.*, 2011),

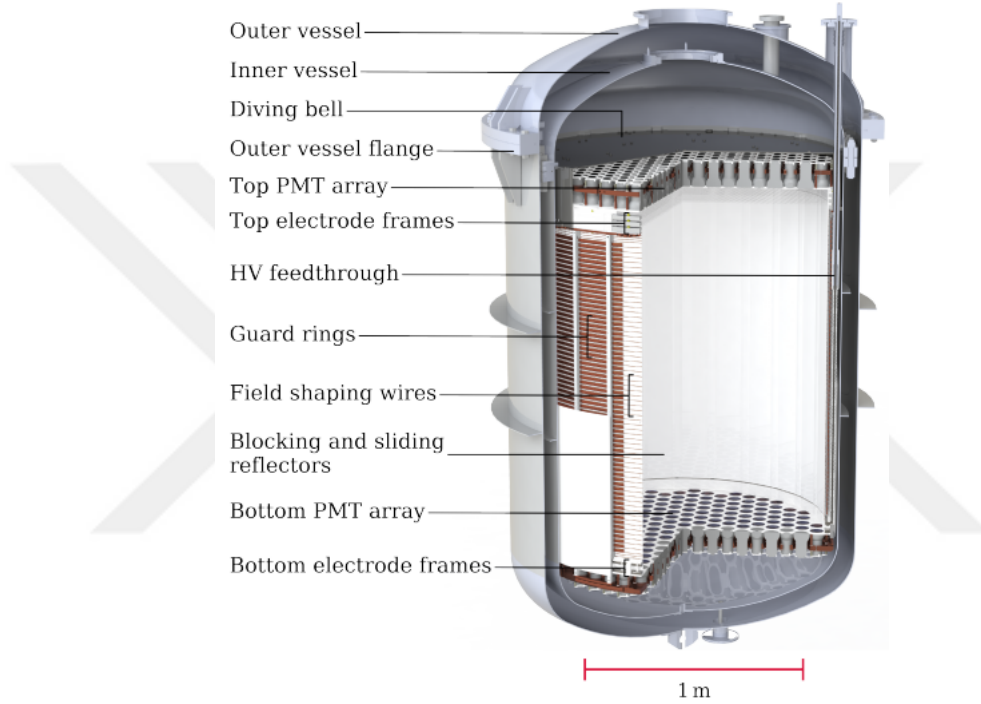


Figure 14. The detector scheme of the XENON program, taken from (Aprile *et al.*, 2020).

followed by XENON100 (Aprile *et al.*, 2012), and upgraded to XENON1T (Aprile *et al.*, 2017). Competitive results have been found by the experiments regarding spin-dependent, spin-independent, and other WIMP-nucleon interactions. Currently, the most recent detector and the largest so far is XENONnT (Aprile *et al.*, 2020), which has been constructed in order to enhance the sensitivity of the previous phase by more than an order of magnitude.

The main scientific objective of XENONnT is to detect the WIMP as a DM candidate, which is expected to generate signals of nuclear recoil. Energy depositions in liquid xenon can produce detectable prompt scintillation photons, labelled S1, and ionization electrons drift under an external electric field. These signals are then extracted into the gaseous xenon above the liquid, where secondary photon scintillations are generated, labeled S2, through

electroluminescence at the same wavelength. Between the S1 and S2 signals, the time interval is proportional to the interaction depth.

The XENONnT is equipped with 1.16 ton-year exposure with a 5.9 ton of LXe fiducial mass. The experiment has reached an energy threshold of around 1 keV_{ee}. The experiment recently published a blinded analysis of low-energy electronic recoils, where it observed no excess above background. This achievement makes the facility an ideal framework that involves spectral deviation at low energies, such as to probe BSM physics phenomena that we are focused on in this work.



3. BEYOND THE STANDARD MODEL PHYSICS WITH NEUTRINO INTERACTIONS

While the SM provides a highly successful framework for describing electroweak and strong interactions, certain limitations suggest the need to extend the theory. Below the electroweak symmetry-breaking scale, additional types of neutrino interactions with fundamental particles may become relevant for new physics phenomena. In various frameworks of BSM, particles from hidden sectors with low mass are commonly introduced to account for such assumptions. In this work, we focus on new light mediator models, the sterile neutrino dipole portal, and the Higgs triplet model. These scenarios may potentially account for the existence of nonzero neutrino masses.

3.1. New Light Mediator Models

In experiments involving neutrinos with low-energy scales, new physics effects are potentially probed through events of $E\nu$ ES or $CE\nu$ NS in the presence of new light mediators. These new mediators interact with the lepton component in the SM. Such models are designed to introduce a minimal set of new particles and interactions, effectively serving as simplified limits of broader BSM scenarios. These models can be characterized by just several free parameters, like the coupling constants and masses associated with the new particles. These models enable the study of a wide range of new physics signatures without the need to specify a complete high-energy theory. In this work, we concentrate on extensions of the SM with universal light mediator models and anomaly-free $U(1)'$ gauge symmetry. The universal models include scalar, vector, and tensor interactions that are consistent with Lorentz invariance. Meanwhile, the anomaly-free $U(1)'$ considers only a new vector type of mediator. Particularly for this case, we consider the $B - L$, $B - 3L_e$, $B - 3L_\mu$, and $B - 3L_\tau$ models as well as the leptophilic $L_e - L_\mu$, $L_e - L_\tau$, and $L_\mu - L_\tau$ models.

The Feynman diagrams for both $E\nu$ ES and $CE\nu$ NS in the presence of light mediator models under consideration are shown in Figure 15. We calculate contributions from the scalar, vector, tensor, and the $U(1)'$ mediators to the $E\nu$ ES and $CE\nu$ NS. The mediator mass

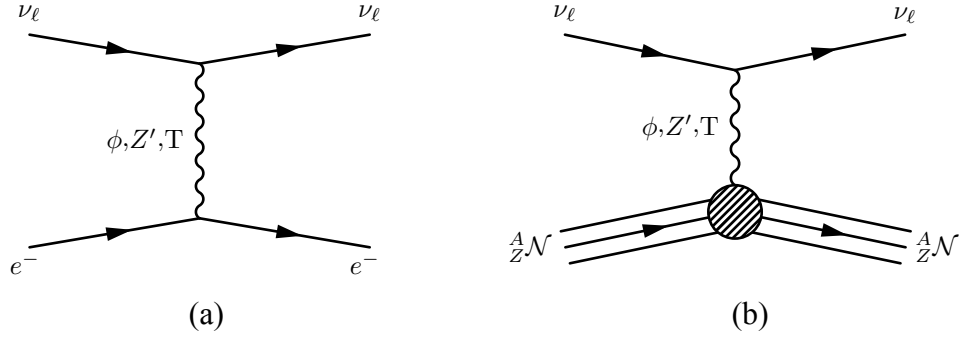


Figure 15. Feynman diagrams of (a) $E\nu\text{ES}$ and (b) $CE\nu\text{NS}$ in the presence of light mediators.

is denoted by m_X . For the coupling constant, we define $g_X = \sqrt{g_X^f g_X^\nu}$, with g_X^f denotes either the electron $f = e$ or quark $f = q$ coupling interaction, and g_X^ν the neutrino coupling interaction with the X mediator model. We explore how well the light mediators can be constrained utilizing data from recent DM-DD facilities. Especially, the low-energy recoil data from the considered experiments potentially impose competitive or even stronger limits on the $m_X - g_X$ parameter space. It is because, as the mediator mass from new interactions becomes comparable to the momentum transfer during interaction, its presence may suggest modification of the recoil spectrum, resulting in tighter bounds.

3.1.1. Universal Scalar Mediator

We examine a universal mediator with scalar type, ϕ . This model can mediate interactions between neutrinos and quarks as well as electrons. The corresponding general Lagrangian to explain this interaction is (Cerd no *et al.*, 2016)

$$\mathcal{L}_\phi \supset -\phi \left[\sum_q g_\phi^q \bar{q}q + g_\phi^e \bar{e}e + g_\phi^{\nu_\ell} \bar{\nu}_{\ell R} \nu_{\ell L} + \text{h.c.} \right], \quad (3.1)$$

where g_ϕ^q is the quark, g_ϕ^e the electron, and $g_\phi^{\nu_\ell}$ the neutrino coupling constant of the universal scalar model, where the label ℓ represents e, μ , or τ . For the $E\nu\text{ES}$ process, we can write the amplitude from this Lagrangian as

$$-i\mathcal{M}_\phi^{\nu e} = [-ig_\phi^{\nu_\ell} \bar{\nu}(p_3) P_L \nu(p_1)] \left[\frac{-i}{q^2 - m_\phi^2} \right] [-ig_\phi^e \bar{e}(p_4) e(p_2)]. \quad (3.2)$$

Meanwhile, for the CE ν NS process, we can write it as

$$-i\mathcal{M}_\phi^{\nu\mathcal{N}} = [-ig_\phi^{\nu\ell}\bar{\nu}(p_3)P_L\nu(p_1)] \left[\frac{-i}{q^2 - m_\phi^2} \right] [-ig_\phi^e\bar{N}(p_4)N(p_2)F(q^2)]. \quad (3.3)$$

We can see the difference appears only in the electron and the nucleus current part. Particularly, it is due to the complex structure of the nucleus that is incorporated within $F(q^2)$, the nucleus form factor. Hence, in the calculation, we can interchange between the result from E ν ES to CE ν NS by the inclusion of the form factor, as well as other quantities that characterize the process, such as mass and recoil energy.

The average of the absolute amplitude square can then be calculated, leading to the differential cross-section contribution. Doing so, for the CE ν NS process, we have

$$\begin{aligned} \langle |\mathcal{M}_\phi^{\nu\mathcal{N}}|^2 \rangle &= \frac{Q_\phi^2 |F(q^2)|^2}{8(q^2 - m_\phi^2)^2} \text{Tr}[\not{p}_3(1 - \gamma_5)\not{p}_1(1 + \gamma_5)] \text{Tr}[(\not{p}_4 + m_\mathcal{N})(\not{p}_2 + m_\mathcal{N})] \\ &= \frac{Q_\phi^2 |F(q^2)|^2}{8(2m_\mathcal{N}T_{nr} + m_\phi^2)^2} (4p_1 \cdot p_3)(4p_2 \cdot p_4 + 4m_\mathcal{N}^2). \end{aligned} \quad (3.4)$$

In the E ν ES process, we need to change the mass $m_\mathcal{N} \rightarrow m_e$, recoil energy $T_{nr} \rightarrow T_e$, and the multiplication factor of the coupling with form factor $Q_\phi^2 |F(q^2)|^2 \rightarrow (g_\phi^e g_\phi^{\nu\ell})^2$. The latter consists of the multiplication of either a neutrino with electron coupling or a neutrino with quark components of the nucleus. Implementing the kinematics of Equation (A.8), we can calculate the differential cross section. We obtain contribution to the E ν ES from the universal light scalar mediator as

$$\left[\frac{d\sigma_{\nu\ell}}{dT_{er}} \right]_\phi^{\nu e} = \frac{(g_\phi^e g_\phi^{\nu\ell})^2}{4\pi(m_\phi^2 + 2m_e T_{er})^2} \left[\frac{T_{er} m_e^2}{E_\nu^2} \right]. \quad (3.5)$$

Meanwhile, the contribution to the CE ν NS process can be written as

$$\left[\frac{d\sigma_{\nu\ell}}{dT_{nr}} \right]_\phi^{\nu\mathcal{N}} = \frac{Q_\phi^2}{4\pi(m_\phi^2 + 2m_\mathcal{N} T_{nr})^2} \left[\frac{T_{nr} m_\mathcal{N}^2 |F(|\vec{q}|^2)|^2}{E_\nu^2} \right], \quad (3.6)$$

where Q_ϕ represents the nucleus scalar charge

$$Q_\phi = \left[Z \sum_q g_\phi^q \frac{m_p}{m_q} f_{T_q}^p + N \sum_q g_\phi^q \frac{m_n}{m_q} f_{T_q}^n \right], \quad (3.7)$$

with $f_{T_u}^p = 0.0208$, $f_{T_u}^n = 0.0189$, $f_{T_d}^p = 0.0411$, and $f_{T_d}^n = 0.0451$ are parameters of the hadronic structure (Hoferichter *et al.*, 2015). In the above relation, m_p is the proton mass, m_n is the neutron mass, and m_q is the quark mass. In our calculation, identical couplings for the up quark and down quark are assumed. Since the cross section of this model has no interference with the SM, the effect of the scalar interaction is simply added to the standard E ν ES and CE ν NS.

3.1.2. Universal Vector Mediator

Next, we evaluate an extra vector mediator, which interacts with quarks, electrons, and neutrinos in the SM. The mediator is generally expressed as Z' . The general Lagrangian associated with this mediator can be expressed as (Cerd no *et al.*, 2016)

$$\mathcal{L}_{Z'} \supset Z'_\mu \left[\sum_q Q_{Z'}^q g_{Z'}^q \bar{q} \gamma^\mu q + Q_{Z'}^e g_{Z'}^e \bar{e} \gamma^\mu e + Q_{Z'}^{\nu_\ell} g_{Z'}^{\nu_\ell} \bar{\nu}_{\ell L} \gamma^\mu \nu_{\ell L} \right]. \quad (3.8)$$

In this relation, $g_{Z'}^q$ is the coupling constant for quark, $g_{Z'}^e$ for electron, and $g_{Z'}^{\nu_\ell}$ for neutrino to the new gauge-boson. The quarks, electrons, and neutrinos vector charges are given by $Q_{Z'}^q$, $Q_{Z'}^e$, and $Q_{Z'}^{\nu_\ell}$, respectively.

Calculation of the differential cross section follows similar steps to the scalar case. In the case of E ν ES, we can write the amplitude as

$$-i\mathcal{M}_{Z'}^{\nu e} = [-iQ_{Z'}^{\nu_\ell} g_{Z'}^{\nu_\ell} \bar{\nu}(p_3) \gamma^\mu P_L \nu(p_1)] \left[\frac{-ig_{\mu\nu}}{q^2 - m_{Z'}^2} \right] [-iQ_{Z'}^e g_{Z'}^e \bar{e}(p_4) \gamma^\nu e(p_2)]. \quad (3.9)$$

Meanwhile, for the CE ν NS process, we can write the amplitude as

$$-i\mathcal{M}_{Z'}^{\nu N} = [-iQ_{Z'}^{\nu_\ell} g_{Z'}^{\nu_\ell} \bar{\nu}(p_3) \gamma^\mu P_L \nu(p_1)] \left[\frac{-ig_{\mu\nu}}{q^2 - m_{Z'}^2} \right] [-iQ_{Z'} \bar{N}(p_4) \gamma^\nu N(p_2) F(q^2)]. \quad (3.10)$$

The complex structure of the nucleus can be accommodated by introducing a form factor $F(q^2)$ and redefining the overall charge factor. Considering vector current conservation, the contributions occur from valence quarks only, where the corresponding charges are simply

added. It is encapsulated by the weak vector charge

$$\begin{aligned} Q_{Z'} &= \left[Z \sum_q Q_{Z'}^q g_{Z'}^q + N \sum_q Q_{Z'}^q g_{Z'}^q \right] Q_{Z'}^\ell g_{Z'}^{\nu_\ell} \\ &= \left[Z(2Q_{Z'}^u + Q_{Z'}^d) + N(Q_{Z'}^u + 2Q_{Z'}^d) \right] g_{Z'}^q g_{Z'}^{\nu_\ell} Q_{Z'}^\ell. \end{aligned} \quad (3.11)$$

This is the general form in terms of charges $Q_{Z'}^u, Q_{Z'}^d$, and $Q_{Z'}^\ell$. For simplicity, we show here for $\text{CE}\nu\text{NS}$, which can be converted to the $\text{E}\nu\text{ES}$ process. We obtain the average of the absolute square of the amplitude as

$$\begin{aligned} \langle |\mathcal{M}_{Z'}^{\nu\mathcal{N}}|^2 \rangle &= \frac{Q_{Z'}^2 |F(q^2)|^2}{8(q^2 - m_{Z'}^2)} \text{Tr}[\not{p}_3 \gamma^\mu (1 - \gamma_5) \not{p}_1 \gamma^\nu (1 + \gamma_5)] \text{Tr}[(\not{p}_4 + m_{\mathcal{N}}) \gamma_\mu (\not{p}_2 + m_{\mathcal{N}}) \gamma_\nu] \\ &= \frac{Q_{Z'}^2 |F(q^2)|^2}{8(2m_{\mathcal{N}} T_{nr} + m_{Z'}^2)} 32 [2(p_1 \cdot p_2)(p_3 \cdot p_4) + 2(p_1 \cdot p_4)(p_2 \cdot p_3) - 2(p_1 \cdot p_3)m_{\mathcal{N}}^2]. \end{aligned} \quad (3.12)$$

Implementing the kinematics of Equation (A.8), we can evaluate the differential cross section of the universal light vector mediator. Doing so, we can write the contribution to the $\text{CE}\nu\text{NS}$ cross section as

$$\left[\frac{d\sigma}{dT_{nr}} \right]_{Z'}^{\nu\mathcal{N}} = \frac{Q_{Z'}^2 m_{\mathcal{N}} |F(|\vec{q}|^2)|^2}{2\pi(m_{Z'}^2 + 2m_{\mathcal{N}} T_{nr})^2} \left(1 - \frac{m_{\mathcal{N}} T_{nr}}{2E_\nu^2} \right), \quad (3.13)$$

where $Q_{Z'}$ denotes the weak vector charge of the nucleus. For the case of universal interaction, this type of mediator couples to all SM fermions. In this case $Q_{Z'}^u = Q_{Z'}^d = Q_{Z'}^\ell = 1$. Notice that there is an interference of this model with the SM. We can then express the cross-section as

$$\left[\frac{d\sigma}{dT_{nr}} \right]_{\text{SM}+Z'}^{\nu\mathcal{N}} = \left[1 + \frac{Q_{Z'}}{\sqrt{2}G_F Q_{SM}(m_{Z'}^2 + 2m_{\mathcal{N}} T_{nr})} \right]^2 \left[\frac{d\sigma}{dT_{nr}} \right]_{\text{SM}}. \quad (3.14)$$

The contribution from the Z' mediator to the $\text{E}\nu\text{ES}$ process is determined by applying

$$g_V^e \longrightarrow g_V^e + \frac{Q_{Z'}^e g_{Z'}^e Q_{Z'}^{\nu_\ell} g_{Z'}^{\nu_\ell}}{2\sqrt{2}G_F(m_{Z'}^2 + 2m_e T_{er})} \quad (3.15)$$

into Equation (2.12). The relation applies to low- and high-mass regimes, since the typical

momentum transfer is approximately 100 keV in DM-DD experiments. For the universal model, $Q_{Z'}^e = Q_{Z'}^{\nu_\ell} = 1$. This set up leads to a similar observable with the $B - L$ model where $Q_{Z'}^e = Q_{Z'}^{\nu_\ell} = -1$. Discrimination between these two scenarios becomes apparent in the CE ν NS framework, in which vector charges of the quark have a charge factor of 1/3. Meanwhile, both models yield identical predictions for E ν ES.

3.1.3. Universal Tensor Mediator

We next evaluate the contribution of the universal light tensor mediator. For a mediator with a tensor interaction T, the relevant Lagrangian can be written as (Barranco *et al.*, 2012)

$$\mathcal{L}_T \supset \left[\sum_q g_T^q \bar{q} \sigma^{\mu\nu} q + g_T^e \bar{e} \sigma^{\mu\nu} e - g_T^{\nu_\ell} \bar{\nu}_{\ell R} \sigma^{\mu\nu} \nu_{\ell L} \right] T_{\mu\nu} \quad (3.16)$$

with $\sigma_{\mu\nu} = i(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu)/2$. Here, g_T^q is the coupling constant for quark, g_T^e is for electron, and $g_T^{\nu_\ell}$ is for neutrino to the new gauge-boson. We can write the amplitude of the tensor contribution to the E ν ES process as

$$-i\mathcal{M}_T^{\nu e} = [-ig_T^{\nu_\ell} \bar{\nu}(p_3) \sigma^{\mu\nu} P_L \nu(p_1)] \left[\frac{-ig_{\mu\mu'} g_{\nu\nu'}}{q^2 - m_T^2} \right] [-ig_T^e \bar{e}(p_4) \sigma^{\mu'\nu'} e(p_2)] \quad (3.17)$$

and the contribution to the CE ν NS process as

$$-i\mathcal{M}_T^{\nu N} = [-ig_T^{\nu_\ell} \bar{\nu}(p_3) \sigma^{\mu\nu} P_L \nu(p_1)] \left[\frac{-ig_{\mu\mu'} g_{\nu\nu'}}{q^2 - m_T^2} \right] [-iQ_T \bar{N}(p_4) \sigma^{\mu'\nu'} N(p_2)]. \quad (3.18)$$

The tensor charge in this case is defined by

$$Q_T = \left(Z \sum_q g_T^q \delta_q^p + N \sum_q g_T^q \delta_q^n \right) g_T^{\nu_\ell}. \quad (3.19)$$

In this work, we use $\delta_u^p = \delta_d^n = 0.54$ and $\delta_u^n = \delta_d^p = -0.23$, from the previous calculation of effective field theory (Cirelli *et al.*, 2013). Working on the E ν ES process, we obtain the

average of the absolute square of the amplitude as

$$\begin{aligned} \langle |\mathcal{M}_T^{\nu e}|^2 \rangle &= \frac{(g_T^{\nu \ell} g_T^e)^2 m_e}{128(q^2 - m_T^2)^2} \text{Tr}[\not{p}_3(\gamma^\mu \gamma^\nu - \gamma^\nu \gamma^\mu)(1 - \gamma^5)\not{p}_1(1 + \gamma^5)(\gamma^{\mu'} \gamma^{\nu'} - \gamma^{\nu'} \gamma^{\mu'})] \\ &\quad \times \text{Tr}[\not{p}_4(\gamma_\mu \gamma_\nu - \gamma_\nu \gamma_\mu)\not{p}_2(\gamma_{\mu'} \gamma_{\nu'} - \gamma_{\nu'} \gamma_{\mu'})], \end{aligned} \quad (3.20)$$

where the contraction of the two traces gives

$$128 [(p_1 \cdot p_3)(p_2 \cdot p_4) - 2((p_1 \cdot p_2)(p_3 \cdot p_4) + (p_1 \cdot p_4)(p_2 \cdot p_3))]. \quad (3.21)$$

Therefore, we obtain the contribution of the universal light tensor mediator to $E\nu$ ES to be

$$\left[\frac{d\sigma}{dT_{er}} \right]_T^{\nu e} = \frac{m_e (g_T^{\nu \ell} g_T^e)^2 Z_{\text{eff}}(T_{er})}{2\pi(m_T^2 + 2m_e T_{er})^2} \left[1 + 2 \left(1 - \frac{T_{er}}{E_\nu} \right) + \left(1 - \frac{T_{er}}{E_\nu} \right)^2 - \frac{m_e T_{er}}{E_\nu^2} \right]. \quad (3.22)$$

We can transform the formula into the $CE\nu$ NS case by replacing $m_e \rightarrow m_{\mathcal{N}}$, $T_{er} \rightarrow T_{nr}$, $Z_{\text{eff}}(T_{er}) \rightarrow F(q^2)$, and the coupling factor $(g_T^{\nu \ell} g_T^e)$ to a form of tensor charge that include quark component Q_T . Then, we obtain the tensor mediator contribution to the $CE\nu$ NS

$$\left[\frac{d\sigma}{dT_{nr}} \right]_T^{\nu \mathcal{N}} = \frac{2Q_T^2 m_{\mathcal{N}} |F(q^2)|^2}{\pi(m_T^2 + 2m_{\mathcal{N}} T_{nr})^2} \left[1 - \frac{m_{\mathcal{N}} T_{nr}}{4E_\nu^2} \right]. \quad (3.23)$$

In the SM, the nuclear effects in the $CE\nu$ NS process are captured by a form factor with spin-independent. It indicates the distribution of the nucleus weak charge. Since the point-nucleus approximation successfully explains the SM with vector interaction, the same formalism is generally adopted in most studies of tensor interactions. In the approximation, an identical form factor is typically used. However, research on DM using the framework of effective field theory suggests modeling the nucleus as a collection of interacting nucleons. The approach captures additional internal effects between a pointlike particle and the nucleus in the process of mapping relativistic interactions to non-relativistic operators. Consequently, it is reasonable to go beyond the point-nucleus approximation and use functions of nuclear response to study nuclear effects in tensor interactions. We focus in this work on universal tensor interactions, but we also briefly discuss developments in the framework of nuclear response. In this approach, the differential cross section of the $CE\nu$ NS process is decomposed

into the transversal and longitudinal terms of the nuclear response. The relevant contribution to the CE ν NS process is (Hoferichter *et al.*, 2020)

$$\left[\frac{d\sigma}{dT_{nr}} \right]_{\text{T}}^{\nu\mathcal{N}} = \frac{g_{\text{T}}^4}{(m_{\text{T}}^2 + 2m_{\mathcal{N}}T_{nr})^2} \frac{m_{\mathcal{N}}}{2J+1} \left[\left(2 - \frac{m_{\mathcal{N}}T_{nr}}{E_{\nu}^2} - \frac{2T_{nr}}{E_{\nu}} \right) S^{\mathcal{T}}(q^2) + \left(1 - \frac{T_{nr}}{E_{\nu}} \right) S^{\mathcal{L}}(q^2) \right]. \quad (3.24)$$

with the nucleus total angular momentum in the ground state represented by J . Notice that here both the transversal $S^{\mathcal{T}}$ and longitudinal $S^{\mathcal{L}}$ nuclear response are involved, as discussed in the axial-vector contribution in the SM. The general form of these parameters, setting $\kappa = \mathcal{T}, \mathcal{L}$, can be written as

$$S^{\kappa}(q^2) = g_f^2 [(\tilde{g}_{\text{T}}^0)^2 S_{00}^{\kappa} + (\tilde{g}_{\text{T}}^1)^2 S_{11}^{\kappa} + \tilde{g}_{\text{T}}^0 \tilde{g}_{\text{T}}^1 S_{01}^{\kappa}], \quad (3.25)$$

where

$$\tilde{g}_{\text{T}}^0 = \frac{\delta_u^{(p)} + \delta_d^{(p)} + \delta_u^{(n)} + \delta_d^{(n)}}{2} \quad \text{and} \quad \tilde{g}_{\text{T}}^1 = \frac{\delta_u^{(p)} + \delta_d^{(p)} - \delta_u^{(n)} - \delta_d^{(n)}}{2}. \quad (3.26)$$

Comparing with the Lagrangian of the tensor interaction, we have $g_f = G_F \epsilon_{\ell}^{\text{T}} \sqrt{2}$, where in the case of new light mediator model we have

$$g_f^2 = \frac{g_{\text{T}}^4}{(m_{\text{T}}^2 + 2m_{\mathcal{N}}T_{nr})^2}. \quad (3.27)$$

The nuclear response for the transversal case follows the same with the axial-vector. Regarding the longitudinal nuclear response, it satisfies

$$S_{00}^{\mathcal{L}} = \sum_L \left[\mathcal{F}_+^{\Sigma_L''}(q^2) \right]^2, \quad (3.28)$$

$$S_{11}^{\mathcal{L}} = \sum_L \left[[1 + \delta''(q^2)] \mathcal{F}_-^{\Sigma_L''}(q^2) \right]^2, \quad (3.29)$$

$$S_{01}^{\mathcal{L}} = \sum_L 2[1 + \delta''(q^2)] \mathcal{F}_+^{\Sigma_L''}(q^2) \mathcal{F}_-^{\Sigma_L''}(q^2), \quad (3.30)$$

with

$$\mathcal{F}_{\pm}^{\Sigma_L''}(q^2) = \mathcal{F}_p^{\Sigma_L''}(q^2) \pm \mathcal{F}_n^{\Sigma_L''}(q^2). \quad (3.31)$$

Here, the correction beyond the leading spin-dependent coupling given by

$$\delta''(\mathbf{q}^2) = -\frac{g_{\phi NN} F_{\pi}}{g_A m_{\mathcal{N}}} \frac{\mathbf{q}^2}{\mathbf{q}^2 + m_{\pi}^2} \quad (3.32)$$

with $\frac{g_{\pi NN} F_{\pi}}{g_A m_{\mathcal{N}}} = 1 + 1\%$ (Hoferichter *et al.*, 2020) and $m_{\phi} = 135.0$ MeV. In the relations, we have neglected two-body current contributions. As for the axial-vector contribution to the SM, the calculation of nuclear response functions follows the description in (Hoferichter *et al.*, 2020). Notice that interference terms with the SM are absent from this tensor interaction.

3.1.4. Anomaly-free $U(1)'$ Models

Various studies have explored SM extensions that incorporate a $U(1)'$ gauge-symmetry along with its corresponding neutral boson, Z' (He *et al.*, 1991; Langacker, 2009). A key necessity is that the condition of anomaly-free theory needs to be satisfied. Such extensions can be realized by extending the SM to include three right-handed neutrinos (Allanach *et al.*, 2019). Furthermore, these models can address several unresolved issues within the SM, such as grand unified theories, the mechanism of baryogenesis via leptogenesis, the properties of dark matter, and experimental anomalies.

Recall that the NC Lagrangian for fermion interactions in the SM are explained by

$$-\mathcal{L}_{\text{NC}}^{\text{SM}} = g \sin \theta_W J_{\text{EM}}^{\mu} A_{\mu} + \frac{g^2}{\cos^2 \theta_W} J^{\mu} Z_{\mu}, \quad (3.33)$$

where A_{μ} denotes the photon field and Z_{μ} denotes the massive neutral boson. The weak angle is $\theta_W = \tan^{-1}(g'/g)$, with g represents the gauge coupling of $SU(2)_L$ and g' represents the coupling of $U(1)_Y$. The currents in the Lagrangian are

$$J_{\text{EM}}^{\mu} = \sum_i \bar{f}_i \gamma^{\mu} f_i, \quad J^{\mu} = \sum_i \bar{f}_i \gamma^{\mu} [g_L^i P_L + g_R^i P_R] f_i, \quad (3.34)$$

with q_i is the electric charge of the i -th fermion f_i . The chiral couplings are $g_L^i = T_{3L}^i - q_i \sin^2 \theta_W$ and $g_R^i = T_{3R}^i - q_i \sin^2 \theta_W$, where T_{3L}^i and T_{3R}^i are the weak isospin third component for the left and right chiral component of fermion f_i . The electric charge of f_i is denoted by q_i , where $q_i = T_{3L,R}^i + y_W^i/2$ with y_W^i denotes the weak hypercharges. As we extent the SM to $SU(2)_L \times U(1)_Y \times U(1)'$, we obtain an additional term

$$-\mathcal{L}'_{\text{NC}} = g' J'^{\mu} Z'_{\mu}, \quad (3.35)$$

where g' and Z'_{μ} are the additional gauge coupling and boson of the $U(1)'$, while its current is

$$J'^{\mu} = \frac{1}{2} \sum_i \bar{f}_i \gamma^{\mu} [g'_V - g'_A \gamma_5] f_i, \quad (3.36)$$

where the $g'_{V,A} = g'_L \pm g'_R$ are their vector and axial couplings. More conveniently, charge of the $U(1)'$ is specified by the fermion f and its antifermion f^c (c means conjugate) using $Q'_f = g'_L$ and $Q'_{f^c} = g'_R$. The gauge bosons interaction terms with scalar multiplet ϕ emerge from the kinetic term of the Lagrangian $\mathcal{L}_{\phi}^{\text{kin}} \sim (D^{\mu} \phi)^{\dagger} D_{\mu} \phi$, where the NC term is

$$D_{\mu} \phi_i = (\partial_{\mu} + ieq_i A_{\mu} + igQ_i Z_{\mu} + ig'Q'_i Z'_{\mu}) \phi_i. \quad (3.37)$$

As an example for the Higgs doublet $\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}$, the neutral component has $T_3^{\phi^0} = -1/2$ and then $y_W^{\phi^0} = 1/2$. When the ϕ^0 neutral component acquires vacuum expectation value (vev), the photon field A_{μ} stays massless while the Z and Z' bosons develop a mass term. The Lagrangian would be $\mathcal{L}_Z^{\text{mass}} = \frac{1}{2} M_{ab}^2 Z_{a\mu} Z_b^{\mu}$ where the mass term general form is

$$M_{ab}^2 = 2g_a g_b \sum_i Q_i^a Q_i^b |\langle \phi_i \rangle|^2. \quad (3.38)$$

The index a, b runs for Z and Z' . In the SM case where there is no extra Z' , then

$$M_Z^2 = \frac{1}{2} g^2 \sum_i |\langle \phi_i \rangle|^2 = \frac{1}{4} g^2 v^2 = \frac{M^2}{\cos^2 \theta_W}, \quad (3.39)$$

with $\sum_i |\langle \phi_i \rangle|^2 = v^2/2 \sim (\sqrt{2}G_F)^{-1} \sim (246 \text{ GeV})^2$. For the Higgs in the SM, we have

$$M_Z^2 = \frac{g^2}{4\sqrt{2}G_F\rho_0} = \frac{M_W^2}{\rho_0 \cos^2 \theta_W}, \quad (3.40)$$

where $\rho_0 \approx 1$. Incorporating the Z' leads to the following form of the mass matrix

$$M_{ZZ'}^2 \equiv \begin{pmatrix} M_Z^2 & \Delta^2 \\ \Delta^2 & M_{Z'}^2 \end{pmatrix} = \begin{pmatrix} 2g^2 \sum_i T_3^{i2} |\langle \phi_i \rangle|^2 & 2gg' \sum_i T_3^i Q^i |\langle \phi_i \rangle|^2 \\ 2gg' \sum_i T_3^i Q^i |\langle \phi_i \rangle|^2 & 2g'^2 \sum_i Q^{i2} |\langle \phi_i \rangle|^2 \end{pmatrix}. \quad (3.41)$$

The general form of eigenvalues for the $M_{ZZ'}^2$ are

$$M_{1,2}^2 = \frac{1}{2} \left[M_Z^2 + M_{Z'}^2 \mp \sqrt{(M_Z^2 - M_{Z'}^2)^2 + 4\Delta^4} \right]. \quad (3.42)$$

The Z and Z' can be connected with the rotation operator $U = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix}$, where the angle θ satisfies

$$\tan(2\theta) = \frac{2\Delta^2}{M_Z^2 - M_{Z'}^2} \quad \text{and} \quad \tan^2 \theta = \frac{M_Z^2 - M_1^2}{M_2^2 - M_Z^2}. \quad (3.43)$$

A consistent low-energy theory has to be free of anomalies, meaning it should satisfy the condition of anomaly cancellation. This condition is generally caused by a one-loop triangle diagram, which includes the mixed gauge-gravitation. An anomaly can be represented by

$$\mathcal{A}_{abc} = \text{Tr}[\{T_a, T_b\}T_c], \quad (3.44)$$

where $T_{a,b,c}$ denotes the group generator for the corresponding symmetry. Theory is anomalous if this relation does not vanish as we sum over all chiral fermions in the theory with a relative minus sign for the left- and right-handed components (Allanach *et al.*, 2019). In fact, the SM which constructed by $SU(3)_C \times SU(2)_L \times U(1)_Y$ is anomaly-free. We can check for $[SU(2)_L]^2[U(1)_Y]$ this condition. From the anomaly relation, $\text{Tr}[\{\sigma_a, \sigma_b\}Y_c] = \delta_{ab} \sum_{\text{left}} Y_c = 6Y_Q + 2Y_L = 0$ needs to be satisfied. In its development, gauge boson from gravitation is also incorporated in the anomaly-free condition. With this consideration, we

find the constraints for anomaly-free condition of the SM as

$$[SU(2)_L]^2[U(1)_Y] \rightarrow 3Y_Q + Y_L = 0, \quad (3.45)$$

$$[SU(3)_C]^2[U(1)_Y] \rightarrow 2Y_Q - Y_u - Y_d = 0, \quad (3.46)$$

$$[U(1)_Y]^3 \rightarrow 3(2Y_Q^3 - Y_u^3 - Y_d^3) + (2Y_L^3 - Y_e^3 - Y_\nu^3), \quad (3.47)$$

$$[\text{gravitation}]^2[U(1)_Y] \rightarrow 3(2Y_Q - Y_u - Y_d) + (2Y_L - Y_e - Y_\nu). \quad (3.48)$$

When the hypercharge values in the Table 1 are substituted, we can see that these are anomaly-free. The general solution to these nonlinear equations are

$$Y_L = -\frac{a}{2} - b, \quad Y_e = -a - b, \quad Y_\nu = -b \quad (3.49)$$

$$Y_Q = \frac{a}{6} + \frac{b}{3}, \quad Y_u = \frac{2a}{3} + \frac{b}{3}, \quad Y_d = -\frac{a}{3} + \frac{b}{3}, \quad (3.50)$$

for any a and b . We obtain the SM hypercharge values for $a = 1$ and $b = 0$.

The SM extension by incorporating $U(1)'$ with hypercharge Y' generates extra anomaly-free conditions. This new gauge symmetry satisfies anomaly-free conditions of the SM with $Y \rightarrow Y'$. Additionally, the anomaly cancellation also needs to be satisfied by $[U(1)_Y]^2[U(1)']$ as well as $[U(1)_Y][U(1)']^2$ combinations. Using the general solutions above for $a = 0$ and $b = 1$, the hypercharge assignments for this additional $U(1)'$ become

$$Y'_Q = Y'_u = Y'_d = \frac{1}{3}, \quad Y'_L = Y'_e = Y'_\nu = -1. \quad (3.51)$$

These mean that the charges for quarks are $+1/3$ and the charges for leptons are -1 . A baryon, for example, has $+1$ charge. This feature motivates labeling this group as $U(1)_{B-L}$. This is important to notice since the local symmetry of $B = 1/3$ and $L = 1$ alone leads to an anomalous theory. However, the global symmetry $B - L$ is anomaly-free, which implies that it is possible to have a gauge boson associated with this symmetry. An important process with this symmetry is a proton decay $p^+ \rightarrow \pi^0 e^+$, where B - and L - number conservation are violated but $B - L$ is preserved. This symmetry can also incorporate the addition of right-handed neutrinos, an extension that could explain the massive neutrino.

A more broader model of constructing the $U(1)'$ by considering a general form of the

charge to this extension with (Foot, 1993)

$$Y' = B + \alpha L_e + \beta L_\mu + \gamma L_\tau, \quad (3.52)$$

where the coefficient α, β, γ are arbitrary and L_e, L_μ, L_τ, B are the lepton number for electron, muon, tau, and baryon number, respectively. We focus here on the case where the baryon sector is left untouched. Adding right-handed neutrinos is required to guarantee vectorial symmetry of the $U(1)'$, making it possible to write a mass term for neutrinos. Assignments of charge in this $U(1)'$ extension of the SM for leptons, then

$$\begin{pmatrix} \nu_{eL} \\ e_L \end{pmatrix} \sim (1, 2, -1/2, \alpha), \quad e_R \sim (1, 1, -1, \alpha), \quad \nu_{eR} \sim (1, 1, 0, \alpha) \quad (3.53)$$

$$\begin{pmatrix} \nu_{\mu L} \\ \mu_L \end{pmatrix} \sim (1, 2, -1/2, \beta), \quad \mu_R \sim (1, 1, -1, \beta), \quad \nu_{\mu R} \sim (1, 1, 0, \beta) \quad (3.54)$$

$$\begin{pmatrix} \nu_{\tau L} \\ \tau_L \end{pmatrix} \sim (1, 2, -1/2, \gamma), \quad \tau_R \sim (1, 1, -1, \gamma), \quad \nu_{\tau R} \sim (1, 1, 0, \gamma). \quad (3.55)$$

An interesting part is when we consider only one right-handed neutrino is added, say for the electron case. The anomaly cancellation gives

$$[SU(2)_L]^2[U(1)'] \rightarrow (3)(3)Y'_Q + \alpha = 0 \quad (3.56)$$

$$[U(1)_Y]^2[U(1)'] \rightarrow (3)(3)\left(2(1/6)^2 - (2/3)^2 - (-1/3)^2\right)Y'_Q \quad (3.57)$$

$$+ \left(2(-1/2)^2 - (-1)^2\right)\alpha = 0. \quad (3.58)$$

Since $Y'_Q = 1/3$, we have $\alpha = -3$, and then this model is generally labeled as $U(1)_{B-3L_e}$. Similar steps for the other neutrino-flavor lead us to $U(1)_{B-3L_\mu}$ and $U(1)_{B-3L_\tau}$. Such scenarios enable us to particularly study how each neutrino flavor appears in the $E\nu$ ES and $CE\nu$ NS process we are interested in.

It is interesting to discuss the effect on the leptonic sector by the $U(1)'$ extension. Such models attract attention to be studied in the elastic neutrino-electron scattering process. It can be realized to consider the lepton numbers L_e, L_μ, L_τ . For this model, we can consider again

the Y' assignment with no B . The anomaly-free conditions of this extension satisfy

$$[SU(2)_L]^2[U(1)'] \rightarrow \alpha + \beta + \gamma = 0 \quad (3.59)$$

$$[U(1)']^3 \rightarrow \alpha^3 + \beta^3 + \gamma^3 = 0, \quad (3.60)$$

where the anomaly-free condition for gravitation is also considered (Foot, 1993). From these equations, the case where either α , β , or γ equals zero gives rise to an anomaly-free model. For example, the case $\gamma = 0$ gives $\alpha = -\beta$ and we have $Y' = L_e - L_\mu$. Following the same steps leads to $Y' = L_e - L_\tau$ and $Y' = L_\mu - L_\tau$. These are anomaly-free, which leads to phenomenological consequences of the possible existence of an additional Z' gauge boson for each model. Another important aspect is that, in contrast with previous $U(1)'$ models, the Z' boson could exist without the necessity for any exotic fermions such as the right-handed neutrino.

We focus in this work on the extra vector mediator Z' with an associated $U(1)'$ gauge group. The models are $U(1)_{B-L}$, $U(1)_{B-3L_i}$, and $U(1)_{L_i-L_j}$, with $i, j = e, \mu, \tau$ and $i \neq j$. The first model, $U(1)_{B-L}$, applies both for $E\nu$ ES and $CE\nu$ NS processes, while for $U(1)_{B-3L_i}$ we consider only for $CE\nu$ NS and $U(1)_{L_i-L_j}$ for $E\nu$ ES. These kind of models vary in terms of the fermion charges for each gauge group. The charges of each model, including the universal ones, are summarized in the Table 5. We can see that the mediator Z' has different

Table 5. The universal and flavor-dependent $U(1)'$ vector charges.

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

charges as it couples to either quarks, electrons, or neutrinos. This distinction specifies how each model contributes to the $E\nu$ ES and $CE\nu$ NS processes in the SM. The differential cross sections are derived within the framework of both processes by incorporating these charges. As a result, contributions of the extra Z' are coherently added to the NC of weak interactions

in the SM, mediated by the Z boson. Specifically, in the $U(1)_{L\mu-L\tau}$ scenario, no direct tree-level coupling to electrons, so its effects appear only at the loop level. That is, the Z' in this model contributes first at one-loop level to the $\nu_\mu - e^-$ and $\nu_\tau - e^-$ scattering. We can get the contributions by substituting in the standard formulation (Amaral *et al.*, 2020)

$$g_V \longrightarrow g_V - \frac{\sqrt{2}\alpha_{\text{em}}g_{Z'}^2(Q_{Z'}^\mu\delta_{\ell\mu} + Q_{Z'}^\tau\delta_{\ell\tau})}{\pi G_F(m_{Z'}^2 + 2m_e T_{er})}\epsilon_{\tau\mu}(|\vec{q}|). \quad (3.61)$$

It should be noted that there is a relative minus sign difference to the contributions from $\nu_\mu - e^-$ and $\nu_\tau - e^-$ scattering. Furthermore, the factor α_{em} represents the fine-structure constant and $\epsilon_{\tau\mu}$ takes the form

$$\epsilon_{\tau\mu}(|\vec{q}|) = \int_0^1 x(1-x) \ln \left(\frac{m_\tau^2 + x(1-x)|\vec{q}|^2}{m_\mu^2 + x(1-x)|\vec{q}|^2} \right) \approx \frac{1}{6} \ln \left(\frac{m_\tau^2}{m_\mu^2} \right). \quad (3.62)$$

3.2. Sterile Neutrino Dipole Portal

Future direct detection experiments of dark matter aim to observe recoils at extremely low energy thresholds. These experiments are expected to significantly enhance precision tests of $\text{E}\nu\text{ES}$ and $\text{CE}\nu\text{NS}$ within the SM, while also providing constraints on—or potential indications of—physics beyond the SM. One scenario is the occurrence of heavy sterile neutrinos from the active ones through a transition magnetic moment (Pontecorvo, 1957; Dasgupta & Kopp, 2021). The process is known as the Primakoff up-scattering (Primakoff, 1951), which was originally proposed to explain photoproduction of neutral meson within a nuclear electric field. A sizable quantity of transition magnetic moment increases the possibilities of producing sterile neutrinos via electromagnetic upscattering of neutrino beams on electron clouds or nuclei in the target material. To constrain this quantity, the sterile neutrino dipole portal has been widely investigated, utilizing data from a variety of experimental sources, including accelerator neutrino facilities (Bolton *et al.*, 2022), spallation neutron sources (De Romeri *et al.*, 2023), nuclear reactors (Kosmas *et al.*, 2017), and direct detection experiments (Sierra *et al.*, 2020).

The sterile neutrino dipole portal can be described by effective Lagrangian (Brdar *et*

al., 2020)

$$\mathcal{L}_{\text{int}} \supset \frac{\mu_{\nu_{\ell 4}}}{2} \bar{\nu}_{\ell L} \sigma^{\mu\nu} P_R \nu_4 F_{\mu\nu} + \text{h.c.}, \quad (3.63)$$

where $\mu_{\nu_{\ell 4}}$ represents the sterile neutrino dipole coupling of transition magnetic moment, ν_4 the sterile neutrino, $\nu_{\ell L}$ the SM (active) neutrino with flavor $\ell = e, \mu, \tau$, and $F_{\mu\nu}$ the tensor of the electromagnetic field. In this context, both neutrino states (active neutrino and sterile neutrino) may possess the nature of either Dirac or Majorana particles. We note that the above Lagrangian is valid only below the EW energy scale. Since the process of $E\nu\text{ES}$ and $\text{CE}\nu\text{NS}$ at DM-DD experiments occurs well below this scale, the Lagrangian remains applicable.

We show in Figure 16 Feynman diagrams that represent the presence of the dipole coupling of the active-sterile transition magnetic moment $\mu_{\nu_{\ell 4}}$ for (a) $E\nu\text{ES}$ and (b) $\text{CE}\nu\text{NS}$. We focus first for up-scattering process at $\text{CE}\nu\text{NS}$, $\nu_{\ell}(p_1) + \mathcal{N}(p_2) \rightarrow \nu_4(p_3) + \mathcal{N}(p_4)$. The

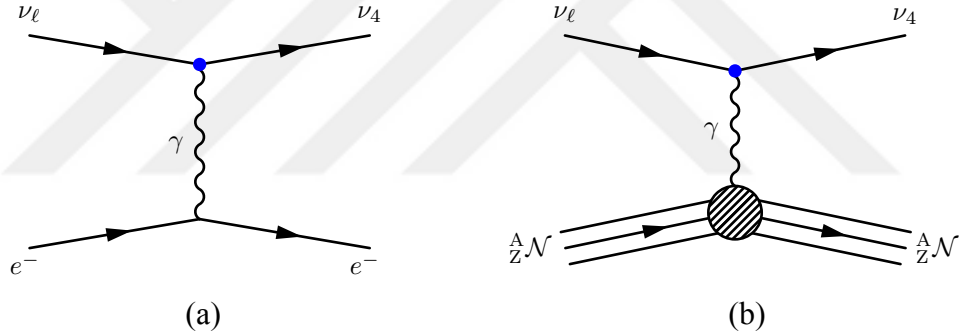


Figure 16. Feynman diagram representation for the up-scattering for (a) $E\nu\text{ES}$ and (b) $\text{CE}\nu\text{NS}$.

incoming neutrino ν_{ℓ} interacts with the nucleus $\mathcal{N}$ in detector material via photon exchange. The incoming active neutrino then upscattered into a sterile neutrino ν_4 . The amplitude for the process can be described by

$$-i\mathcal{M}_{\nu_4}^{\nu_{\ell}\mathcal{N}} = (\mu_{\nu_{\ell 4}})^* [-i\bar{\nu}_4(p_3)\sigma^{\mu\nu}P_L q_{\nu}\nu_{\ell}(p_1)] \left[\frac{-ig_{\mu\lambda}}{q^2} \right] [-ieZ\bar{N}(p_4)\gamma^{\lambda}N(p_2)F(|\vec{q}|^2)]. \quad (3.64)$$

We consider the spin-1/2 target nucleus in this work. Notice that we have used ν_4 to represent the spinor of the scattered sterile state. We emphasize that for the $E\nu\text{ES}$ case the hadronic current in the third parenthese is replaced by electron current $-i(e\bar{e}(p_4)\gamma^{\lambda}e(p_2))$. We can

obtain the analytical formula of the cross section by evaluating the absolute square of the amplitude in Equation (3.64). By averaging the absolute square of the amplitude, we have

$$\begin{aligned} \langle |\mathcal{M}_{\nu_4}^{\nu\mathcal{N}}|^2 \rangle &= \frac{|\mu_{\nu\ell_4}|^2 e^2 Z^2 F(|\vec{q}|^2)^2}{8q^4} \text{Tr} \left[(\not{p}_3 + m_4) \sigma_{\mu\rho} (1 - \gamma^5) \not{p}_1 \sigma_{\nu\lambda} \right] q^\rho q^\lambda \\ &\quad \times \text{Tr}[(\not{p}_4 + m_{\mathcal{N}}) \gamma^\mu (\not{p}_2 + m_{\mathcal{N}}) \gamma^\nu]. \end{aligned} \quad (3.65)$$

Performing the trace calculation, we obtain

$$\begin{aligned} \langle |\mathcal{M}_{\nu_4}^{\nu\mathcal{N}}|^2 \rangle &= \frac{|\mu_{\nu\ell_4}|^2 e^2 Z^2 F(|\vec{q}|^2)^2}{8q^4} \left[-16(p_1 \cdot p_3)(p_2 \cdot p_4)q^2 + 48(p_1 \cdot p_3)q^2 m_{\mathcal{N}}^2 \right. \\ &\quad - 32(p_1 \cdot p_3)(p_2 \cdot q)(p_4 \cdot q) + 32(p_3 \cdot q)(p_1 \cdot q)(p_2 \cdot p_4) \\ &\quad \left. - 64(p_3 \cdot q)(p_1 \cdot q)m_{\mathcal{N}}^2 \right]. \end{aligned} \quad (3.66)$$

By substituting the kinematics of Equation (A.8) and Equation (A.10) relevant in this process, we have

$$\begin{aligned} \left[\frac{d\sigma}{dT_{nr}} \right]_{\nu_4}^{\nu\mathcal{N}} &= \frac{\pi\alpha_{\text{EM}}^2}{m_e^2} \left| \frac{\mu_{\nu\ell_4}}{\mu_B} \right|^2 Z^2 F^2(|\vec{q}|^2) \left[\frac{1}{T_{nr}} - \frac{1}{E_\nu} - \frac{m_4^2}{2T_{nr}E_\nu m_{\mathcal{N}}} \right. \\ &\quad \left. \times \left(1 - \frac{T_{nr}}{2E_\nu} + \frac{m_{\mathcal{N}}}{2E_\nu} \right) - \frac{m_4^4}{8m_{\mathcal{N}}T_{nr}^2 E_\nu^2} \left(1 - \frac{T_{nr}}{m_{\mathcal{N}}} \right) \right], \end{aligned} \quad (3.67)$$

where subdominant terms have been neglected. The factor of $\pi\alpha_{\text{EM}}/m_e^2$ arises from the standard normalization using the Bohr magneton μ_B . We emphasize that this form is consistent with those reported by Bolton *et al.* (2022) and Li & Xia (2022). For $m_4 = 0$, the expression reduces to the conventional cross section for active neutrino magnetic moments, as discussed in the next section. Another important aspect is that the mass of the new sterile neutrino needs to satisfy the following constraint (De Romeri *et al.*, 2023)

$$m_4^2 \leq 2m_{\mathcal{N}}T_{nr} \left(\sqrt{\frac{2}{m_{\mathcal{N}}T_{nr}}} E_\nu - 1 \right). \quad (3.68)$$

We substitute $T_{nr} \rightarrow T_{er}$, $m_{\mathcal{N}} \rightarrow m_e$, and $Z^2 F^2(|\vec{q}|^2) \rightarrow Z_{\text{eff}}(T_{er})$ in the equation of cross section and the mass constraint above to obtain the corresponding quantities for the $\text{E}\nu\text{ES}$.

Correspondingly, we obtain the cross-section of the process as

$$\left[\frac{d\sigma}{dT_{er}} \right]_{\nu_4}^{\nu e} = \frac{\pi \alpha_{\text{EM}}^2}{m_e^2} \left| \frac{\mu_{\nu_{\ell 4}}}{\mu_B} \right|^2 \left[\frac{1}{T_{er}} - \frac{1}{E_\nu} - \frac{m_4^2}{2T_{er}E_\nu m_e} \left(1 - \frac{T_{er}}{2E_\nu} + \frac{m_e}{2E_\nu} \right) - \frac{m_4^4}{8m_e T_{er}^2 E_\nu^2} \left(1 - \frac{T_{er}}{m_e} \right) \right]. \quad (3.69)$$

3.3. Triplet Higgs Model

The Higgs triplet model (HTM) is constructed on a similar symmetry group to the SM (Arhrib *et.al.*, 2011). The only distinction is an extra triplet field, denoted by Δ , with hypercharge $Y = 1$, embedded into the Higgs sector of the SM, which contents already includes an isospin doublet field Φ ($Y = 1/2, L = 2$). Correspondingly, the Higgs sector of the model contains the SM Higgs $H \sim (1, 2, 1/2)$ and the extra $SU(2)_L$ scalar triplet $\Delta \sim (1, 3, 1)$. In general, the matrix representation of the model can be expressed as

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix} \quad (3.70)$$

The Yukawa interaction Lagrangian for leptons with the triplet Higgs is

$$\mathcal{L}_{\text{Yukawa}} = -f_{ij} L_i^T C i \sigma_2 \Delta L_j + \text{h.c.}, \quad (3.71)$$

where f_{ij} represents the complex and symmetric coupling matrix ($i, j = e, \mu, \tau$) with C denotes the Dirac charge conjugation operator, $L_i = (\nu_i, l_i)_L^T$ is a doublet of the left-handed lepton, and σ_2 is a Pauli matrix.

The most general $SU(2)_L \times U(1)_Y$ gauge invariant renormalizable Lagrangian in the scalar sector is

$$\mathcal{L} = (D_\mu H)^\dagger (D^\mu H) + \text{Tr}(D_\mu \Delta)^\dagger (D^\mu \Delta) - V(H, \Delta) + \mathcal{L}_{\text{Yukawa}}, \quad (3.72)$$

where the covariant derivatives satisfy

$$D_\mu H = \partial_\mu H + igT_a W_\mu^a H + i\frac{g'}{2}B_\mu H, \quad (3.73)$$

$$D_\mu \Delta = \partial_\mu \Delta + ig[T_a W_\mu^a, \Delta] + ig'\frac{Y_\Delta}{2}B_\mu \Delta, \quad (3.74)$$

where gauge fields and couplings are denoted by (W_μ^a, g) for the $SU(2)_L$ and (B_μ, g') for the $U(1)_Y$, while $T_a \equiv \sigma_a/2$ with $\sigma_a (a = 1, 2, 3)$ represents the Pauli matrices. Moreover, the potential $V(H, \Delta)$ is given by (Arhrib *et.al.*, 2011)

$$\begin{aligned} V(H, \Delta) = & -m_H^2 H^\dagger H + \frac{\lambda}{4}(H^\dagger H)^2 + m_\Delta^2 \text{Tr}(\Delta^\dagger \Delta) + [\mu(H^T i\sigma^2 \Delta^\dagger H) + \text{h.c.}] \\ & + \lambda_1(H^\dagger H)\text{Tr}(\Delta^\dagger \Delta) + \lambda_2(\text{Tr}\Delta^\dagger \Delta)^2 \\ & + \lambda_3 \text{Tr}(\Delta^\dagger \Delta)^2 + \lambda_4(H^\dagger \Delta \Delta^\dagger H), \end{aligned} \quad (3.75)$$

where the trace is over 2×2 matrices. The Yukawa sector of the SM plus one extra term is contained in $\mathcal{L}_{\text{Yukawa}}$. The additional term arises after spontaneous symmetry breaking, generating (Majorana) mass terms for the neutrinos without requiring right-handed neutrinos.

The simple form of the tree-level Majorana mass terms to explain massive neutrinos is the central motivation for exploring the HTM. A non-zero triplet vacuum expectation value (vev), $\langle \Delta^0 \rangle = v_\Delta/\sqrt{2}$, generate the mass matrix for neutrinos $m_{ij} = 2f_{ij}\langle \Delta^0 \rangle = \sqrt{2}f_{ij}v_\Delta$. Realistic neutrino masses may emerge with a perturbative f_{ij} by the condition $v_\Delta \lesssim 1$ eV. At tree level, the existence of v_Δ with a non-zero value leads to $\rho (= m_W^2/(m_Z^2 \cos\theta_W)) \neq 1$. Therefore, the requirement is that $v_\Delta \lesssim 1$ GeV to satisfy the $\rho \sim 1$ measurement result. The neutral Higgs triplet interaction term at tree level can be obtained from the Lagrangian

$$\mathcal{L} = -f_{ij}\nu_{iL}^T C \Delta^0 \nu_{jL} + \sqrt{2}f_{ij}\nu_{iL}^T C \Delta^+ l_{jL} + f_{ij}l_{iL}^T C \Delta^{++} l_{jL} + \text{h.c.} \quad (3.76)$$

Note that for $\nu^c(p) = C\bar{\nu}^T$ or $\bar{\nu}^c = \nu^T C$, implying that the first term in the Lagrangian can be expressed as $f_{ij}\bar{\nu}_{iL}^c \Delta^0 \nu_{jL}$, the second term as $\sqrt{2}f_{ij}\bar{\nu}_{iL}^c \Delta^+ l_{jL}$, and the third term as $f_{ij}\bar{l}_{iL}^c \Delta^{++} l_{jL}$. These terms indicate that both lepton flavor conserving (FC) and flavor violation (FV) processes can be described with this prescription. Notice that $\bar{\nu}_L^c = \bar{\nu}^c P_L$ from the properties of charge conjugation and chiral projection. With this notation, we can write the

interaction term of the Lagrangian as

$$\mathcal{L} = -f_{ij}\bar{\nu}_i^c P_L \Delta^0 \nu_j + \sqrt{2}f_{ij}\bar{\nu}_i^c P_L \Delta^+ l_j + f_{ij}\bar{l}_i^c P_L \Delta^{++} l_j + \text{h.c.} \quad (3.77)$$

In this dissertation, we study the contribution of the charged Higgs boson to $E\nu$ ES process using solar neutrinos. The suitable interaction term is the second term of the Lagrangian and its Hermitian conjugate in Equation (3.77).

We can see the Feynman diagram of the charged Higgs contribution to the $E\nu$ ES process in Figure 17. From Equation (3.77), the couplings of leptons to charged Higgs $H^\pm$ can

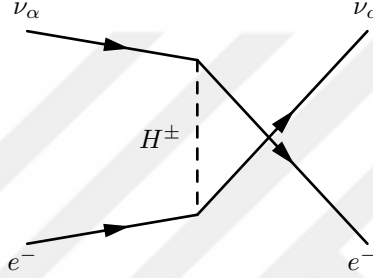


Figure 17. Feynman diagram representation of the charged Higgs contribution to $E\nu$ ES.

be expressed as

$$\mathcal{L} = -f_{ij}\sqrt{2}(l_i^T C P_L \nu_j + \nu_i^T C P_L l_j) \cos \theta^+ H^+ + \text{h.c.} \quad (3.78)$$

with the mixing angle $\cos \theta^+$ is approximated to unity. It is sourced by the suppression of v_Δ/v_0 from the ρ parameter, so that the mixing between the Higgs doublet and triplet is insignificance (Ong *et al.*, 2013). The amplitude of the charged Higgs contribution is

$$-i\mathcal{M}_{H^\pm} = [-if_{\alpha e}\bar{e}^c(p_4)P_L \nu_\alpha(p_1)] \left[\frac{-i}{q^2 - m_{H^\pm}^2} \right] [-if_{e\alpha}\bar{\nu}_\alpha^c(p_3)P_L e(p_1)]. \quad (3.79)$$

Simplifying this form by applying Fierz identities in Equation (A.50), we have

$$\mathcal{M}_{H^\pm} = \frac{f_{\alpha e}^2}{2m_{H^\pm}^2} [\bar{\nu}_\alpha(p_3)\gamma^\mu P_L \nu_\alpha(p_1)][\bar{e}(p_4)\gamma_\mu P_L e(p_2)] \quad (3.80)$$

This form is now relevant to the amplitude of the $E\nu$ ES process in the SM. We can express the total amplitude of the process with the charged Higgs correction as (Coarasa Perez *et al.*,

1995)

$$\mathcal{M}_{H^\pm} = \frac{G_F}{\sqrt{2}} \bar{\nu}_\alpha(p_3) \gamma^\mu (1 - \gamma^5) \nu_\alpha(p_1) \bar{e}(p_4) \gamma_\mu (g_{V_{H^\pm}} - g_{A_{H^\pm}} \gamma^5) e(p_2). \quad (3.81)$$

We have the couplings corrected by the the singly charged Higgs contribution to be

$$g_{V_{H^\pm}} = g_V - \frac{f_{\alpha e}^2}{m_{H^\pm}^2} \frac{1}{2\sqrt{2}G_F} \quad \text{and} \quad g_{A_{H^\pm}} = g_A - \frac{f_{\alpha e}^2}{m_{H^\pm}^2} \frac{1}{2\sqrt{2}G_F} \quad (3.82)$$

with $m_{H^\pm}$ is mass of the charged Higgs. Accordingly, taking the singly charged Higgs boson into consideration, we can obtain the total differential cross section by substituting $g_V \rightarrow g_{V_{H^\pm}}$ and $g_A \rightarrow g_{A_{H^\pm}}$ into Equation (2.12). We concentrate our work here only on the charged Higgs boson pure contribution. Therefore, we ignore the interference terms in this work. Since solar neutrinos oscillate and arrive as a mixture of three neutrino flavor, relevant coupling constants are f_{ee} , $f_{e\mu}$, or $f_{e\tau}$ in the presence of the charged Higgs boson.

4. ANALYSIS DETAILS

Signatures of new physics may become apparent through deviations from the background components in DM-DD experiments. We examine various BSM scenarios in light of recent results from CDEX-10, PandaX-4T, and XENONnT. This chapter provides a detailed description of our analysis aimed at constraining new physics parameters. We begin by outlining the methodology used to evaluate event rates arising from solar neutrino interactions. The treatment includes the implementation of detector efficiencies and energy smearing functions necessary for accurate rate predictions. Next, we describe the χ^2 analysis performed to extract constraints on the new physics parameters. Finally, we present the projected sensitivity of our approach, offering insight into the potential reach of future experiments under the considered scenarios.

4.1. Expected Event Rates

We calculate the expected event rates by convoluting the neutrino flux with the differential cross section. It is then adjusted to account for the specific properties of each experiment. Depending on the interaction channel under consideration, the detected signal may correspond to either electron recoil T_{er} or nuclear recoil T_{nr} . For a general recoil energy T_r , the predicted event rate can be expressed as

$$\left[\frac{dR}{dT_r} \right]^i = \int_{E_\nu^{\min}}^{E_\nu^{\max}} dE_\nu \frac{d\Phi^i}{dE_\nu} \left[\frac{d\sigma}{dT_r} \right], \quad (4.1)$$

where $d\Phi^i/dE_\nu$ is the flux with $i = pp, {}^7\text{Be}$, etc. refers to the solar neutrino flux main components, $d\sigma/dT_r$ is the differential cross-section, while E_ν denotes the neutrino energy. The neutrino energy has a minimum value

$$E_\nu^{\min} = \frac{1}{2} \left(T_r + \sqrt{T_r^2 + 2m_t T_r} \right), \quad (4.2)$$

where m_t is either electron or nucleus mass. For the kinematic end-points of each solar neutrino, we consider the maximum energy $E_{\nu,i}^{\max}$ in the solar neutrino fluxes. The minimum energy is the required initial energy to trigger the electron or nuclear recoil.

The predicted number of events expected in the experiments can be evaluated using

$$R^k = \varepsilon N_t \int_{T_r^k}^{T_r^{k+1}} dT'_r \mathcal{A}(T'_r) \int_0^{T_r^{\max}} dT_r \mathcal{R}(T'_r, T_r) \sum_{i=pp, {}^7\text{Be}, {}^8\text{B}} \left[\frac{dR}{dT_r} \right]^i. \quad (4.3)$$

Here, N_t represents the number of target nuclei per unit mass in the detector material and ε denotes the exposure factor. Notice that the reconstructed electron recoil energy is denoted by T'_r . As for the maximum recoil energy, we have

$$T_r^{\max} = \frac{2E_\nu^2}{2E_\nu + m_t}. \quad (4.4)$$

This relationship indicates that milder target nuclei improve the maximum recoil energy values in experiments. The detector efficiency is given by the factors $\mathcal{A}(T'_r)$ and the smearing functions by $\mathcal{R}(T'_r, T_r)$. Both are required to obtain the signal accurately. For the smearing function, we consider the Gaussian smearing function

$$\mathcal{R}(T'_r, T_r) = \frac{1}{\sigma(T_r)\sqrt{2\pi}} \exp \left[-\frac{(T'_r - T_r)^2}{2\sigma(T_r)^2} \right]. \quad (4.5)$$

The factor $\sigma(T_r)$ denotes detector resolutions that are different from experiments. The resolution for the CDEX-10 is (Xu *et al.*, 2022)

$$\sigma = 35.8 + 16.6\sqrt{T_r} \text{ (keV)}, \quad (4.6)$$

for the PandaX-4T (Zhang *et al.*, 2022)

$$\sigma = 0.073 + 0.173T_r - 0.0065T_r^2 + 0.00011T_r^3, \quad (4.7)$$

and for the XENONnT (Aprile *et al.*, 2022)

$$\sigma = 0.31\sqrt{\text{keV}}\sqrt{T_r} + 0.0037T_r. \quad (4.8)$$

Furthermore, we use the efficiency $\mathcal{A}(T'_r)$ from each experiment in deriving the BSM parameters. The CDEX-10 efficiency is shown in Figure 18, which we adapt from (Jiang *et. al.*, 2019). It is obtained by combining two efficiencies: the trigger and the physics-noise (PN)

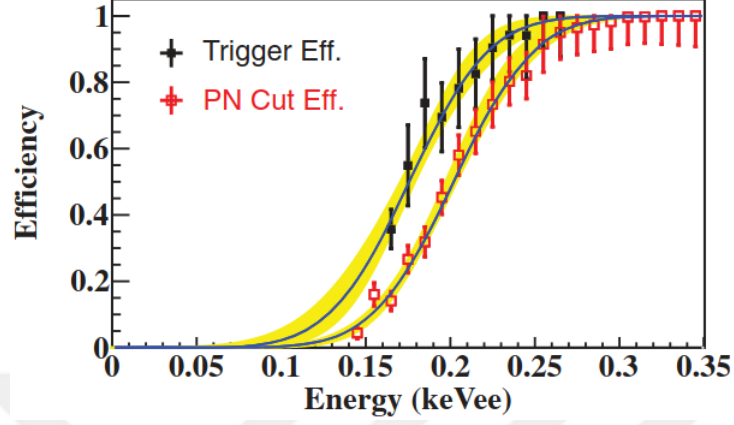


Figure 18. Efficiency of CDEX-10, taken from (Jiang *et. al.*, 2019).

cut. The two curves of efficiency were derived from the source of ^{137}Cs events. It is further fitted by an error function, as can be seen as a yellow band representing 1σ . For PandaX-4T, the efficiency is shown in Figure 19. The efficiencies for Run0 is shown by solid-line and for Run1 by dashed-line. Notice that there are also spectra of the theoretical (red) and

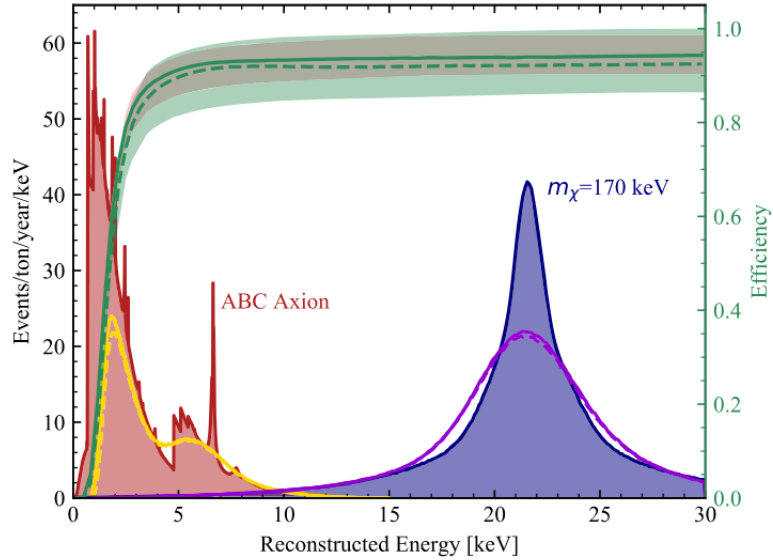


Figure 19. Efficiency of PandaX-4T, taken from (Zeng *et. al.*, 2025).

expected (yellow) energy spectra for the ABC axion and DM- ν conversion (blue) scenarios. The uncertainties for each efficiency of the dataset are shown by pink and green bands. We

consider these two efficiencies in our calculation. As for XENONnT, Figure 20 presents the combined detection and event selection efficiency of the experiment, including its associated uncertainties. The total efficiency is given by the black curve. It is a combination of two ef-

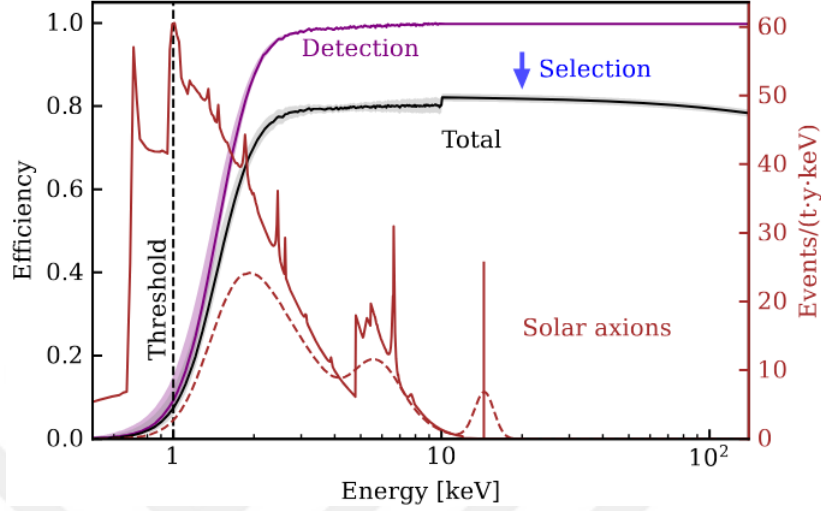


Figure 20. Efficiency of XENONnT, taken from (Aprile *et al.*, 2022).

ficiencies: the detection and event selection. Meanwhile, the detection efficiency dominated by the coincidence requirement of the S1 signal is represented by the purple curve (Aprile *et al.*, 2020). For a perspective, it is shown also the solar axion signal model before (red solid) and after incorporating energy smearing and efficiency loss (dashed lines).

4.2. Likelihood Analysis

We constrain the parameters of the models under consideration using the minimum χ^2 -method (Fieldman & Cousins, 1998). The published experimental data include uncertainties in the observed best-fit values. This technique allows us to place confidence level (C.L.) bounds on the model parameters. Unless stated otherwise, we focus on deriving 90% C.L. limits for each model throughout this work. The corresponding critical values of χ^2 depend on the number of degrees of freedom (d.o.f.) in the fit. These values are listed in Table 6. Since each model includes two free parameters in our analysis, the χ^2 thresholds are taken as follows:

- For 68% C.L., we use $\chi_{\min}^2 + 2.30$,

Table 6. Significance value and its corresponding degree of freedom.

C.L. (%)	degree of freedom				
	1	2	3	4	5
68.30	1.00	2.30	3.53	4.72	5.89
90.00	2.71	4.61	6.25	7.78	9.24
95.40	4.00	6.17	8.02	9.70	11.3
99.00	6.63	9.21	11.3	13.3	15.1
99.73	9.00	11.8	14.2	16.3	18.2
99.99	15.1	18.4	21.1	23.5	25.7

- For 90% C.L., we use $\chi_{\min}^2 + 4.61$, and so on, according to standard $\chi_{\min}^2$ distribution properties.

Apart from their primary goal of detecting DM candidates, data from these experiments have also been employed to explore various new physics scenarios, such as light mediator models, dark photons, and neutrino electromagnetic properties. As part of the background in these detectors, $E\nu$ ES and $CE\nu$ NS contribute predictable and well-understood signals. These backgrounds are typically subtracted in standard DM analyses. However, the presence of beyond the SM interactions may alter these signals, providing a promising framework to investigate potential effects of BSM physics.

We analyze experimental data using electron recoil measurements from XENONnT (Aprile *et al.*, 2022) and PandaX-4T (Zeng *et al.*, 2025), as well as nuclear recoil data from CDEX-10 (Geng *et al.*, 2023), employing a binned analysis approach. The CDEX-10 collaboration has reported results on both neutrino-electron and neutrino-nucleus interactions using p -type contact germanium detector. The collaboration explores various BSM scenarios accordingly. The data taking began in February 2017, with a physics analysis threshold is 160 eVee. The PandaX-4T recently published their unblinding Run0 and Run1 data with electron recoil < 30 keV. The Run0 data was obtained after operating from 28 November 2020 to 16 April 2021, while the Run1 data was obtained from 16 November 2021 to 15 May 2022. The two facilities are part of the China Jinping Underground Laboratory (CJPL). Meanwhile, the XENONnT is the latest detector program by the XENON collaboration. The update has a 5.9 tonne of liquid xenon active target. The electron recoil SR0 dataset, collected from 6 July 2021 to 10 November 2021, was published by the experiment with no excess above the

background (Aprile *et al.*, 2022).

In order to obtain sensitivity of the new physics parameter(s) $\mathcal{S}$, we apply a Gaussian χ^2 function for statistical analysis (Baker & Cousins, 1984; Fogli *et al.*, 2002) as follows:

$$\chi^2(\mathcal{S}) = \min_{(\alpha_i, \beta_i)} \left[\sum_{k=1}^{30} \left(\frac{R_{\text{exp}}^k(\mathcal{S}; \alpha, \beta) - R_{\text{obs}}^k}{\sigma^k} \right)^2 + \sum_i \left(\frac{\alpha_i}{\sigma_{\alpha_i}} \right)^2 + \sum_i \left(\frac{\beta_i}{\sigma_{\beta_i}} \right)^2 \right]. \quad (4.9)$$

Here, the R_{obs}^k represents the observed event rate and R_{exp}^k the expected event rate, both in the k -th energy bin. The expected rate is decomposed with R_{SM}^k added with $R_{\text{BSM}}^k(\phi, Z', \dots)$ and R_{Bkg} . The latter represents other background components. The explicit form is calculated by

$$R_{\text{exp}}^k(\mathcal{S}; \alpha, \beta) = (1 + \alpha)(R_{\text{SM}}^k + R_{\text{BSM}}^k(\mathcal{S})) + (1 + \beta)R_{\text{Bkg}}^k \quad (4.10)$$

In these expressions, α denotes the nuisance parameters associated with uncertainties in the solar neutrino flux, while β denotes the normalization factor for the background component. We use background data provided by PandaX-4T (Zeng *et al.*, 2025) and XENONnT (Aprile *et al.*, 2022). For CDEX-10 (Geng *et al.*, 2023), we assume a constant background of 1 kpcc. The experimental uncertainty is represented by σ^k . Meanwhile, σ_{α} denotes the solar neutrino flux uncertainty, taken from Table 6 of Vinyoles *et al.* (2017), and σ_{β} denotes the experimental background. The SM with new physics contributions is evaluated by using the solar neutrino fluxes from Bahcall's spectrum (Bahcall, 1989), normalized according to the B16-GS98 SSM (Vinyoles *et al.*, 2017). By adding all the energy bins, we obtain total neutrino-electron as well as neutrino-nucleus event rates that are in good agreement with experimental results.

4.3. Projected Sensitivity

Direct detection experiments are advancing into the multiton scale, offering enhanced sensitivity to $\text{CE}\nu\text{NS}$ and enabling the exploration of a wide range of new physics scenarios. Xenon-based detectors have already demonstrated the capability to detect $\text{CE}\nu\text{NS}$ signals from solar neutrinos, as shown by PandaX-4T (Bo *et al.*, 2024) and XENONnT (Aprile *et al.*, 2024). Looking forward, the DARWIN project (Aalbers *et al.*, 2016) is expected to operate with a liquid xenon target mass of approximately 50 tons, significantly improving sensitivity

to both SM and BSM signals. In parallel, next-generation low-threshold detectors utilizing solid-state targets, such as silicon or germanium, are under development. These technologies offer complementary sensitivity, particularly in the low-energy recoil regime, further expanding the discovery potential of future experiments. Examples of such next-generation detectors include EDELWEISS (Armengaud *et al.*, 2022), SuperCDMS (Agnese *et al.*, 2017), and SENSEI (Barak *et al.*, 2020). These experiments are specifically designed to improve sensitivity to low-mass WIMPs and are capable of detecting extremely low-energy recoils induced by solar neutrinos. As a result, they offer promising avenues for placing more stringent constraints on BSM physics.

Currently, the CDEX collaboration is progressing toward the CDEX-50 phase (Geng *et al.*, 2024), following the completion of its CDEX-1 and CDEX-10 stages. In its third phase, the experiment plans to deploy a detector comprising a 50 kg array of high-purity germanium. This setup targets to reach a background level of about $0.01 \text{ events keV}^{-1}\text{kg}^{-1}\text{day}^{-1}$, with a $150 \text{ kg} \cdot \text{year}$ target exposure and a 160 eVee analysis threshold (Geng *et al.*, 2024). Ultimately, the experiment seeks to establish a Ge detector in a ton-scale, relying on PCGe technology (Kang *et al.*, 2013). It is anticipated to operate in the near future with an energy threshold in a sub-keV regime. Such an improvement is expected to enhance the sensitivity to detect solar neutrino interactions, potentially providing measurements with high precision.

Building on the motivations discussed above, we configure several scenarios to explore the potential of upcoming experimental advancements. For the new light mediator models in the $\text{CE}\nu\text{NS}$ process, we consider two possibilities. The first is the realistic scenario, assuming a 10% achievable uncertainty. The second configuration is termed the optimistic scenario, reduces the uncertainty down to 1.5%. Such low uncertainties are anticipated through improvements in future experiments regarding detector material purity, tightly controlled exposure, and enhanced shielding. In both cases, a flat background is assumed based on these advancements. The scenarios are applied within the expected ROI of future recoil energies. In this case, the nuclear recoils are assumed to reach about 0.1 keV in the realistic scenario and 0.015 keV in the optimistic one. Furthermore, for the sterile neutrino dipole portal, we define two categories. These are next-generation and future scenarios, with the latter further divided into future 1 and future 2. In the next-generation scenario, the exposure is set to $150 \text{ kg} \cdot \text{year}$, whereas both future scenarios assume an exposure of $1.5 \text{ ton} \cdot \text{year}$. The

expected energy threshold is 1 keVnr for the next-generation and future 1 cases, and it is 0.1 keVnr for the future 2 case. We assume a 50 kg target mass in all scenarios. In all cases, we assume a flat background similar to the new light mediator case. These are expected to lower the uncertainty for the next-generation scenario to about 10% and for the future scenarios to about 1%. Based on these projections, we study the connection between bounds of the sterile neutrino mass and experimental precision, which potentially provides an intuitive scaling for future direct detection progress.



5. RESULTS AND DISCUSSIONS

In this chapter, we present the derived constraints on three classes of BSM scenarios: light mediator models, the sterile neutrino dipole portal, and the charged Higgs boson. Our primary goal is to identify potential deviations from the Standard Model (SM) expectations in $\text{CE}\nu\text{NS}$ and $\text{E}\nu\text{ES}$ processes. To this end, we analyze recent low-energy electronic recoil data from XENONnT, PandaX-4T (Run 0 and Run 1), and CDEX-10. We begin by presenting the predicted event rates for both the SM and the BSM scenarios under consideration. Subsequently, we provide detailed results for each scenario, together with comparisons to existing limits reported in previous studies.

5.1. Predicted Event Rates

We consider Ge and Xe as the target materials to probe new physics scenarios within the $\text{E}\nu\text{ES}$ and $\text{CE}\nu\text{NS}$ framework. It is important to note that $\text{CE}\nu\text{NS}$ event rates scale with the square of the number of neutrons in the target nucleus, while $\text{E}\nu\text{ES}$ rates depend on the effective electron charge, which varies depending on the atomic structure of the material. In our analysis, we consider all stable isotopes of both germanium and xenon to accurately account for the material-dependent effects on the interaction rates.

The cross-sections for the two processes are presented in Figure 21. We set the initial neutrino energy to be 10 MeV, which is typical for solar neutrino fluxes. It can be inferred that the heavier target provides a higher spectrum. We can also see the effect of neutron number in the $\text{CE}\nu\text{NS}$ process. Meanwhile, the stepped shape in the $\text{E}\nu\text{ES}$ process shows the effective electron charge effect that improved for the outer nuclear shell.

We now evaluate the predicted event rates. Taking into account all relevant solar neutrino flux components, the expected event rates are presented in Figure 22 for $\text{E}\nu\text{ES}$ and in Figure 23 for $\text{CE}\nu\text{NS}$. Only specific components of the solar neutrino flux contribute significantly to the processes under consideration. In the case of $\text{E}\nu\text{ES}$, the dominant contributions arise from pp and ${}^7\text{Be}$ neutrinos. In contrast, ${}^8\text{B}$ and hep neutrinos are the primary contribu-

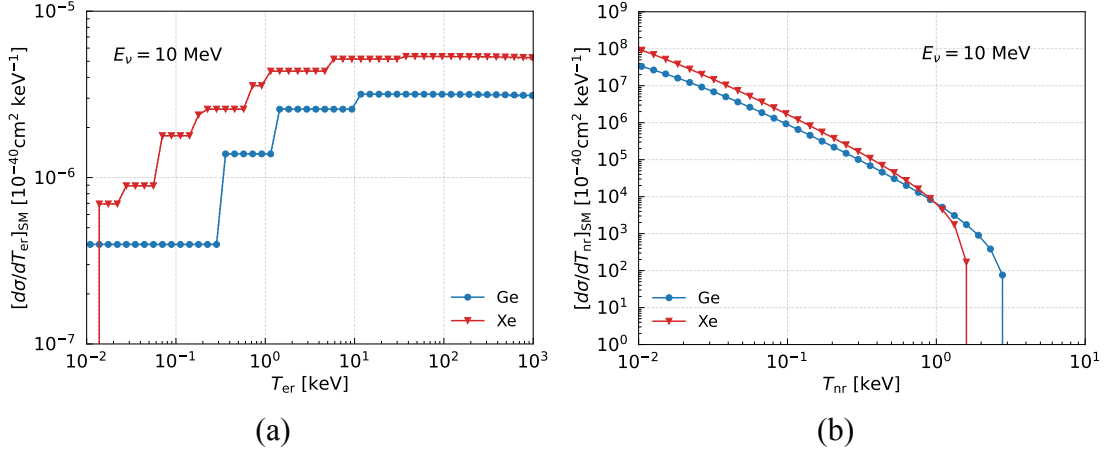


Figure 21. Differential cross-sections for (a) $E\nu$ ES and (b) $CE\nu$ NS with germanium and xenon targets.

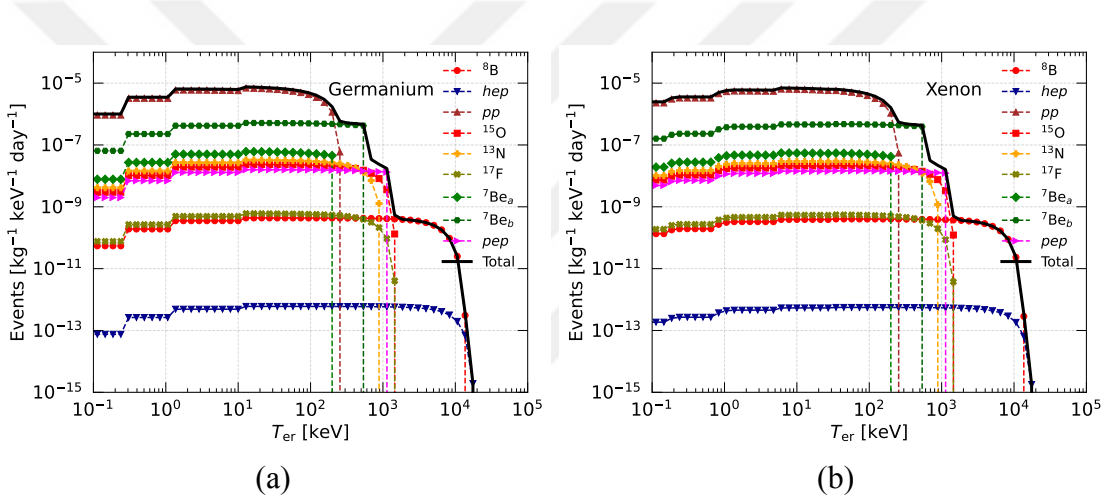


Figure 22. Predicted event rates of $E\nu$ ES using (a) germanium and (b) xenon target.

tors in the $CE\nu$ NS process due to their higher energies. These relevant flux components are incorporated into our calculations when deriving parameter constraints for each BSM model.

We present the predicted event rates in the presence of new physics for the CDEX-10 in Figure 24. We show contributions from a new light mediator model, the sterile neutrino dipole portal, and the charged Higgs boson. We can see that data from the CDEX-10 is shown as a dotted point. We have chosen the benchmark for each model to show the possible effect of these BSM scenarios. In the $E\nu$ ES, we set $g_{Z'} = 2 \times 10^{-6}$, $m_{Z'} = 0.01$ MeV for the universal light vector mediator, $\mu_{\nu_{\ell 4}} = 3 \times 10^{-9}$, $m_4 = 0.01$ for the dipole portal, and $f_{ae}^2/m_{H^\pm}^2 = 5 \times 10^{-3}$ for the charged Higgs boson. Meanwhile, in the $CE\nu$ NS, we set $g_{Z'} = 9 \times 10^{-5}$, $m_{Z'} = 0.01$ MeV for the universal light vector mediator and $\mu_{\nu_{\ell 4}} =$

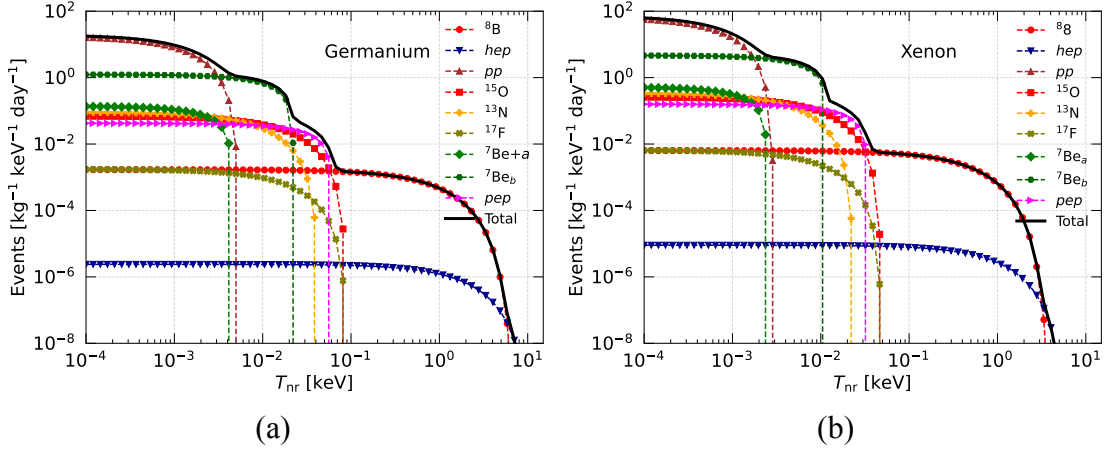


Figure 23. Predicted event rates of CE ν NS using (a) germanium and (b) xenon target.

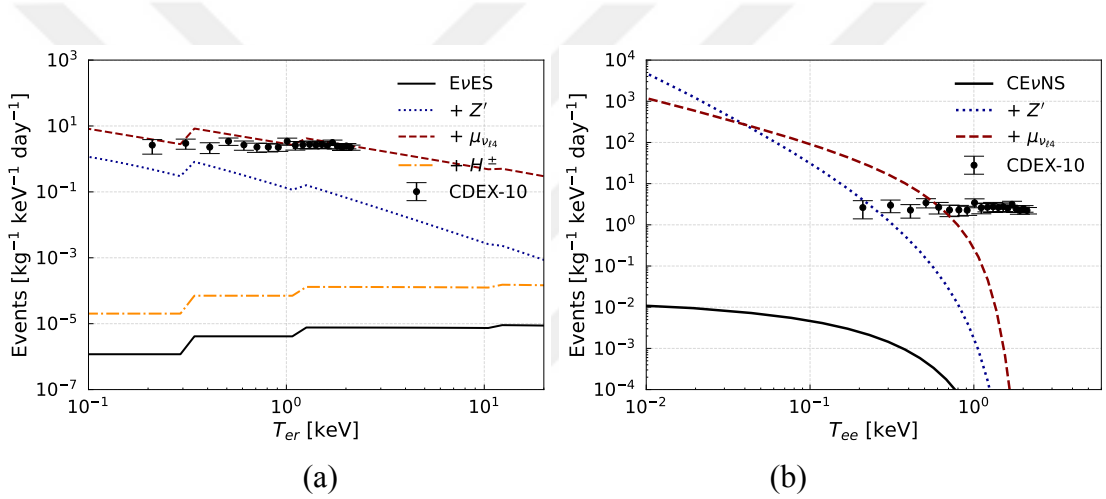
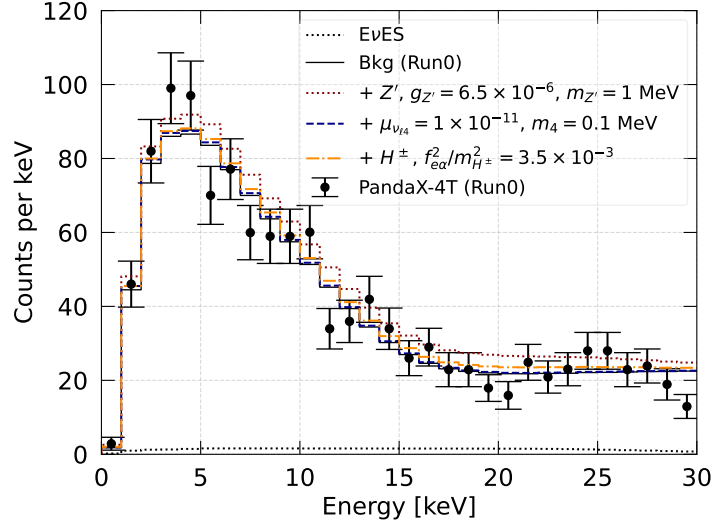


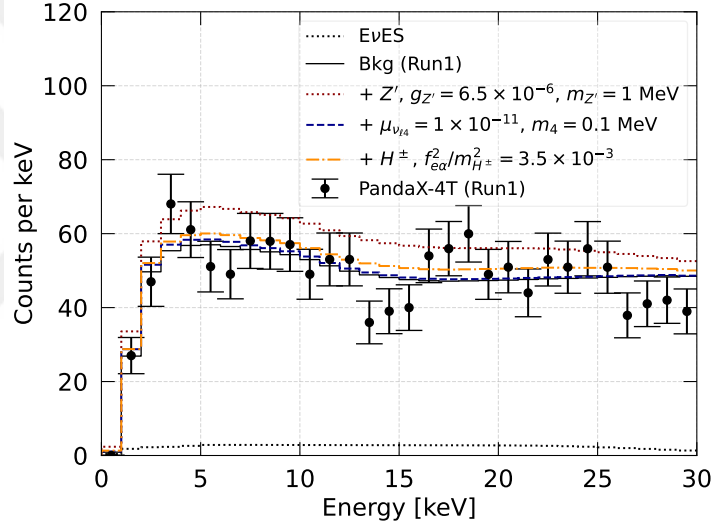
Figure 24. Event rate predictions from CDEX-10 results for (a) E ν ES and (b) CE ν NS with contribution from new physics.

$1 \times 10^{-7} m_4 = 0.01$ for the dipole portal. We can infer that the CDEX-10 data have low sensitivity to the CE ν NS process, while the data potentially indicate the new physics effect to the E ν ES process in the considered ROI of the experiment.

We next show the predicted event rates for the Run0 and Run1 datasets of PandaX-4T in Figure 25. We present contributions from a universal light vector mediator model, active-sterile neutrino transition magnetic moment, and charged Higgs boson. Benchmark values for the models are $g_{Z'} = 6.5 \times 10^{-6}$, $m_{Z'} = 1$ MeV, $\mu_{\nu\ell_4} = 1 \times 10^{-11} \mu_B$, $m_4 = 0.1$ MeV, and $f_{\alpha e}^2/m_{H^\pm}^2 = 5 \times 10^{-3}$, respectively. For both Run0 and Run1 data sets, we can clearly observe enhancement of signals due to these scenarios.



(a)



(b)

Figure 25. Event rate predictions from PandaX-4T (a) Run0 and (b) Run1 with contribution from light mediator and transition magnetic moment.

For the XENONnT, we present the predicted event rates in Figure 26. Similar to the previous case, we present our calculated spectra of the SM $E\nu$ ES with a universal light mediator model, active-sterile neutrino transition magnetic moment, and charged Higgs boson. In the evaluation, we set $g_{Z'} = 1 \times 10^{-6}$ with $m_{Z'} = 1$ MeV, $\mu_{\nu_{\ell 4}} = 2 \times 10^{-10} \mu_B$ with $m_4 = 0.1$ MeV, and $f_{\alpha e}^2/m_{H^\pm}^2 = 2.5 \times 10^{-3}$. It is observed that applying these benchmarks to the models enhances the predicted event rates from the experiment, enabling to reach of high-value counts that unable to be reached by solely considering the SM.

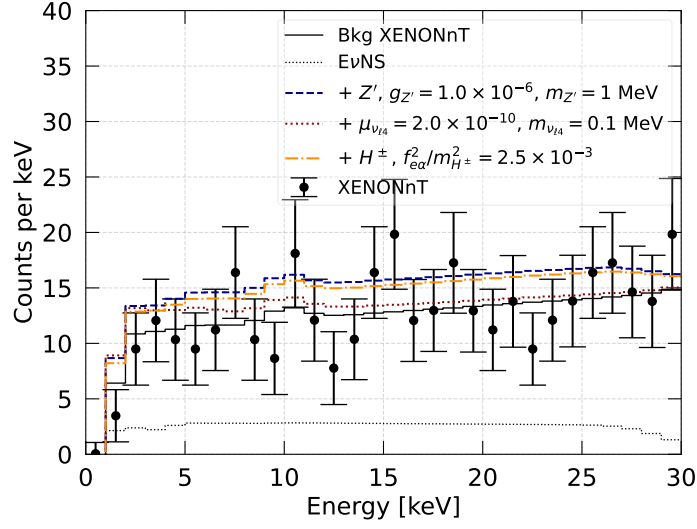


Figure 26. Event rate predictions from XENONnT electronic recoil signals with contribution from light mediator and transition magnetic moment are shown.

5.2. Previous Limits

Several previous studies have placed constraints on the parameters considered in this work, using neutrino data from both artificial and natural sources. Artificial sources include neutron spallation facilities and nuclear reactors, while natural sources involve solar neutrino experiments and DM-DD experiments. Some also investigated neutrinos from supernovae, stellar cooling, and in the study of cosmology. In addition, there are projected results from planned facilities. In this study, we superimpose our results onto the existing limits and compare them within the context of the specific BSM scenarios under investigation

For the neutron spallation source, we consider the COHERENT result, which first announced the discovery of the $CE\nu NS$ process (Akimov *et al.*, 2016, 2017). Experiments who use neutrino fluxes from nuclear reactors are TEXONO (Deniz *et al.*, 2010), CONNIE (Aguilar-Arevalo *et al.*, 2020), Dresden-II (Colaesi *et al.*, 2022), and CONUS (Bonet *et al.*, 2020). In these reactor experiments, neutrinos produced are electron-antineutrinos and relevant for the $E\nu ES$ process. Meanwhile, BOREXINO (Bellini *et al.*, 2010) was exclusively constructed to search for solar neutrino interaction with electrons. Regarding DM-DD experiments, there are the LZ (Akerib *et al.*, 2013), XENON (Aprile *et al.*, 2017) and PandaX (Zhang *et al.*, 2019). Limits from collider studies such as Mainz (Merkel *et al.*, 2011), KLOE (Abelev *et al.*, 2013), BaBar (Lees *et al.*, 2014), PHENIX (Adare *et al.*, 2015), NA48/2 (Bat-

ley *et al.*, 2015). Results from oscillation data (Coloma *et al.*, 2021) and the allowed bound of $(g - 2)_\mu$ (Aguilarrrd *et al.*, 2023) are also presented. Moreover, predicted results from future facilities such as DUNE (Abi *et al.*, 2020), NUCLEUS (Rothe *et al.*, 2019), FLArE (Batell *et al.*, 2021), and SHiP (Alekhn *et al.*, 2016) are also included. For complementary, astrophysical limits (Big Bang Nucleosynthesis (Blinov *et al.*, 2019), supernova (Croon *et al.*, 2021), cosmology (Escudero *et al.*, 2019), and stellar cooling (Li & Xu, 2023)) are also presented. Limits derived from previous results are either for $E\nu$ ES or $CE\nu$ NS. Additional details on other limits are mentioned in the subsequent discussions.

5.3. New Light Mediator Models

In this section, we present the exclusion regions derived for various new light mediator models. We begin by discussing the results for universal mediator scenarios, including scalar, vector, and tensor mediators that couple identically to all fermion species. Following this, we examine a class of anomaly-free $U(1)'$ models, which introduce mediator couplings based on specific charge assignments. The results are organized by interaction type: we first present the constraints obtained from electron recoil data, followed by those derived from nuclear recoil observations.

5.3.1. Results from Electron Recoil Data

The exclusion limits at 90% C.L. derived from the analysis of PandaX-4T data for the universal scalar mediator scenario are presented in Figure 27. We observe from Figure 27 (a) as $m_\phi \leq 1$ keV Run0 reaches $g_\phi \lesssim 1.1 \times 10^{-6}$ and Run1 reaches $g_\phi \lesssim 8.3 \times 10^{-7}$. The latter limit provides a more sensitive limit than the Run0. It is around 1.3 times the limit from the Run0 data. It is anticipated that this behavior comes from the larger exposure of the Run1. Next, our results are then superimposed with previous bounds for this model in Figure 27 (b). We observe that results from Run0 and Run1 datasets indicate more sensitive constraints in the low-mass region than limits from COHERENT (CsI+LAr), CONNIE (95% C.L.), CONUS+ ($CE\nu$ NS+ $E\nu$ ES), Dresden-II, BOREXINO, and CDEX-10 ($CE\nu$ NS), as well as, for the whole region, more sensitive than the limit derived from CONUS ($E\nu$ ES).

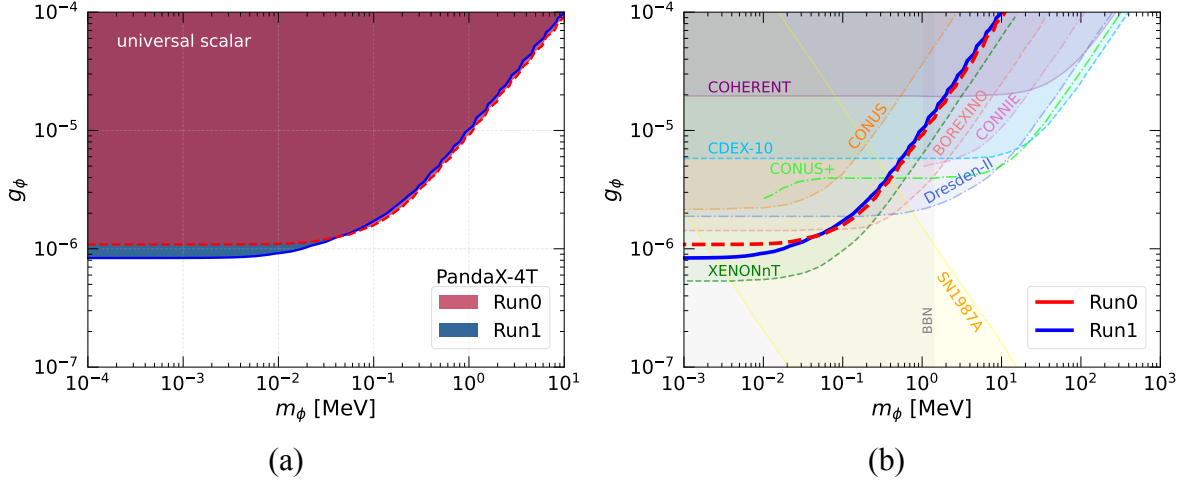


Figure 27. (a) Exclusion limits on scalar mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

Meanwhile, limits obtained from the PandaX-4T two datasets have an identical order of magnitude value to the one from XENONnT. However, the obtained limits have yet to give an improvement relative to this limit. For complementarity, the derived limits cover the limit from BBN for $m_\phi > 1.4$ MeV and the SN1987A in the low-mass region.

Moving on to the universal vector mediator, our derived limits at 90% C.L. are presented in Figure 28. We observe from Figure 28 (a) that $g_{Z'} \lesssim 5.1 \times 10^{-7}$ for Run0 data and $g_{Z'} \lesssim 2.4 \times 10^{-7}$ for Run1 data, when $m_{Z'} \leq 1$ keV. In the low-mass region, there is about 2.1 times enhancement on the limit from Run1 data over the Run0 data. Next in

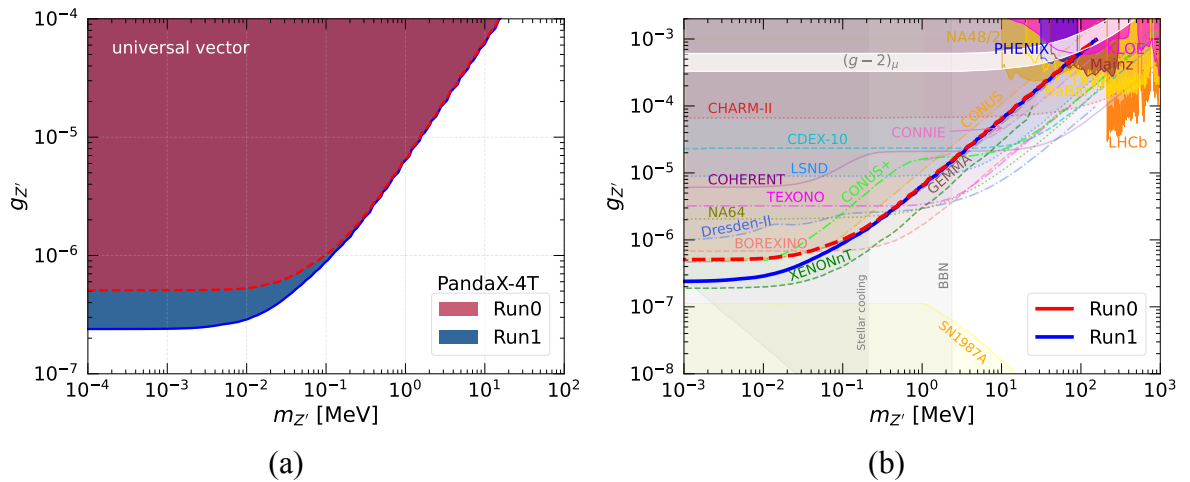


Figure 28. (a) Exclusion limits on vector mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

Figure 28 (b), we present the superimposed plot with available bounds from analyzing other experiments mentioned previously. It can be seen that the derived limits at low-mass regions provide more stringent results than most of the previous results, which emphasizes the novelty of our approach. Particularly, in the region $m_{Z'} \lesssim 4$ MeV, our obtained limits completely cover results of COHERENT (CsI+LAr), indicating about one order of magnitude improvements. We also observe that the constraints we obtain are tighter than those previously derived from reactor power plant (CONNIE (95% C.L.), TEXONO, Dresden-II, and CONUS+), other DM-DD facility (CDEX-10 with CE ν NS), and other experiments regarding electron-neutrino scattering (CHARM-II, LSND, and NA64). Specifically, the limit from the Run0 data extends along limits from GEMMA and CONUS (E ν ES). Meanwhile, a few times improvement is seen regarding the limit from the Run1 relative to those results. For $m_{Z'} \leq 70$ keV, the limit from the Run0 provides about 25% while the limit from the Run1 65% improvement over the limit from the BOREXINO. Moreover, the obtained limits are yet to reach the limit from the XENONnT, but possess the same order of magnitude. Note that for $m_{Z'} > 100$ MeV, there is a small region of the $(g - 2)_\mu$ allowed bound yet to be excluded by the obtained limits. We can also see that our results are to complement limits from collider experiments, stellar cooling, SN1987A, and BBN.

We then discuss the universal light tensor mediator model. We present our limits at 90% C.L. in Figure 29. In Figure 29 (a), we observe for $m_T \leq 1$ keV, $g_T \lesssim 2.8 \times 10^{-7}$ for the Run0 data and $g_T \lesssim 1.4 \times 10^{-7}$ for the Run1 data. Concerning the low mass region, the limit from the Run1 indicates an improvement of around 2 times over the limit from the Run0 data. We further provide in Figure 29 (b) a comparison of our derived limits with the available ones from previous studies. We can see that our limits cover the majority of bounds derived using COHERENT (CsI+LAr) and CDEX-10 datasets, corresponding to about one and two orders of magnitude improvements, respectively. Furthermore, compared to the limit from XENONnT, our results have the same order of magnitude but have yet to reach it.

The next discussion is about the new light mediator from the anomaly-free $U(1)'$ models, namely the ones obtained from $U(1)_{L_e-L_\mu}$, $U(1)_{L_e-L_\tau}$, and $U(1)_{L_\mu-L_\tau}$ gauge symmetry. Concerning the 90% C.L. exclusion limits, results of the $L_e - L_\mu$ is shown in Figure 30, the $L_e - L_\tau$ in Figure 31, and the $L_\mu - L_\tau$ in Figure 32. We see in Figure 30 (a) that $g_{L_e-L_\mu} \lesssim 8.2 \times 10^{-7}$ from the Run0 data and $g_{L_e-L_\mu} \lesssim 3.9 \times 10^{-7}$ from the Run1 data,

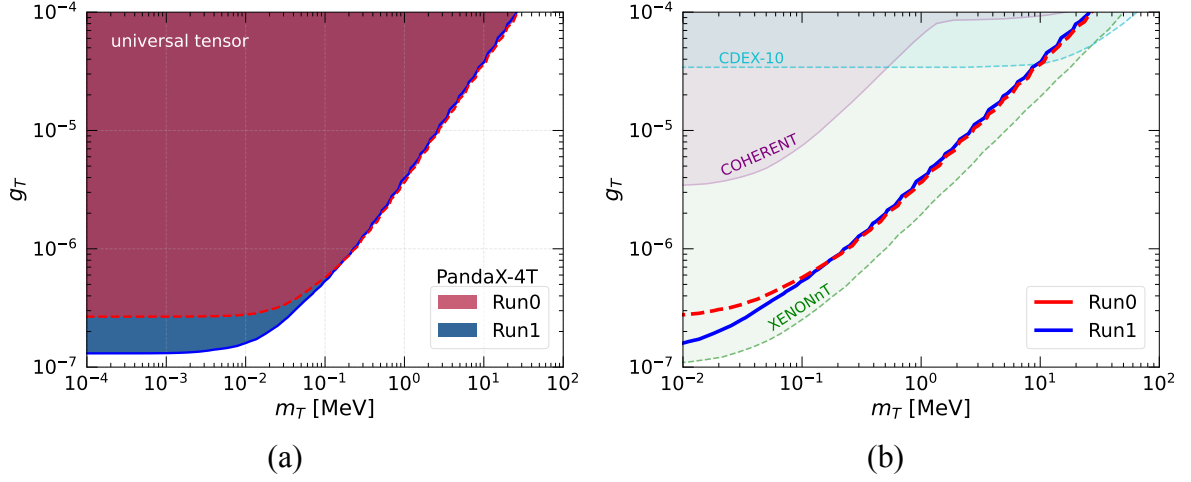


Figure 29. (a) Exclusion limits on tensor mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

indicating that the Run1 limit is more stringent of about 2.1 times over the Run0 limit. For

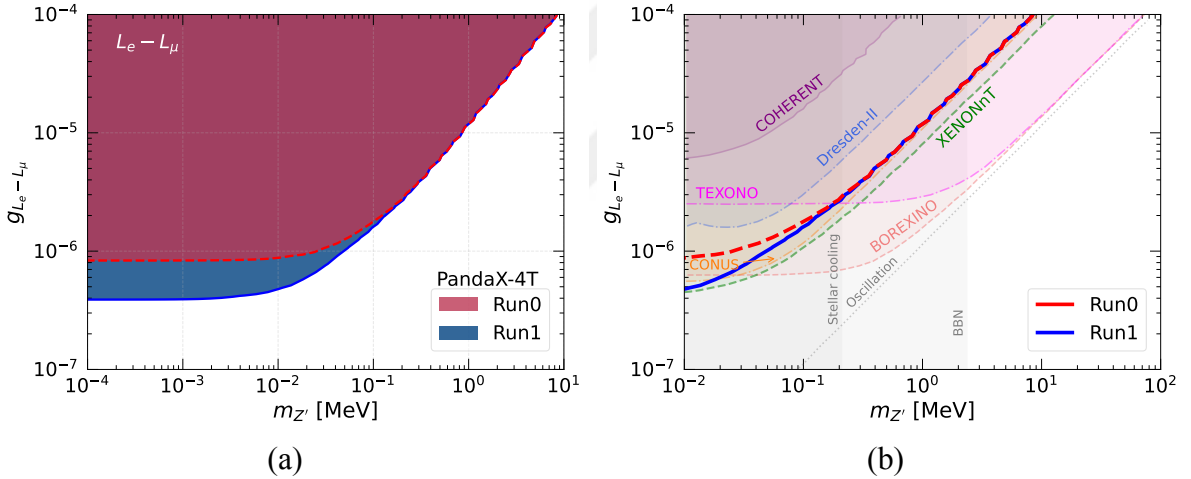


Figure 30. (a) Exclusion limits on $U(1)_{L_e-L_\mu}$ mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

$g_{L_e-L_\tau}$, we see in Figure 31 (a) that $g_{L_e-L_\tau} \lesssim 1.2 \times 10^{-6}$ for the Run0 and $g_{L_e-L_\tau} \lesssim 3.8 \times 10^{-7}$ for the Run1. The limit from Run1 is around 3.2 times more stringent than the limit from Run0. As for the $g_{L_\mu-L_\tau}$, we see in Figure 32 (a) that $g_{L_\mu-L_\tau} \lesssim 6.0 \times 10^{-6}$ for the Run0 data and $g_{L_\mu-L_\tau} \lesssim 3.3 \times 10^{-6}$ for the Run1 data. We find that the Run1 limit is about 1.8 times more stringent than the Run0 limit.

We further evaluate new limits using the XENONnT recent electronic recoil data (Aprile *et al.*, 2022) for each new light mediator from anomaly-free $U(1)'$ models. We emphasize

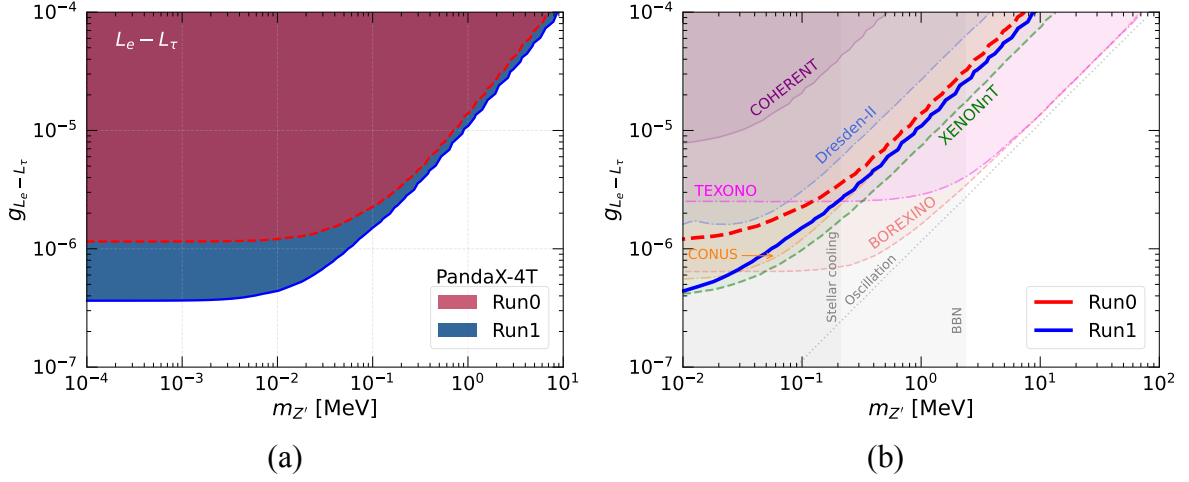


Figure 31. (a) Exclusion limits on $U(1)_{L_e-L_\tau}$ mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

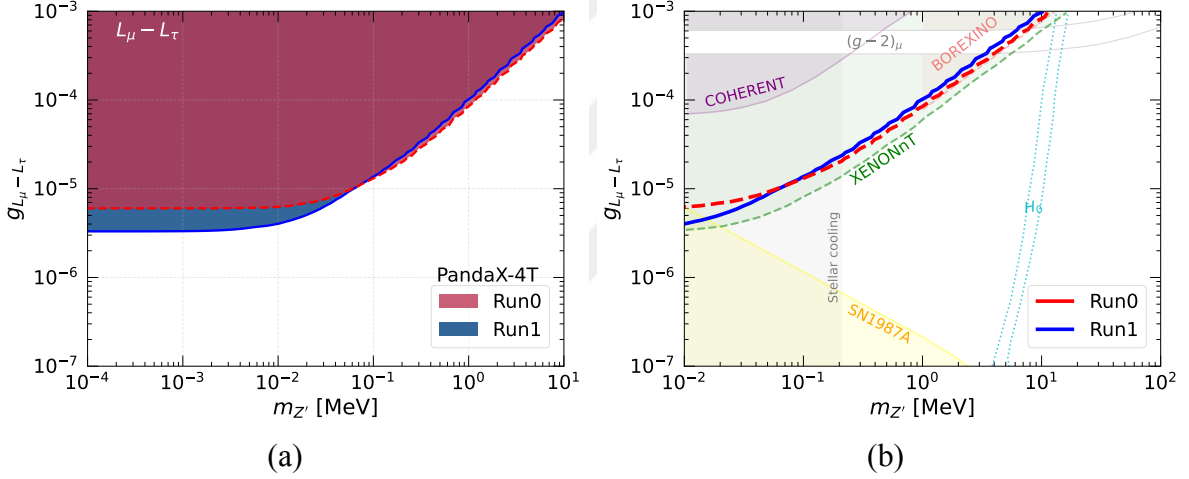


Figure 32. (a) Exclusion limits on $U(1)_{L_\mu-L_\tau}$ mediator using PandaX-4T Run0 and Run1 data with E ν ES and (b) comparison with previous limits.

that, for the first time, limits for $U(1)_{L_e-L_\mu}$ and $U(1)_{L_e-L_\tau}$ models are derived using this data. Results for $L_\mu - L_\tau$ have been derived previously in (Melas *et al.*, 2023), and we find compatible results with this work. Notice that the limits of the considered scenarios have been evaluated by combining datasets from DM-DD experiments in (De Romeri *et al.*, 2024). However, it did not present the limits separately for XENONnT. To accommodate this, we apply the same analysis procedure when handling the XENONnT data.

For comparison with available limits, we consider here the region $10 \text{ keV} \leq m_{Z'} \leq 100 \text{ MeV}$ for the flavor-dependent models. We include our derived limits using XENONnT data for each model. In the case of $L_e - L_\mu$ model and $L_e - L_\tau$ model, we present the

available constraints derived from the solar neutrino facility (BOREXINO), nuclear reactors (TEXONO, CONUS, Dresden-II), stopped-pion source (COHERENT), oscillation data, stellar cooling, and BBN. Meanwhile, in the case of $L_\mu - L_\tau$, we present additional bounds derived from analyzing SN1987A and cosmology (H0) results.

Furthermore, we discuss a comparison of the new light mediator from the $U(1)_{L_e-L_\mu}$ and $U(1)_{L_e-L_\tau}$ gauge symmetries with earlier studies. Figure 30 (b) shows the $L_e - L_\mu$ and Figure 31 (b) shows the $L_e - L_\tau$. Our analysis indicates around an order-of-magnitude improvement relative to the limit from COHERENT and a threefold enhancement over the Dresden-II bound. For $m_{Z'} \leq 0.2$ MeV, the TEXONO constraint is fully covered, with improvements reaching a factor of six. Relative to the limit from CONUS, our derived limits remain competitive, exceeding the bound for $m_{Z'} \leq 20$ keV. In addition, for $m_{Z'} \leq 20$ keV, the $L_e - L_\mu$ limit and, for $m_{Z'} \leq 30$ keV, the $L_e - L_\tau$ limit from Run1 yield stronger bounds. Compared to the limits from XENONnT data, constraints from the Run1 are more stringent for $m_{Z'} \lesssim 20$ keV in both models. Although our bounds have not yet reached the oscillation limits, they provide valuable complementarity to bounds obtained from stellar cooling and BBN. For the $L_\mu - L_\tau$, the comparison with previous works can be seen in Figure 32 (b). Our results set tighter bounds than those from COHERENT and BOREXINO, though they remain a few times weaker than the XENONnT constraints within the examined region. In the low-mass regime, the limits from SN1987A and stellar cooling are already saturated. While most limits in this region are accessed, some bounds of the $(g-2)_\mu$ allowed region persist for $m_{Z'} > 4$ MeV. Additionally, the region that can address the Hubble tension H_0 for $\Delta N_{\text{eff}} \sim 0.2 - 0.5$ taken from (Escudero *et al.*, 2019) is also shown.

In general, it is worth stressing that PandaX-4T is anticipated to yield stronger bounds than those from analyzing data of the low-energy solar neutrino experiment, stopped-pion facility, and nuclear power plant, while also serving as a complementary probe to other endeavors such as high-energy particle colliders or astrophysical searches. We emphasize that PandaX-4T, as a DM-DD experiment, benefits from its energy threshold that lies in low recoil, a large amount of exposure, and substantial target mass, allowing it to set more stringent bounds than many earlier results. This is particularly evident when compared with COHERENT, whose threshold of recoil energy is around ~ 5 keV; limits from the PandaX-4T improve upon the limit from COHERENT by up to an order of magnitude. Moreover, limits

from the nuclear reactor surpass those from the stopped-pion facility and are competitive with limits from PandaX-4T data. As an example, TEXONO, CONNIE, CONUS+, and Dresden-II achieve more stringent limits at heavier mediator masses because of their higher recoil energies. Meanwhile limits from CONUS collaboration, which is able to access a sub-keV threshold, provide competitive sensitivity in the low-mass regime. Relative to the limit from BOREXINO, a solar neutrino experiment, limits from PandaX-4T achieve more stringent bounds in the low-mass regime and dominate at larger masses thanks to its ability to probe higher recoil energies. Overall, PandaX-4T constraints are able to match other DM-DD results in the low-mass regime.

5.3.2. Results from Nuclear Recoil Data

Our limits for the universal light scalar mediator from CE ν NS at CDEX-10 are shown in Figure 33. We see that the upper-limit at 90% C.L. is $g_\phi \lesssim 5.68 \times 10^{-6}$ for $m_\phi < 0.5$ MeV. Meanwhile, the upper-limit of the realistic scenario is $g_\phi \lesssim 3.28 \times 10^{-6}$ and for the optimistic

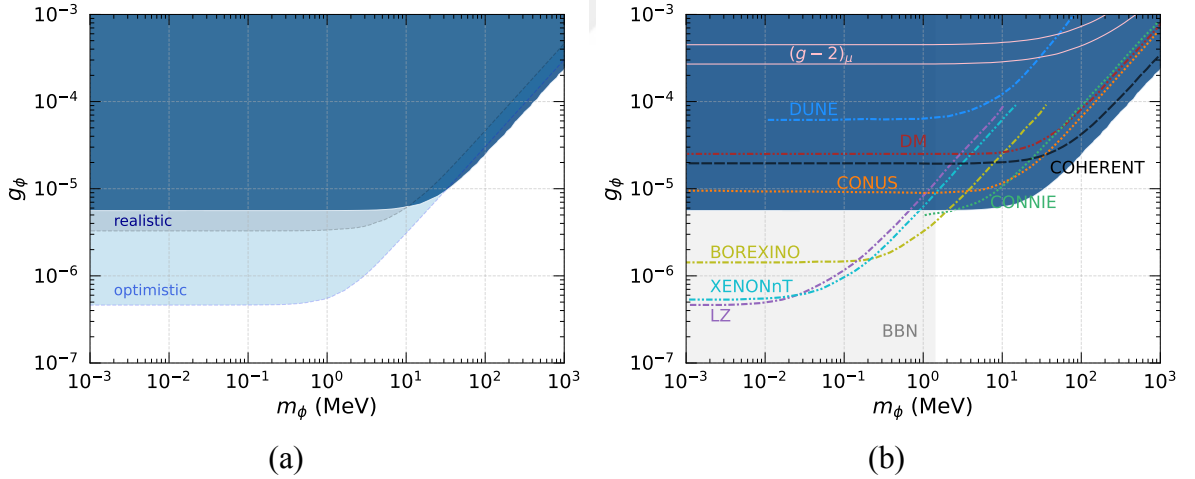


Figure 33. (a) Exclusion limits on universal light scalar mediator from CDEX-10 with CE ν NS and (b) comparison with previous limits.

scenario is $g_\phi \lesssim 4.67 \times 10^{-7}$. These indicate improvements over the current constraint by approximately 42.3% for the realistic scenario and 91.8% for the optimistic scenario. We also overlaid this limit with other limits studied in previous studies. Our results clearly impose stronger constraints than limits obtained from COHERENT, CONNIE, and CONUS. The same behavior is also seen for DUNE and projected dark matter searches. Nevertheless, the

bounds derived from the BOREXINO solar neutrino experiment and LZ and XENONnT are yet to be fully surpassed in the low-mass regions, namely for $m_\phi < 2.0$ MeV, $m_\phi < 0.7$ MeV, and $m_\phi < 0.9$ MeV, respectively. Regarding bounds from the BBN study, our limits reach it when $m_\phi \leq 1.41$ MeV. Notice further that the allowed region for $(g - 2)_\mu$ (2σ) is entirely covered by the derived results. Looking ahead, our projected scenarios suggest that all previously established limits concerning the universal scalar mediator model could be accessed in further advancement to come.

We next show in Figure 34 the derived limits of the universal light vector mediator from analyzing neutrino-nucleus interaction at CDEX-10. We observe from Figure 34 (a) at 90% C.L. that the upper-limit is approximately $g_{Z'} \lesssim 2.28 \times 10^{-5}$ for $m_{Z'} < 1.3$ MeV. For the projected scenarios, the realistic one gives $g_{Z'} \lesssim 5.57 \times 10^{-6}$ and the optimistic one gives $g_{Z'} \lesssim 1.49 \times 10^{-6}$. This yields an improvement of about 75.6% for the realistic scenario and 93.5% for the optimistic scenario. We overlay our results with other existing bounds on the universal vector light mediator in Figure 34 (b). Our analysis shows stronger constraints

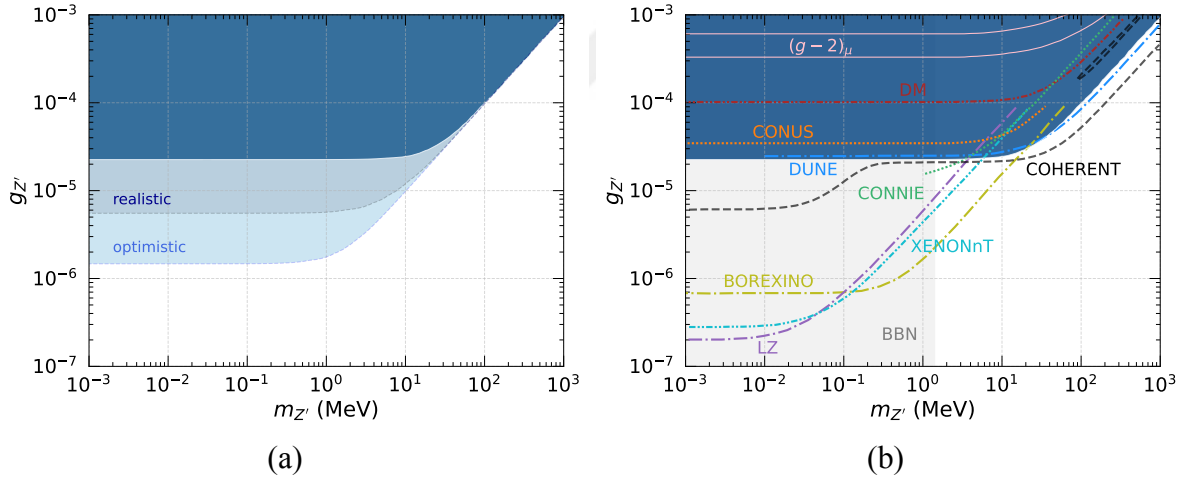


Figure 34. (a) Exclusion limits on universal light vector mediator from CDEX-10 with CEνNS and (b) comparison with previous limits.

than limits derived from the CONUS experiment and projected dark matter searches. The CONNIE limit is accessed as $m_{Z'} < 3.4$ MeV. Meanwhile, the DUNE limit is mostly covered when $m_{Z'} < 40$ MeV. The limit from LZ is surpassed as $m_{Z'} > 3.8$ MeV, XENONnT as $m_{Z'} > 5.8$, and BOREXINO as $m_{Z'} > 20$ MeV. The bound from $(g - 2)_\mu$ preferred region is fully covered. On the other hand, the combined CsI and Ar constraints from COHERENT remain beyond reach. Looking ahead, projected scenarios are expected to access the low-

mass region of COHERENT, though they would still fall short of fully covering the exclusion ranges of BOREXINO, LZ, and XENONnT.

Next, we present our limit for the tensor mediator for the $\text{CE}\nu\text{NS}$ process at CDEX-10 in Figure 35. We find in Figure 35 (a) that the upper-limit at 90% C.L. is $g_T \lesssim 3.35 \times 10^{-5}$ for $m_T \leq 0.8$ MeV. Meanwhile, limits of the realistic scenario yield an improvement of around 64.2%, and the optimistic scenario is about 92.3% relative to the limit obtained from the current data. These correspond to $g_T \lesssim 1.20 \times 10^{-5}$ for the realistic scenario and $g_T \lesssim 2.59 \times 10^{-6}$ for the optimistic scenario. We then overlay available limits from other works of this model in Figure 35 (b). The projected dark matter limit is fully covered by our results.

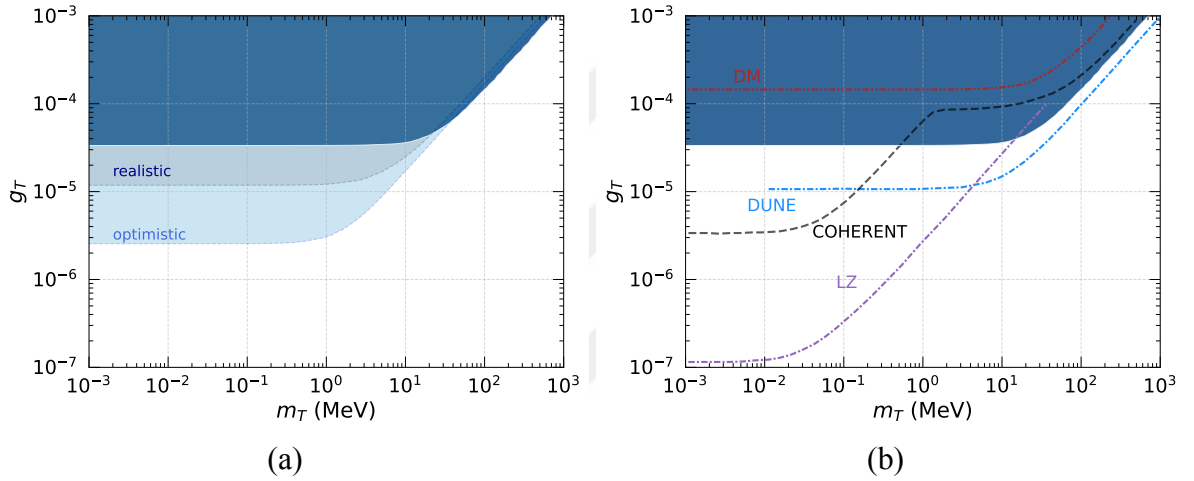


Figure 35. (a) Exclusion limits on universal light tensor mediator from CDEX-10 with $\text{CE}\nu\text{NS}$ and (b) comparison with previous limits.

In addition, we obtain more competitive bounds than COHERENT for $m_T > 0.5$ MeV. Moreover, the limit derived from LZ is surpassed when $m_T > 14$ MeV, while the limit from DUNE remains uncovered. Under projected scenarios, the COHERENT limit can be fully reached, with partial coverage of the LZ as well as DUNE exclusions in the regime with low mass.

We now provide the derived limits for the $U(1)_{B-L}$ light mediator model. The 90% C.L. excluded region is shown in Figure 36 (a). For $m_{Z'} < 0.1$ MeV, the upper bound on the coupling reaches $g_{B-L} \simeq 3.93 \times 10^{-5}$. As anticipated from the $B-L$ model, the limit is relatively less stringent than that of the universal light vector mediator, owing to the different charge of the $U(1)'$. Nevertheless, the CDEX-10 data provide stronger bounds than previous results in certain mass ranges. The projected realistic scenario yields further improvements

of about 75.4%, with $g_{B-L} \lesssim 9.66 \times 10^{-6}$, and the optimistic scenario yields 93.4%, $g_{B-L} \lesssim 2.58 \times 10^{-6}$, over the current bound. We then overlaid our results with available limits in

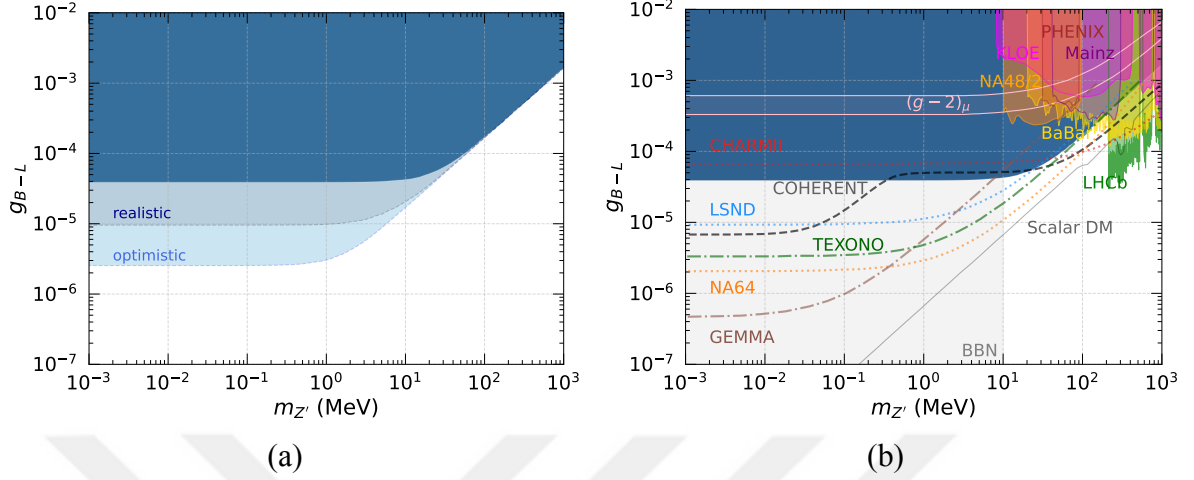


Figure 36. (a) Exclusion limits on $U(1)_{B-L}$ mediator from CDEX-10 with $\text{CE}\nu\text{NS}$ and (b) comparison with previous limits.

Figure 36 (b). These limits are obtained from facility to observe $\text{CE}\nu\text{NS}$ at COHERENT (CsI+LAr) (De Romeri *et al.*, 2023), neutrino-electron scattering experiments such as LSND (Auerbach *et al.*, 2001), CHARMII (Vilain *et al.*, 1994), and NA64 (Andreev *et al.*, 2022), nuclear reactor experiments at TEXONO (Deniz *et al.*, 2010) and GEMMA (Beda *et al.*, 2010), as well as collider experiments (BaBar (Lees *et al.*, 2014), LHCb (Aaij *et al.*, 2020), KLOE (Abelev *et al.*, 2013), Mainz (Merkel *et al.*, 2011), PHENIX (Adare *et al.*, 2015)), and rare-meson decay experiment (NA48/2 (Batley *et al.*, 2015)). We further show the allowed region of $(g-2)_\mu$, which can describe the magnetic moment anomaly of muon (Aguilarrrd *et al.*, 2023) as well as scalar thermal DM (Berlin *et al.*, 2019). With the present data, more stringent limits are obtained than COHERENT in the range $0.33\text{MeV} \leq m_{Z'} \leq 23.33\text{MeV}$. Bound from the CHARM-II is covered for $m_{Z'} \leq 42.4\text{MeV}$, and then the GEMMA constraint is reached for $m_{Z'} \geq 7.68\text{MeV}$. Our results also slightly improve upon the LSND limit as $m_{Z'} \geq 24.91\text{MeV}$. The TEXONO exclusion limit is attained at $m_{Z'} \geq 152.07\text{MeV}$, whereas the NA64 bound remains beyond reach. In addition, the Z' can mediate feeble interactions between scalar dark matter and SM particles (Berlin *et al.*, 2019). Assuming a mass relation $m_{Z'} = 3m_\chi$ between the mediator and scalar DM χ (Andreev *et al.*, 2022), the corresponding bound has not yet been probed by our analysis. The constraint from BBN is accessed for $m_{Z'} < 10$ MeV. Finally, concerning collider searches, the CDEX-10 data are found to cover

parameter space within $8.21 \text{ MeV} \lesssim m_{Z'} \lesssim 1.00 \text{ GeV}$.

Next, we turn to the flavor-dependent light mediator models in the $\text{CE}\nu\text{NS}$ process at CDEX-10, specifically the $B - 3L_e$, $B - 3L_\mu$, and $B - 3L_\tau$ scenarios. For comparison, we also include current limits reported in earlier studies, derived from a variety of experiments: π -DAR (COHERENT (Corona *et al.*, 2022)), nuclear reactors (TEXONO (Heck *et al.*, 2019)), colliders (BaBar (Lees *et al.*, 2014), LHCb Aaij *et al.* (2020)), neutrino trident production (CCFR (Mishra *et al.*, 1991)), and neutron–lead scattering (Barbierri & Ericson, 1975). Constraints from global oscillation analyses for each model are likewise considered (Coloma *et al.*, 2021), along with $(g - 2)_\mu$ and $(g - 2)_e$ (Corona *et al.*, 2022) that can explain the magnetic moment anomalies for each particle. For the $B - 3L_\tau$ case, we additionally include projected bounds from pion and kaon decay data (Heck *et al.*, 2019). Overall, our results set stronger limits than many of these existing constraints, though not yet surpassing all of them.

We show in Figure 37 the excluded region for the $B - 3L_e$ light mediator model, along with its projected sensitivities. From Figure 37 (a), we find that the current data reach $g_{B-3L_e} \lesssim 3.06 \times 10^{-5}$ for $m_{Z'} < 1.6 \text{ MeV}$, while the projected realistic and optimistic scenarios improve this limit by about 75.4% ($g_{B-3L_e} \lesssim 7.50 \times 10^{-6}$) and 93.4% ($g_{B-3L_e} \lesssim 2.00 \times 10^{-6}$). For the comparison, we see in Figure 37 (b) the COHERENT and oscillation bounds are fully excluded, while the $(g - 2)_e$ constraint is accessed for $m_{Z'} < 2.07 \text{ MeV}$. In addition, the BaBar limit is partially covered, whereas the TEXONO bound remains beyond the reach of the current data.

Moving on to the $B - 3L_\mu$ light mediator model, we present the excluded region in Figure 38. We see in Figure 38 (a) that the upper-limit is $g_{B-3L_\mu} \lesssim 3.07 \times 10^{-5}$ when $m_{Z'} < 0.1 \text{ MeV}$, while the realistic scenario gives upper-limit of $g_{B-3L_\mu} \lesssim 7.55 \times 10^{-6}$ and the optimistic scenario gives $g_{B-3L_\mu} \lesssim 2.01 \times 10^{-6}$, correspond to improvements of about 75.4% and 93.4% from the current data, respectively. In Figure 38 (b), we find from the comparison with other constraints that the $(g - 2)_\mu$ and CCFR bounds are largely excluded, while those from COHERENT and oscillation data remain beyond reach. The LHCb constraint is only partially covered in the region of high-mass values.

Next, we provide the exclusion limit in Figure 39 for the $B - 3L_\tau$ light mediator model. From Figure 39 (a), we observe that the current CDEX-10 data set a bound of

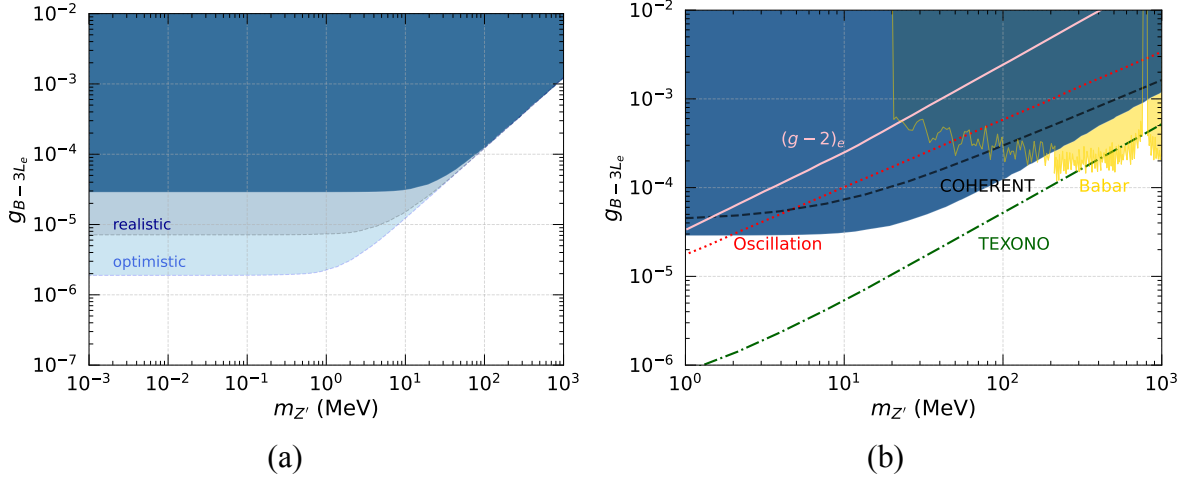


Figure 37. (a) Exclusion limits on $U(1)_{B-3L_e}$ mediator from CDEX-10 with $CE\nu NS$ and (b) comparison with previous limits.

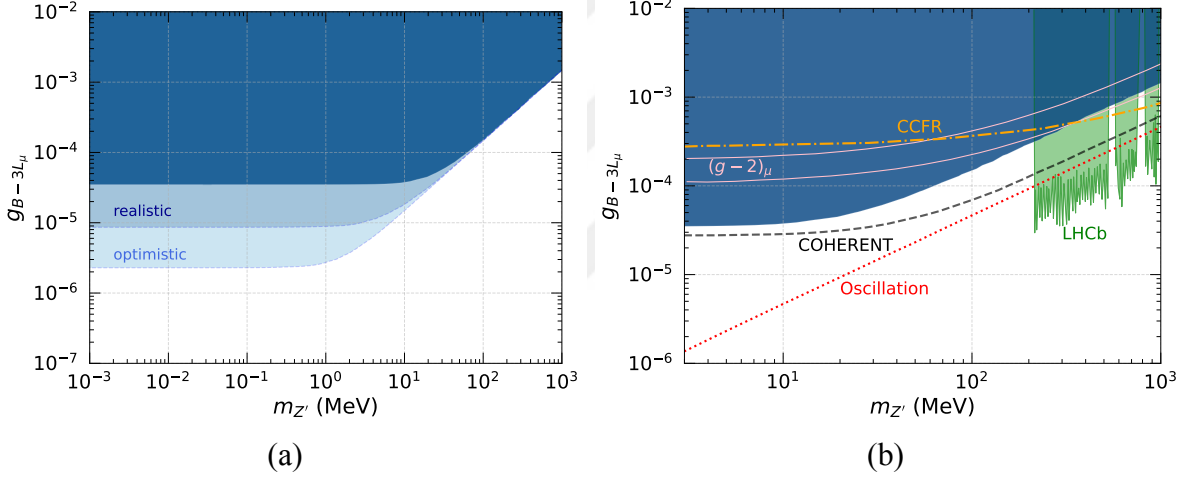


Figure 38. (a) Exclusion limit on $U(1)_{B-3L_\mu}$ mediator from CDEX-10 with $CE\nu NS$ and (b) comparison with previous limits.

$g_{B-3L_\tau} \lesssim 2.84 \times 10^{-5}$ when $m_{Z'} < 0.9$ MeV, while The projected realistic scenario improve this constraint by about 75.1%, and the optimistic scenario by 93.4%. These correspond to $g_{B-3L_\tau} \lesssim 7.07 \times 10^{-6}$ for the realistic and $g_{B-3L_\tau} \lesssim 1.86 \times 10^{-6}$ for the optimistic scenario. When compared with existing limits, we observe in Figure 39 (b) that our results indicate that the derived limits can account for bounds from kaon decay, pion decay, and neutron-lead scattering, though the oscillation constraint remains out of reach.

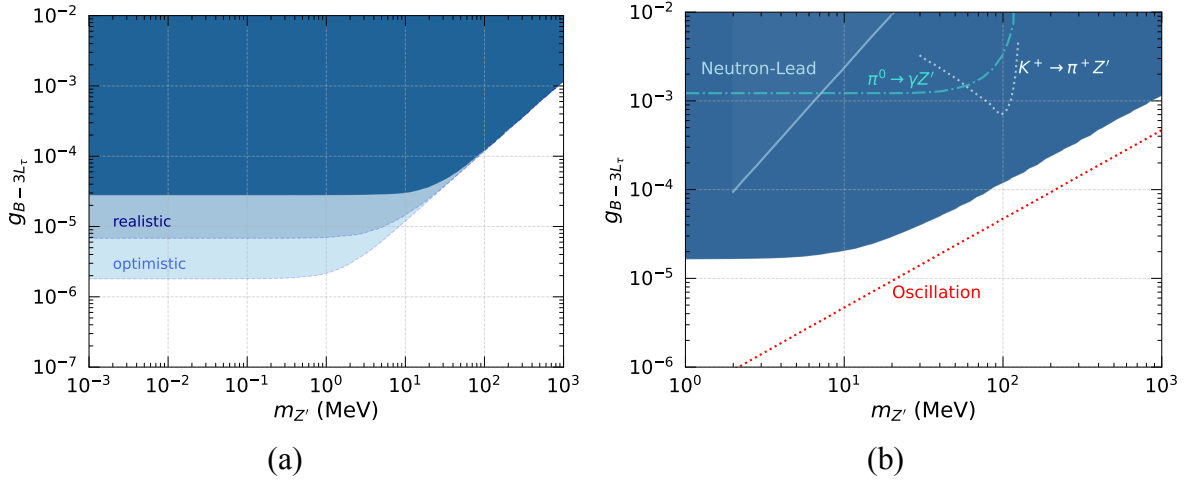


Figure 39. (a) Exclusion limits on $U(1)_{B-3L_\tau}$ mediator from CDEX-10 with $CE\nu NS$ and (b) comparison with previous limits.

5.4. Sterile Neutrino Dipole Portal

In this section, we present our results on the sterile neutrino dipole portal. We first provide limits for the flavor-dependent scenarios, followed by the flavor-independent ones. The results are presented initially for electron recoils, and subsequently for nuclear recoils. We utilize experimental data from CDEX-10, PandaX-4T, and XENONnT to derive the obtained limits. Particular to the CDEX-10, we also provide two projected scenarios regarding future development of the experiment within the $CE\nu NS$ framework.

5.4.1. Results from Electron Recoil Data

We discuss the 2 d.o.f. results at 90% C.L. for the electron-flavor case in Figure 40 (a). For $m_4 \lesssim 0.03$ MeV, the derived upper limits are $\mu_{\nu_{e4}} \lesssim 4.52 \times 10^{-11} \mu_B$ for PandaX-4T Run0, $\mu_{\nu_{e4}} \lesssim 1.91 \times 10^{-11} \mu_B$ for PandaX-4T Run1, and $\mu_{\nu_{e4}} \lesssim 1.86 \times 10^{-11} \mu_B$ for XENONnT in the low-mass region. We observe a substantial improvement from PandaX-4T Run0 to Run1, in which the upper limit is more sensitive by about a factor of two. This enhancement is primarily attributed to the increased exposure and improved background suppression in Run1. Moreover, XENONnT delivers the most stringent constraint in this channel, slightly outperforming PandaX-4T Run1, which reflects the advantage of its larger exposure and lower background levels in the relevant low-energy region. Figure 40 (b) presents

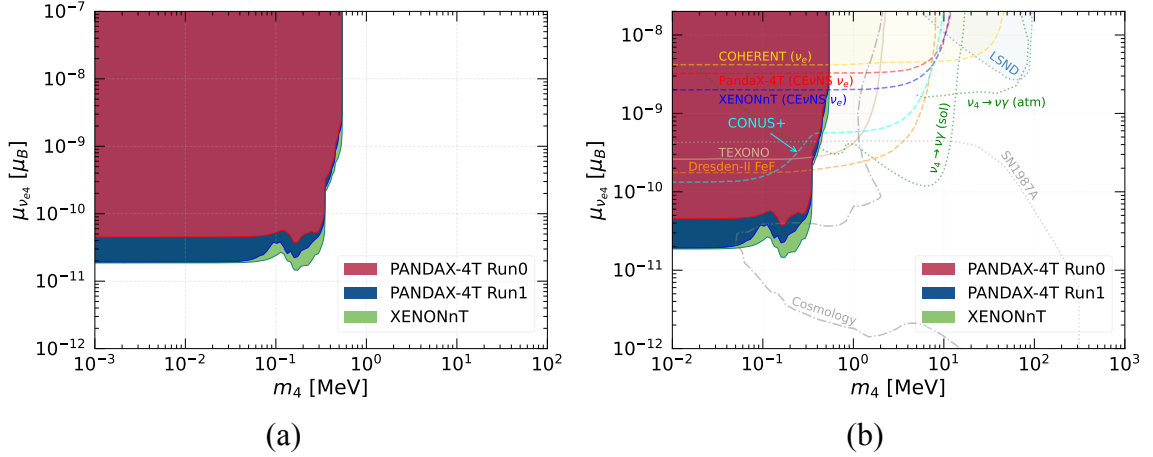


Figure 40. (a) Exclusion limit on μ_{ν_e4} from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.

a comparison of our results with previous experimental limits, highlighting the significant improvements achieved by recent datasets—particularly those from PandaX-4T Run 1 and XENONnT. Notably, our derived limits are more stringent than prior constraints in the sub-MeV mass range. Our results are more stringent, approximately two orders of magnitude, than the limits obtained from CE ν NS processes, such as those reported by COHERENT and from solar neutrino-induced backgrounds in PandaX-4T and XENONnT. In comparison to constraints from nuclear reactor experiments of CONUS+, TEXONO, and Dresden-II FeF, the discrepancy is reduced to about one order of magnitude. However, at mass scales $m_4 \gtrsim 1$ MeV, our results do not surpass the existing limits set by LSND and $\nu_4 \rightarrow \nu\gamma$ decay searches. Additionally, our constraints partially overlap with the parameter space excluded by SN1987A and cosmological observations.

For the muon-flavor case, we present the 2 d.o.f. results at 90% C.L. in Figure 41 (a). The derived upper limits in the region of $m_4 \lesssim 0.03$ MeV are $\mu_{\nu_{\mu4}} \lesssim 7.41 \times 10^{-11} \mu_B$ for PandaX-4T Run0, $\mu_{\nu_{\mu4}} \lesssim 3.15 \times 10^{-11} \mu_B$ for PandaX-4T Run1, and $\mu_{\nu_{\mu4}} \lesssim 3.08 \times 10^{-11} \mu_B$ for XENONnT. We observe a clear improvement in sensitivity from Run0 to Run1, with the upper limit reduced by more than a factor of two. Additionally, XENONnT provides a slightly stronger constraint than PandaX-4T Run1, confirming its competitive performance in this channel. A comparison with previous limits is shown in Figure 41 (b). Apart from the already mentioned limits, here we also superimpose additional limits from the solar neutrino experiment and collider facilities. The PandaX-4T Run0 limit is about as sensitive as the

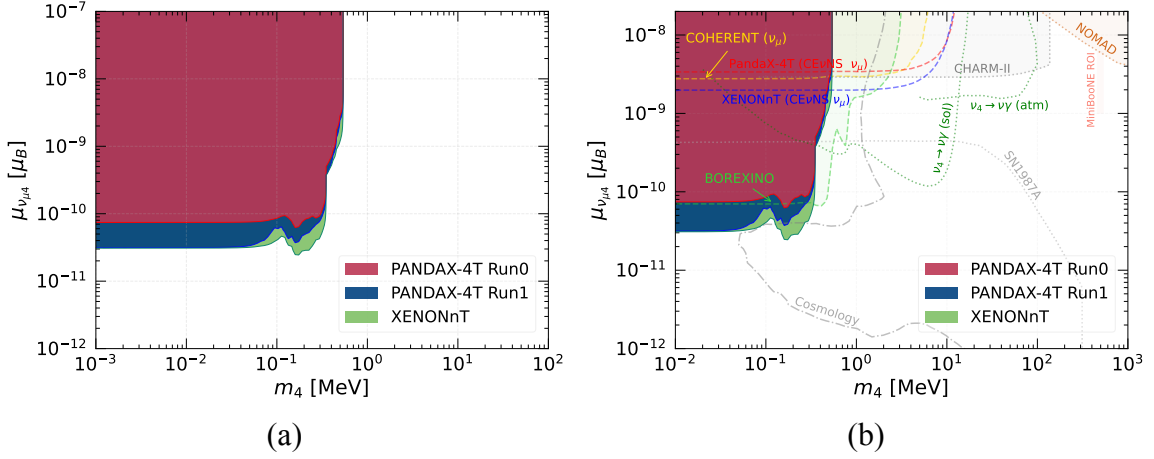


Figure 41. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.

BOREXINO limit, while the PandaX-4T Run1 and XENONnT limits provide a few times stringent bounds than result from this solar neutrino experiment. Limits from CHARM-II, MiniBooNE, and NOMAD dominate at the high-mass region. Regarding the already mentioned limits, similar behavior to the previous case.

For the tau-flavor case, the 2 d.o.f. results at 90% C.L. are shown in Figure 42 (a). The obtained upper limits are $\mu_{\nu_{\tau 4}} \lesssim 6.43 \times 10^{-11} \mu_B$ for PandaX-4T Run0, $\mu_{\nu_{\tau 4}} \lesssim 2.71 \times 10^{-11} \mu_B$ for PandaX-4T Run1, and $\mu_{\nu_{\tau 4}} \lesssim 2.67 \times 10^{-11} \mu_B$ for XENONnT, all in the region of $m_4 \lesssim 0.03 \text{ MeV}$. Similar to the muon-flavor case, we observe a substantial improvement between Run0 and Run1, while XENONnT again delivers the most stringent constraint. This result emphasizes the increasing sensitivity of recent direct detection experiments to the tau-flavor sector, which traditionally suffers from weaker constraints due to the absence of direct tau neutrino sources at these energies. The comparison with existing bounds is provided in Figure 42 (b). In the currently considered parameter space, there are but a few available limits. The induced solar neutrino CE ν NS processes of PandaX-4T and XENONnT are surpassed for $m_4 \lesssim 0.4 \text{ MeV}$, while they dominate the high-mass region up to 10 MeV . Moreover, the sterile-neutrino decay and IceCube/DeepCore bounds occupy the much larger mass region up to the sub-GeV scale.

Lastly, we report the 2 d.o.f. results at 90% C.L. for the flavor-independent case in Figure 43 (a). The corresponding upper limits for $m_4 \lesssim 0.03 \text{ MeV}$ are $\mu_{\nu_{\ell 4}} \lesssim 3.31 \times 10^{-11} \mu_B$ for PandaX-4T Run0, $\mu_{\nu_{\ell 4}} \lesssim 1.40 \times 10^{-11} \mu_B$ for PandaX-4T Run1, and $\mu_{\nu_{\ell 4}} \lesssim 1.38 \times 10^{-11} \mu_B$

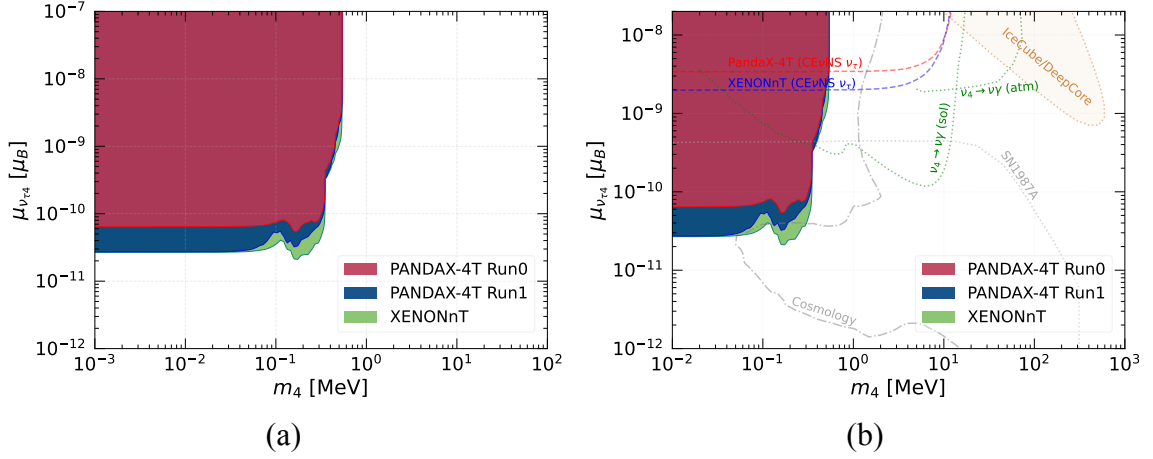


Figure 42. (a) Exclusion limit on $\mu_{\nu\tau 4}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.

for XENONnT. As with the flavor-dependent cases, we observe significant improvements in the limits from Run0 to Run1, with XENONnT providing the tightest constraint overall. A

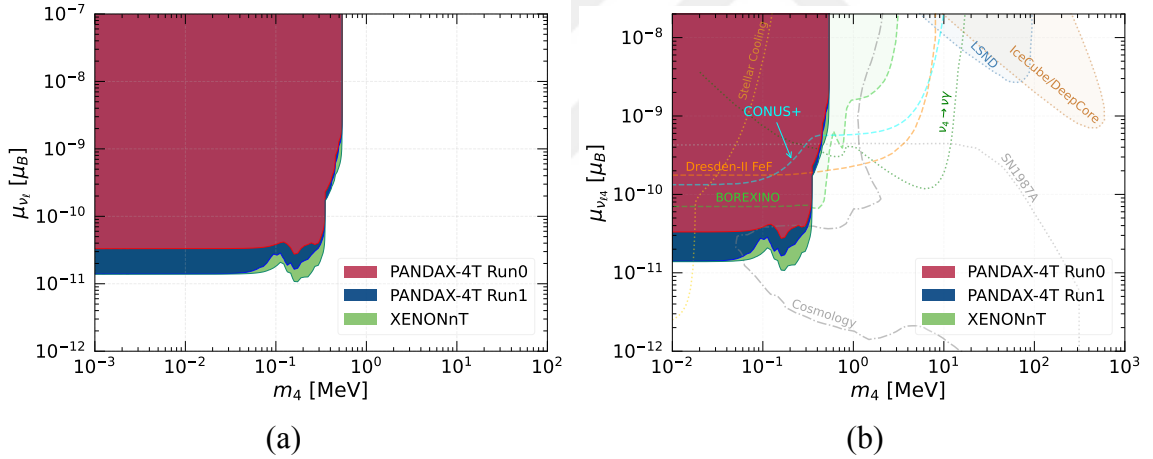


Figure 43. (a) Exclusion limit on $\mu_{\nu\tau 4}$ from PandaX-4T and XENONnT and (b) comparison of the projected results with existing limits.

detailed comparison with earlier results of flavor-dependent cases is presented in Figure 43 (b). In general, similar behavior from previous results is found. Supplementary limit from stellar cooling (Brdar *et al.*, 2020) process is added, indicating that the current data dominate the sub-MeV mass scale and are about one order of magnitude away from probing this region. Overall, the recent low-energy electronic recoil data from DM-DD facilities could provide limit enhancements on the transition magnetic moment compared to other facilities. This comparison shows the progressive tightening of bounds achieved by the latest datasets.

5.4.2. Results from Nuclear Recoil Data

In this section, we present the exclusion limits on the sterile neutrino dipole portal, derived from the $\text{CE}\nu\text{NS}$ process using CDEX-10 nuclear recoil data. We show the electron-flavor limits in Figure 44 (a). It reaches $2.26 \times 10^{-8} \mu_B$ in the region of $m_4 < 0.1$ MeV. Limits for the next-generation reaches $\mu_{\nu_{e4}} \lesssim 1.50 \times 10^{-9} \mu_B$, future 1 $\mu_{\nu_{e4}} \lesssim 1.04 \times 10^{-9} \mu_B$, and future 2 $\mu_{\nu_{e4}} \lesssim 5.08 \times 10^{-10} \mu_B$. We observe that the future 1 configuration provides about 30.5% and future 2 about 66.2%. These are more stringent than the next-generation scenario. As we compare them with the current limit, we see the next-generation provide 15.1, the future 1 provide 21.7, and future 2 provide 44.5 times improvement. In Figure 44 (b), we provide the comparison of the projected sensitivities with the available constraints relevant to $\mu_{\nu_{e4}}$. We infer that the scenarios provide more stringent bounds than limits from COHERENT

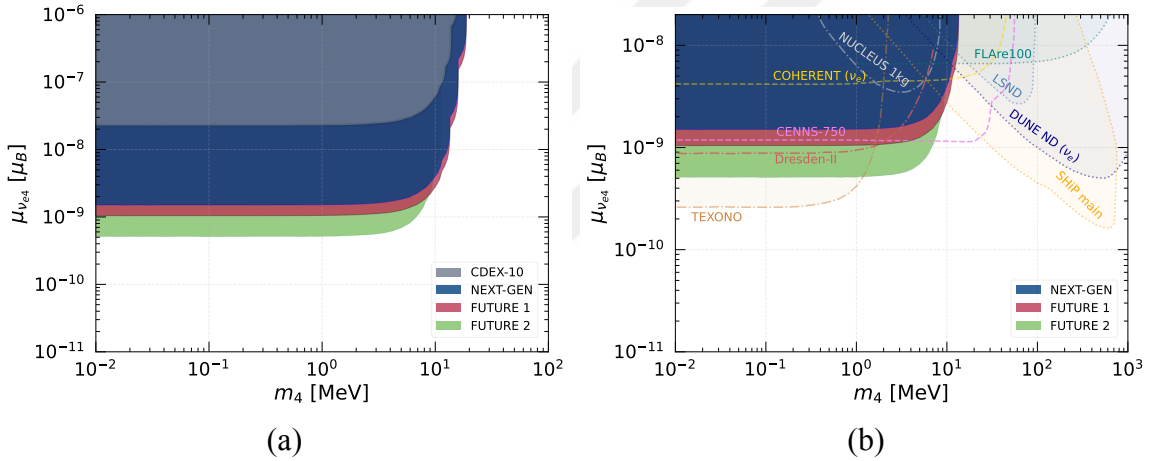


Figure 44. (a) Exclusion limit on $\mu_{\nu_{e4}}$ from CDEX-10 with $\text{CE}\nu\text{NS}$ and the projected future results and (b) comparison of the projected results with existing limits.

(CsI+LAr) in the region with low-mass scale. The improvement arises from the threshold, which is located in the lower recoil energy. It allows for more efficient probing of the $1/T_{nr}$ dependence compared to neutrinos originating from the stopped-pion source. For reactor-based constraints, the next-generation scenario exceeds the Dresden-II limit as $m_4 \gtrsim 3$ MeV, while it is fully covered by the future 2 scenario. Although the projected sensitivities are weaker than TEXONO at masses below 1.2 MeV, at higher mass scales, they offer stronger bounds. As expected, neutrino-electron scattering at reactor facilities dominates in the low-mass regime relative to $\text{CE}\nu\text{NS}$ results obtained from solar neutrino facility, because of the threshold that lies in lower energy, triggered by the incoming neutrino. Looking ahead, the

limit from NUCLEUS 1 kg is fully covered by the projected scenarios and reaches the limits from FLAre100 and CENNS-750 when $m_4 < 10$ MeV. Additionally, a significant section of the previously unreachable region of LSND and DUNE ND and SHiP (main) is accessed in the region with a MeV mass scale, reflecting the high energies of produced neutrinos in accelerators, which are primarily more stringent at higher masses.

We next discuss the muon-neutrino flavor. We present the derived limits of the sterile neutrino dipole portal from the CDEX-10 data, as well as the next-generation and two future scenarios in Figure 45 (a). The CDEX-10 limit for $\mu_{\nu_{\mu 4}}$ is about $2.41 \times 10^{-8} \mu_B$ for $m_4 \lesssim 0.1$ MeV. At the same time, the next-generation limit is $\mu_{\nu_{\mu 4}} \lesssim 1.59 \times 10^{-9} \mu_B$. Moreover, the future 1 limit is $\mu_{\nu_{\mu 4}} \lesssim 1.11 \times 10^{-9} \mu_B$ and future 2 limit is $\mu_{\nu_{\mu 4}} \lesssim 5.64 \times 10^{-10} \mu_B$. This corresponds to approximately 30.1% and 64.4% more stringent limits for the future 1 scenario and future 2 scenario compared to the next-generation scenario. Relative to the CDEX-10 data, the next-generation scenario provides improvements by a factor of 15.2, future 1 by 21.7, and future 2 by 42.7. Figure 45 (b) compares these results with previous limits. The

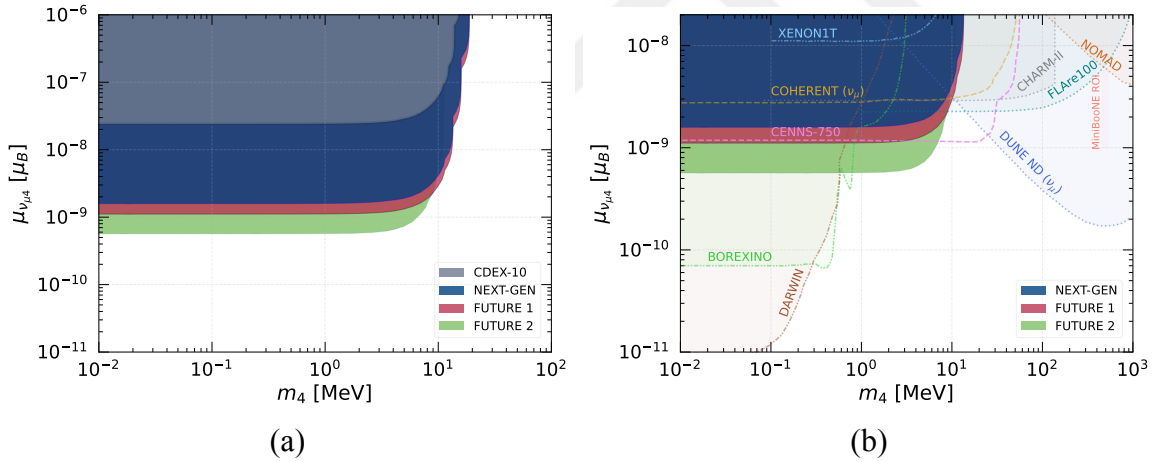


Figure 45. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from CDEX-10 with CEvNS and the projected future results and (b) comparison of the projected results with existing limits.

sensitivities from projected scenarios surpass the limit from XENON1T and improve the limit from COHERENT (CsI+LAr) and CHARM-II by about two orders of magnitude. They also strengthen limits from BOREXINO for $m_4 \gtrsim 0.6$ MeV, though in the very low-mass regime, neutrino–electron scattering still provides stronger constraints because of the lower energy threshold. For future facilities, our scenarios probe limits obtained from FLAre100, CENNS-750, and DUNE ND in the low-mass regions. While the projected limits do not yet

access the expected limit from DARWIN for $m_4 \lesssim 0.6$ MeV, they improve constraints in the higher-mass region. In contrast, NOMAD and the MiniBooNE ROI dominate at high masses, reflecting the higher neutrino energies of accelerator-based experiments.

We provide our limits of the sterile neutrino dipole portal for the tau-neutrino flavor in Figure 46 (a). It includes the limit from the CDEX-10 data and also limits from the next-generation, future 1, and future 2 scenarios. We find that the CDEX-10 data set an upper bound of $\mu_{\nu_{\tau 4}} \simeq 2.08 \times 10^{-8} \mu_B$ for $m_4 \lesssim 0.1$ MeV. Moreover, limit from the next-generation is $\mu_{\nu_{\tau 4}} \lesssim 1.39 \times 10^{-9} \mu_B$, future 1 $\mu_{\nu_{\tau 4}} \lesssim 9.50 \times 10^{-10} \mu_B$, and future 2 $\mu_{\nu_{\tau 4}} \lesssim 4.86 \times 10^{-10} \mu_B$. The future 1 and future 2 scenarios offer roughly 31.6% and 65.0% more sensitive limits than the scenario of next-generation, while all three scenarios improve over the current data by factors of 15.0, 21.9, and 42.8, respectively. We present in Figure 46 (b) a comparison of these

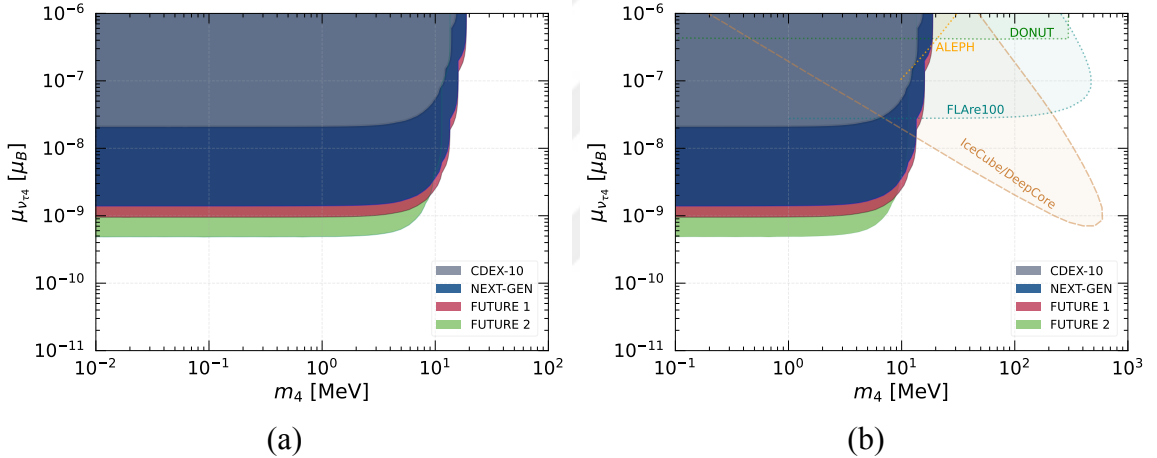


Figure 46. (a) Exclusion limit on $\mu_{\nu_{\mu 4}}$ from CDEX-10 with $\text{CE}\nu\text{NS}$ and the projected future results and (b) comparison of the projected results with existing limits.

results with existing constraints derived from facilities like ALEPH and DONUT, as well as FLArE100 and IceCube/DC. It is indicated that the data from CDEX-10 and the projected scenarios provide stronger bounds compared to these previous limits, especially when $m_4 < 10$ MeV. Specifically, the scenarios improve upon DONUT and ALEPH constraints for $m_4 \lesssim 19$ MeV and FLArE-100 for $m_4 \lesssim 17$ MeV. This enhancement is because of the lower-energy spectrum from solar neutrinos compared to accelerator-produced neutrinos. Moreover, the derived limits cover a larger section of the parameter space relative to the IceCube/DC limits for $m_4 \lesssim 15$ MeV, highlighting that solar neutrinos can probe sterile-neutrino mass with lower energy than the atmospheric neutrinos.

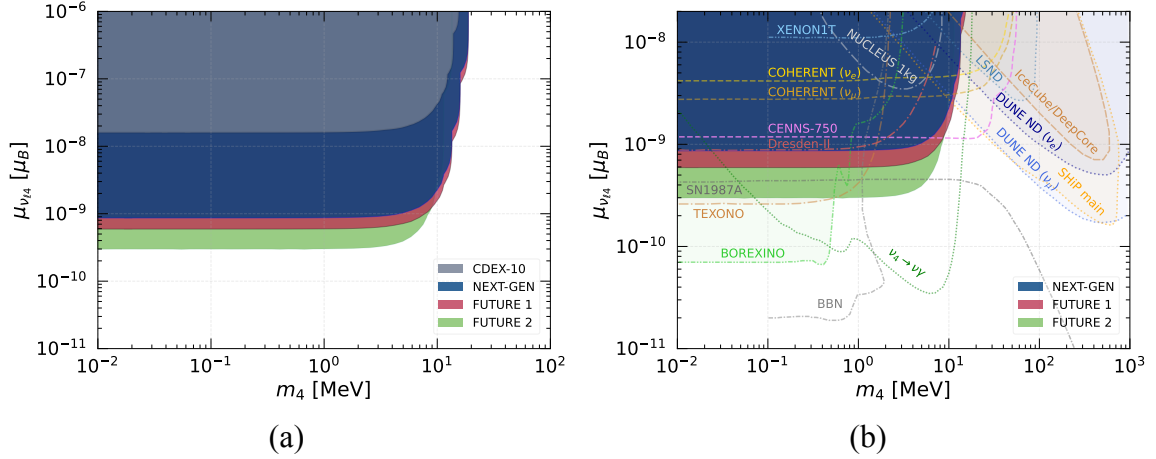


Figure 47. (a) Exclusion limit on $\mu_{\nu_{\ell 4}}$ from CDEX-10 with CE ν NS and the projected future results and (b) comparison of the projected results with existing limits.

Finally, we discuss the limits of the flavor-independent sterile neutrino dipole portal. We present the result in Figure 47 (a). Using the CDEX-10 data, the upper limit is $\mu_{\nu_{\ell 4}} \simeq 1.59 \times 10^{-8} \mu_B$ for $m_4 \lesssim 0.1$ MeV. Moreover, the next-generation limit is $\mu_{\nu_{\ell 4}} \lesssim 8.63 \times 10^{-10} \mu_B$, future 1 is $\mu_{\nu_{\ell 4}} \lesssim 5.91 \times 10^{-10} \mu_B$, and future 2 is $\mu_{\nu_{\ell 4}} \lesssim 2.97 \times 10^{-10} \mu_B$. It can be seen that the scenarios of future 1 and future 2 achieve roughly 31.5% and 65.5% stronger limits when compared to the next-generation case, while improvements over the limit from CDEX-10 reach factors of 18.4, 26.9, and 53.3, respectively. Figure 47 (b) compares these projected sensitivities with existing limits available from previous works. The projected scenarios provide stronger bounds than several previous limits, fully covering the regions constrained by XENON1T, NUCLEUS 1 kg, and Dresden-II. They also probe previously unexplored parameter space from COHERENT (ν_e , ν_μ) and LSND, as well as from CENNS-750, SHiP main, DUNE ND (ν_e , ν_μ), and IceCube/DeepCore. Furthermore, we can see improvements are as $m_4 \gtrsim 0.5$ MeV, where BOREXINO and TEXONO neutrino-electron scattering limits are less sensitive. For the region $m_4 \gtrsim 1$ MeV, CE ν NS dominates the constraints due to kinematic advantages, whereas E ν ES is more sensitive at lower energy scales. For cosmological bounds, our limits reach the BBN bound when $m_4 \gtrsim 1$ MeV, and only the future 2 limit reaches the SN1987A constraint. In addition, we include the $\nu_4 \rightarrow \nu\gamma$ limit to complement our results. These advancements emphasize the crucial role of lowering nuclear energy thresholds and enhancing exposure in DM-DD experiments.

5.5. Triplet Higgs Model

We discuss in this section our results on the charged Higgs boson parameters from analyzing PandaX-4T and XENONnT recent electron recoil data. We provide limits for $f_{ee}/m_{H^\pm}$, $f_{e\mu}/m_{H^\pm}$, and $f_{e\tau}/m_{H^\pm}$ from the model contribution to the $E\nu$ ES process. For the one-dimensional limits, we hold one parameter and set the others to zero. As for the two-

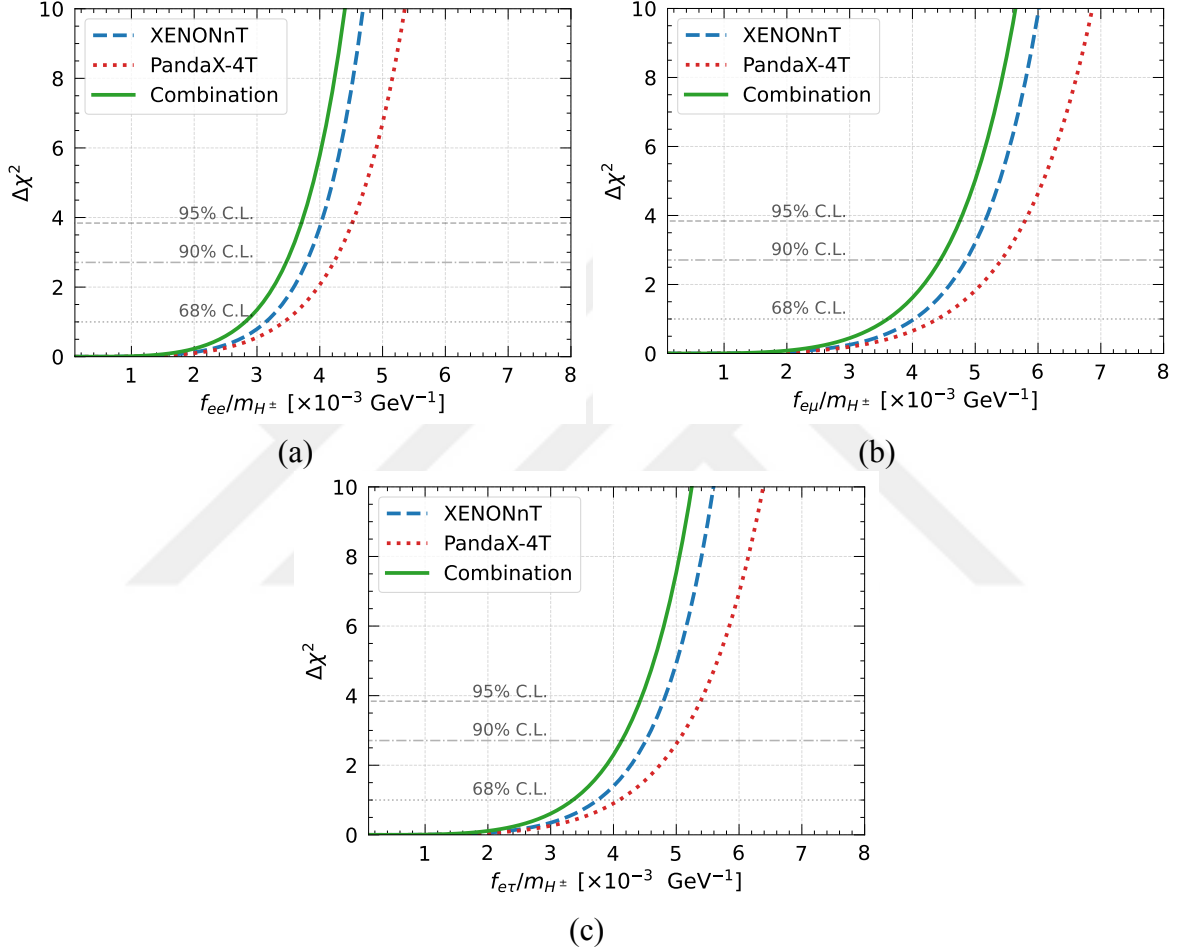


Figure 48. Exclusion limit on the considered parameters from PandaX-4T and XENONnT for (a) $f_{ee}/m_{H^\pm}$, (b) $f_{e\mu}/m_{H^\pm}$, and (c) $f_{e\tau}/m_{H^\pm}$.

dimensional limits, we hold two parameters from two different neutrino flavors and assign the other inactive. Furthermore, we compare our limits with previous bounds available.

We present in Figure 48 the one-dimensional $\Delta\chi^2$ distribution for $f_{ee}/m_{H^\pm}$, $f_{e\mu}/m_{H^\pm}$, and $f_{e\tau}/m_{H^\pm}$. We incorporate the 68% C.L., 90% C.L., and 95% C.L. limits of XENONnT, PandaX-4T, and their combination. For instance, the limits at the 90% C.L. from the data of XENONnT are $f_{ee}/m_{H^\pm} \leq 3.79 \times 10^{-3} \text{ GeV}^{-1}$, $f_{e\mu}/m_{H^\pm} \leq 4.86 \times 10^{-3} \text{ GeV}^{-1}$, and

$f_{e\tau}/m_{H^\pm} \leq 4.52 \times 10^{-3} \text{ GeV}^{-1}$. For the PandaX-4T, the obtained limits are $f_{ee}/m_{H^\pm} \leq 4.22 \times 10^{-3} \text{ GeV}^{-1}$, $f_{e\mu}/m_{H^\pm} \leq 5.42 \times 10^{-3} \text{ GeV}^{-1}$, and $f_{e\tau}/m_{H^\pm} \leq 5.04 \times 10^{-3} \text{ GeV}^{-1}$. As for the two experiment combinations, the limits are $f_{ee}/m_{H^\pm} \leq 3.46 \times 10^{-3} \text{ GeV}^{-1}$, $f_{e\mu}/m_{H^\pm} \leq 4.44 \times 10^{-3} \text{ GeV}^{-1}$, and $f_{e\tau}/m_{H^\pm} \leq 4.13 \times 10^{-3} \text{ GeV}^{-1}$. We can read these limits from these figures. For complementary, we further perform a two-dimensional analysis

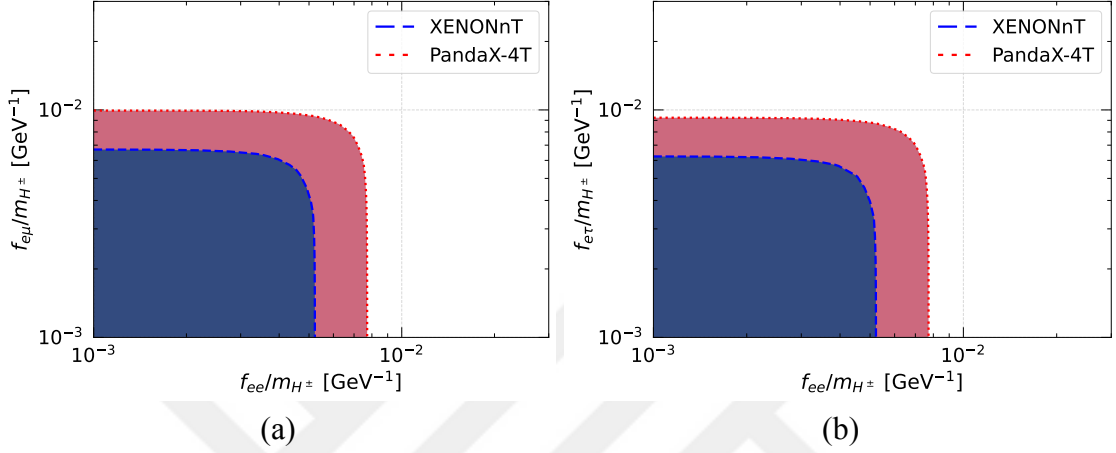


Figure 49. The 90% C.L. allowed regions with 2 d.o.f. in the (a) $f_{e\mu}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ and (b) $f_{e\tau}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ planes.

upon the model parameters. We present our derived limits with 90% C.L. for the allowed regions in Figure 49 for $f_{e\mu}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ and $f_{e\tau}/m_{H^\pm}$ vs $f_{ee}/m_{H^\pm}$ parameter spaces. We find that limits from the data of XENONnT indicate a few times more stringent bounds relative to the limits from PandaX-4T.

Previous studies have investigated bounds on the charged Higgs boson parameters through neutrinos produced in reactor and beam experiments. For $f_{ee}/m_{H^\pm}$, the upper bound from analyzing the data of TEXONO was found to be $f_{ee}/m_{H^\pm} < 2.57 \times 10^{-3} \text{ GeV}^{-1}$ (90% C.L.), more stringent than the other previous result of $f_{ee}/m_{H^\pm} < 2.80 \times 10^{-3} \text{ GeV}^{-1}$ (Ong *et al.*, 2013) concerning the same experiment. As from LSND experiment, the upper bound was found to be $f_{ee}/m_{H^\pm} < 6.48 \times 10^{-3} \text{ GeV}^{-1}$ (90% C.L.). These limits were reported by previous study in (Sevda *et al.*, 2017). Additionally, a previous study using LAMPF (Allen *et al.*, 1993) data at 90% C.L. was obtained $f_{ee}/m_{H^\pm} < 7.20 \times 10^{-3} \text{ GeV}^{-1}$ (Coarasa Perez *et al.*, 1995). Concerning the $f_{e\mu}/m_{H^\pm}$, the limit from anomalous magnetic moment was found to be $f_{e\mu}/m_{H^\pm} < 1.20 \times 10^{-2} \text{ GeV}^{-1}$ (90% C.L.) and from CHARM II was $f_{e\mu}/m_{H^\pm} < 9.90 \times 10^{-4} \text{ GeV}^{-1}$ (90% C.L.) Coarasa Perez *et al.* (1995). Meanwhile for $f_{e\tau}/m_{H^\pm}$, the

limit from the anomalous magnetic moment was found to be $f_{e\tau}/m_{H^\pm} < 0.54 \text{ GeV}^{-1}$ (90% C.L.) (Coarasa Perez *et al.*, 1995).

5.6. Conclusion

In this work, we have demonstrated that $E\nu\text{ES}$ and $CE\nu\text{NS}$, induced by solar neutrinos at direct dark matter detection experiments, can serve as powerful probes for beyond the SM physics. Although these two processes constitute an irreducible background in such experiments, they also provide a unique opportunity to explore new physics within the neutrino sector. We have specifically focused on three BSM scenarios: new light mediator models, active-sterile neutrino transition magnetic moment, and charged Higgs contribution. To constrain the parameters of each model, we analyzed recent low-energy recoil data from the CDEX-10, PandaX-4T, and XENONnT collaborations. Based on our study, we summarize the key conclusions for each scenarios as follows:

5.6.1. New light mediator models

We have derived robust limits using recent CDEX-10, PandaX-4T, and XENONnT data and compared them with available bounds derived from various experimental probes. We have considered the universal light mediator models and mediators from $U(1)'$ models. The derived limits cover some of the previously established limits, while others remain beyond reach. For the universal scalar mediator model and universal vector mediator model, our constraints enhance upon those derived from CONUS, CONNIE, as well as other projected dark matter experiments. For the universal tensor mediator model, even results from the DM-DD projection can be accessed. Furthermore, our results enhance all regions for the scalar mediator of the DUNE limit, partially outperform the high-mass region by the universal vector mediator, and are fully dominated by the universal tensor mediator. For COHERENT, our constraints dominate the limit for the universal scalar mediator, remain below reach for the universal vector mediator, and show partial improvement for the universal tensor mediator. Meanwhile, limits obtained from solar neutrino experiment at BOREXINO and from DM-DD experiments at XENONnT and LZ – at least over a significant portion – remain

unattained across all considered models.

We list our results for light mediator models at $E\nu$ ES in Table 7. These limits can be read from Figure 27, Figure 28, and Figure 29 for the universal model, and from Figure 30, Figure 31, and Figure 32 for the flavor-dependent models. We further provide available limits from analyzing other experiments, such as COHERENT, BOREXINO, and XENONnT data for comparison. We can see that the obtained limits are more stringent as we compare them to the COHERENT limit (neutrino beamline) and BOREXINO limit (solar neutrino), and competitive with the XENONnT limit (DD).

Table 7. Upper limits with 90% C.L. on g_X for light mediator models from PandaX-4T and previous limits from COHERENT ((De Romeri *et al.*, 2024; Coloma *et al.*, 2021; Melas *et al.*, 2023)), BOREXINO (Coloma *et al.*, 2021; Gninenko, 2021), and XENONnT (ShivaSankar *et al.*, 2023).

g_X	PandaX-4T		COHERENT (CsI+LAr)	BOREXINO	XENONnT
	Run0	Run1			
g_ϕ	$\lesssim 1.1 \times 10^{-6}$	$\lesssim 8.3 \times 10^{-7}$	$\lesssim 1.9 \times 10^{-5}$	$\lesssim 1.5 \times 10^{-6}$	$\lesssim 5.3 \times 10^{-7}$
$g_{Z'}$	$\lesssim 5.1 \times 10^{-7}$	$\lesssim 2.4 \times 10^{-7}$	$\lesssim 6.1 \times 10^{-6}$	$\lesssim 6.8 \times 10^{-7}$	$\lesssim 1.9 \times 10^{-7}$
g_T	$\lesssim 2.8 \times 10^{-7}$	$\lesssim 1.4 \times 10^{-7}$	$\lesssim 3.5 \times 10^{-6}$	-	$\lesssim 1.0 \times 10^{-7}$
$g_{L_e-L_\mu}$	$\lesssim 8.2 \times 10^{-7}$	$\lesssim 3.9 \times 10^{-7}$	$\lesssim 6.1 \times 10^{-6}$	$\lesssim 6.2 \times 10^{-7}$	$\lesssim 4.0 \times 10^{-7}$
$g_{L_e-L_\tau}$	$\lesssim 1.2 \times 10^{-7}$	$\lesssim 3.7 \times 10^{-7}$	$\lesssim 7.9 \times 10^{-6}$	$\lesssim 6.4 \times 10^{-7}$	$\lesssim 3.8 \times 10^{-7}$
$g_{L_\mu-L_\tau}$	$\lesssim 6.0 \times 10^{-6}$	$\lesssim 3.3 \times 10^{-6}$	$\lesssim 7.0 \times 10^{-5}$	$\lesssim 7.0 \times 10^{-5}$	$\lesssim 3.2 \times 10^{-6}$

Table 8. Upper limits with 90% C.L. on g_X for light mediator models from CDEX-10 and previous limits from COHERENT (Corona *et al.*, 2022; De Romeri *et al.*, 2023) and XENONnT (ShivaSankar *et al.*, 2023).

g_X	CDEX-10	COHERENT (CsI+LAr)	XENONnT
g_ϕ	$\lesssim 5.68 \times 10^{-6}$	$\lesssim 1.9 \times 10^{-5}$	$\lesssim 5.3 \times 10^{-7}$
$g_{Z'}$	$\lesssim 2.28 \times 10^{-5}$	$\lesssim 6.1 \times 10^{-6}$	$\lesssim 1.9 \times 10^{-7}$
g_T	$\lesssim 3.35 \times 10^{-5}$	$\lesssim 3.5 \times 10^{-6}$	-
g_{B-L}	$\lesssim 3.93 \times 10^{-5}$	$\lesssim 6.74 \times 10^{-6}$	-
g_{B-3L_e}	$\lesssim 3.06 \times 10^{-5}$	$\lesssim 4.54 \times 10^{-7}$	-
g_{B-3L_μ}	$\lesssim 3.07 \times 10^{-5}$	$\lesssim 2.76 \times 10^{-5}$	-
g_{B-3L_τ}	$\lesssim 2.84 \times 10^{-5}$	-	-

We list our results for new light mediator models at $CE\nu$ NS in Table 8. These limits can be directly read from Figure 33, Figure 34, Figure 35, and Figure 36 for the flavor-

independent case, and from Figure 37, Figure 38, and Figure 39 for the flavor-dependent case. We also present available limits from COHERENT and XENONnT for comparative purposes. We can deduce that the results from DM-DD experiments provide a more stringent bound compared to the spallation neutron source. Additionally, we can also see that electron recoil data provides approximately an order of magnitude stringent bound than using nuclear recoil data.

5.6.2. Active-sterile neutrino transition magnetic moment

We have explored both the flavor-dependent case and flavor-independent case of active-sterile neutrino transition magnetic moments, incorporating the neutrino oscillation parameters and flavor conversion during propagation.

We list our results of the active-sterile neutrino transition magnetic moment within the $E\nu$ ES framework in Table 9. These limits can be read directly from Figure 40, Figure 41, and Figure 42 for flavor-dependent cases, as well as Figure 43 for flavor-independent case. It clearly illustrates the sensitivity comparison of our results with available limits from other experiments. We can see that results from PandaX-4T Run1 and XENONnT are currently the most competitive limits, in which XENONnT slightly provides stringent limits. Moving

Table 9. Upper limits with 90% C.L. on the transition magnetic moments from PandaX-4T Run0 and Run1 and previous limits..

$\mu_{\nu\ell 4}$	PandaX-4T		XENONnT
	Run0	Run1	
$\mu_{\nu\ell 4}$	$\lesssim 3.31 \times 10^{-11} \mu_B$	$\lesssim 1.40 \times 10^{-11} \mu_B$	$\lesssim 1.38 \times 10^{-11} \mu_B$
$\mu_{\nu e 4}$	$\lesssim 4.52 \times 10^{-11} \mu_B$	$\lesssim 1.91 \times 10^{-11} \mu_B$	$\lesssim 1.86 \times 10^{-11} \mu_B$
$\mu_{\nu \mu 4}$	$\lesssim 7.41 \times 10^{-11} \mu_B$	$\lesssim 3.15 \times 10^{-11} \mu_B$	$\lesssim 3.08 \times 10^{-11} \mu_B$
$\mu_{\nu \tau 4}$	$\lesssim 6.43 \times 10^{-11} \mu_B$	$\lesssim 2.71 \times 10^{-11} \mu_B$	$\lesssim 2.67 \times 10^{-11} \mu_B$

on to the $CE\nu$ NS framework, we list our results in Table 10. These limits can be read directly from Figure 44, Figure 45, and Figure 46 for flavor-dependent cases, and in Figure 47 for flavor-independent case. Overall, the current data yield competitive constraints compared to other DM-DD experiments, while the projected scenarios offer substantial improvements over several existing limits. We observe that reducing the experimental energy threshold and

Table 10. Upper limits with 90% C.L. on the transition magnetic moments from CDEX-10 and projected limits.

$\mu_{\nu\ell A}$	CDEX-10	Next-Gen	Future 1	Future 2
$\mu_{\nu e A}$	$\lesssim 22.6 \times 10^{-9}$	$\lesssim 1.50 \times 10^{-9}$	$\lesssim 1.04 \times 10^{-9}$	$\lesssim 0.51 \times 10^{-9}$
$\mu_{\nu\mu A}$	$\lesssim 24.1 \times 10^{-9}$	$\lesssim 1.59 \times 10^{-9}$	$\lesssim 1.11 \times 10^{-9}$	$\lesssim 0.56 \times 10^{-9}$
$\mu_{\nu\tau A}$	$\lesssim 20.8 \times 10^{-9}$	$\lesssim 1.39 \times 10^{-9}$	$\lesssim 0.95 \times 10^{-9}$	$\lesssim 0.49 \times 10^{-9}$
$\mu_{\nu\ell A}$	$\lesssim 15.9 \times 10^{-9}$	$\lesssim 0.86 \times 10^{-9}$	$\lesssim 0.59 \times 10^{-9}$	$\lesssim 0.30 \times 10^{-9}$

increasing exposure in the projected setups indicate enhancements of the current data limits by nearly two orders of magnitude.

Sterile neutrinos continue to represent one of the most intriguing extensions to the SM, especially given their potential to resolve the origin of neutrino mass. Our study demonstrates the capabilities of current-generation direct detection experiments, using solar neutrinos via $E\nu$ ES and $CE\nu$ NS, to explore this phenomenological scenario in a complementary and competitive manner. Looking ahead, future improvements in detector exposure, background suppression, and lower recoil energy thresholds are anticipated to further enhance the sensitivity to sterile neutrino dipole portal. Such developments will potentially enable the next generation of experiments to probe transition magnetic moments at levels previously accessible only via astrophysical or cosmological observations, offering valuable new tests of physics beyond the SM.

5.6.3. Charged Higgs boson

Our limits on the charged Higgs boson content of the HTM generally provide more stringent bounds compared to previous studies. These results indicate the minimum sensitivity required to push the limits of the charged Higgs parameter further to reveal direct observation or rule out the proposal. Therefore, the currently developed DM-DD facilities can potentially be used as a testing ground to search for the charged Higgs boson of the HTM.

The limits of the $f_{e\alpha}^2/m_{H^\pm}^2$ is summarized in Table 11. We present the 1 d.o.f. limits with 68% C.L., 90% C.L., and 95% C.L. of each parameter. We provide here limits from analyzing the electron recoil signal from XENONnT and PandaX-4T, as well as combinations of them. The XENONnT data is clearly able to provide more stringent limits than the PandaX-

Table 11. Upper limits with 1 d.o.f. on the charged Higgs parameters obtained from analyzing data of XENONnT, PandaX-4T, and their combination.

$f_{e\alpha}^2/m_{H^\pm}^2$	Experiment	Limits ($\times 10^{-3} \text{ GeV}^{-1}$)		
		68% C.L.	90% C.L.	95% C.L.
$f_{ee}/m_{H^\pm}$	XENONnT	≤ 3.14	≤ 3.79	≤ 4.02
	PandaX-4T	≤ 3.43	≤ 4.22	≤ 4.52
	Combination	≤ 2.81	≤ 3.46	≤ 3.71
$f_{e\mu}/m_{H^\pm}$	XENONnT	≤ 4.03	≤ 4.86	≤ 5.16
	PandaX-4T	≤ 4.40	≤ 5.42	< 5.79
	Combination	≤ 3.61	≤ 4.44	≤ 4.75
$f_{e\tau}/m_{H^\pm}$	XENONnT	≤ 3.75	≤ 4.52	≤ 4.80
	PandaX-4T	≤ 4.09	≤ 5.04	≤ 5.39
	Combination	≤ 3.35	≤ 4.13	≤ 4.42

4T concerning each experiment result, while the combination of both data improves the upper limits a few times.

All in all, $E\nu\text{ES}$ and $\text{CE}\nu\text{NS}$ processes induced by solar neutrinos provide improvements to the existing limits on the considered scenarios. This emphasizes the value of integrating solar neutrino observations with direct detection experiments to gain a deeper concept of how neutrinos interact with ordinary matter. We anticipate that our findings will provide complementary insights for exploring new physics through present and upcoming experiments with solar and other neutrino sources.

6. SUGGESTIONS

The phenomenological studies presented in this work are expected to provide valuable guidance for future efforts in probing beyond the SM physics, particularly in the context of direct dark matter detection experiments, which are rapidly advancing worldwide. Based on our findings, we offer the following suggestions:

- Incorporating radiative corrections to both $E\nu$ ES and $CE\nu$ NS processes to improve the precision of theoretical predictions.
- Extending the analysis to other BSM scenarios within the $E\nu$ ES and $CE\nu$ NS frameworks, such as neutrino electromagnetic properties and non-standard neutrino interactions.
- For $CE\nu$ NS specifically, conducting a more detailed treatment of nuclear response contributions, considering a variety of nuclear targets to improve modeling accuracy.
- Utilizing future low-threshold data from upcoming direct detection experiments to refine the sensitivity projections for the BSM models investigated.
- Exploring the effects of the considered BSM models using alternative neutrino sources, such as nuclear reactors, atmospheric neutrinos, or supernova neutrinos, to enhance discovery potential across different energy regimes.

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8. APPENDICES

CONVENTION AND NOTATION

We have adapted natural units $\hbar = 1$, $c = 1$ throughout this work. In this convention, the Lagrangian density of the Dirac fields

$$\mathcal{L}_{Dirac} = \bar{\psi}(i\not{\partial} + m)\psi, \quad (\text{A.1})$$

where $\not{\partial} = \gamma^\mu \partial_\mu$ represents the Feynman slash. The equation of motion of this Lagrangian is the Dirac equation

$$(i\gamma^\mu \partial_\mu - m)\psi = 0, \quad (\text{A.2})$$

The γ^μ denotes the Dirac gamma-matrices which obey the Clifford algebra

$$\{\gamma^\mu, \gamma^\nu\} = 2g^{\mu\nu}. \quad (\text{A.3})$$

The metric tensor $g^{\mu\nu}$ is defined as

$$g^{\mu\nu} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}. \quad (\text{A.4})$$

Explicitly, the gamma-matrices are

$$\gamma^0 = \begin{pmatrix} \mathbb{1} & 0 \\ 0 & -\mathbb{1} \end{pmatrix}, \quad \gamma^i = \begin{pmatrix} 0 & \sigma^i \\ -\sigma^i & 0 \end{pmatrix}, \quad (\text{A.5})$$

where $\mathbb{1}$ is a 2×2 identity matrix while σ^i , with $i = 1, 2, 3$, are the Pauli matrices

$$\sigma^1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma^2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma^3 = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \quad (\text{A.6})$$

These matrices are Hermitian and satisfy the following relations

$$[\sigma^i, \sigma^j] = 2i\epsilon^{ijk}\sigma^k, \quad \{\sigma^i, \sigma^j\} = 2\delta^{ij}, \quad (\text{A.7})$$

where ϵ^{ijk} and δ^{ij} are the Levi-Civita symbol and Kronecker delta, respectively.

KINEMATICS

In this work, we work in the laboratory frame. The kinematic of the $\nu(p_1) + X(p_2) \rightarrow \nu(p_3) + X(p_4)$ in the rest frame of the target particle is shown in Figure 50. Note that here X can be nucleus or electron. We consider the initial and final mass for the target and incoming neutrino to be the same. Particularly, $m_1 = m_3 = m_\nu$ for neutrino and $m_2 = m_4 = m_X$ for the target. Relations between four-vector momentum $p_i = (E_i, \vec{p}_i)$ with $i = 1, 2, 3, 4$ can be proceed further. By using conservation of momentum and energy, we find

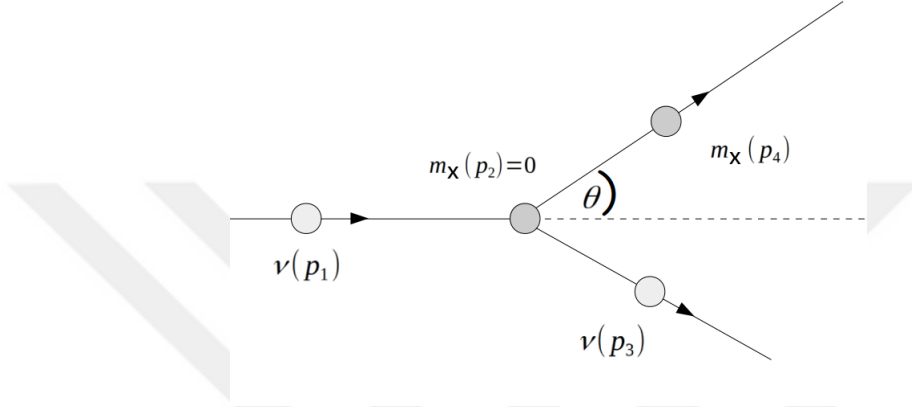


Figure 50. Kinematics for $\nu(p_1) + X(p_2) \rightarrow \nu(p_3) + X(p_4)$ process.

$$\begin{aligned}
 p_1 \cdot p_2 &= p_3 \cdot p_4 = E_\nu m_X, \\
 p_1 \cdot p_3 &= m_\nu^2 + m_X T \approx m_X T \\
 p_2 \cdot p_3 &= p_1 \cdot p_4 = m_X (E_\nu - T) \\
 p_2 \cdot p_4 &= m_X E_4 = m_X (T + m_\ell),
 \end{aligned} \tag{A.8}$$

with the recoil energy transfer, which can be electron or nuclear, is defined by

$$T = E_3 - E_1 = E_4 - E_2. \tag{A.9}$$

Note that for the case of active-sterile neutrino transition magnetic moment, the scattered sterile-neutrino mass is considered. In this case, the following relations are satisfied

$$\begin{aligned}
 p_1 \cdot p_3 &= m_X T + \frac{1}{2} m_4^2, \\
 p_1 \cdot q &= E_\nu^2 - m_X T - \frac{1}{2} m_4^2, \\
 p_3 \cdot q &= m_X T - \frac{1}{2} m_4^2.
 \end{aligned} \tag{A.10}$$

where $q = p_1 - p_3 = p_4 - p_2$ have been used. From the 4-momentum relation, we have

$$|\vec{p}_4| = \sqrt{T(T + 2m_X)}. \tag{A.11}$$

Using $E_3 = E_\nu - T$, $E_4 = T + m_X$, and using this $|\vec{p}_4|$, we find

$$\left(\frac{m_X}{E_\nu - T} - 1\right) = \sqrt{T(T + 2m_X)}. \quad (\text{A.12})$$

Squaring this relation, we obtain the maximum recoil energy as

$$T_{\max} = \frac{2E_\nu^2}{m_X + 2E_\nu}. \quad (\text{A.13})$$

We can also find the minimum energy by noting that $2E_\nu^2 = m_X T + 2E_\nu T$, we find

$$E_\nu^{\min} = \frac{1}{2} \left(T + \sqrt{2m_X T + T^2} \right). \quad (\text{A.14})$$

The general form of differential cross section for 2+2 process satisfies

$$d\sigma = \frac{\langle |\mathcal{M}^2| \rangle}{2E_1 2E_2 |\vec{v}_1 - \vec{v}_2|} \frac{d^3 \vec{p}_3}{(2\pi^3) 2E_3} \frac{d^3 \vec{p}_4}{(2\pi^3) 2E_4} (2\pi)^4 \delta^4(p_1 + p_2 - p_3 - p_4). \quad (\text{A.15})$$

Using the velocity relation $\vec{v} = \vec{p}/E$, the incident flux can be written as

$$2E_1 2E_2 |\vec{v}_1 - \vec{v}_2| = 4(E_2 |\vec{p}_1| - E_1 |\vec{p}_2|), \quad (\text{A.16})$$

where in the center of mass frame, $\vec{p}_1 = -\vec{p}_2$ and $\vec{p}_3 = -\vec{p}_4$, so we can integrate

$$\begin{aligned} \int \int \frac{d^3 \vec{p}_3}{(2\pi)^3 2E_3} \frac{d^3 \vec{p}_4}{(2\pi)^3 2E_4} (2\pi)^4 \delta(E_1 + E_2 - E_3 - E_4) \delta^3(\vec{p}_1 + \vec{p}_2 - \vec{p}_3 - \vec{p}_4) \\ = \frac{|\vec{p}_3|^2 d\Omega}{16\pi^2 E_3 E_4} \int d|\vec{p}_3| \delta(E_{CM} - E_3 - E_4), \end{aligned} \quad (\text{A.17})$$

where now $E_3 = \sqrt{|\vec{p}_3|^2 + m_3^2}$ and $E_4 = \sqrt{|\vec{p}_3|^2 + m_4^2}$, while we have defined $E_{CM} = E_1 + E_2$. Note also that the solid angle is $d\Omega = \sin \theta d\theta d\phi$. The Jacobian of the last line for $x = E_{CM} - E_3(\vec{p}_3) - E_4(\vec{p}_3)$

$$J = \left| \frac{dx}{dp_3} \right| = |\vec{p}_3| \frac{E_3 + E_4}{E_3 E_4}, \quad (\text{A.18})$$

changing variable $dp_3 = dx |J|^{-1}$ leads us to

$$\frac{|\vec{p}_3|^2 d\Omega}{16\pi^2 E_3 E_4} \int dx \delta(x) \left(\frac{E_3 E_4}{|\vec{p}_3| (E_3 + E_4)} \right) = \frac{|\vec{p}_3| d\Omega}{16\pi^2 E_{CM}} \theta(E_{CM} - m_3 - m_4) \quad (\text{A.19})$$

where the delta function gives $E_{CM} = E_3 + E_4$ and note that $\theta(x)$ is a Heaviside function. The differential cross-section then becomes

$$\frac{d\sigma}{d\Omega} = \frac{\langle |\mathcal{M}^2| \rangle}{64\pi^2 E_{CM}^2} \frac{|\vec{p}_3|}{|\vec{p}_1|}. \quad (\text{A.20})$$

Instead of solid angle $d\Omega$, we aim to convert this into nuclear recoil energy dT . We note that $q = p_3 - p_1$, hence $q^2 = p_1^2 + p_3^2 - 2|\vec{p}_1||\vec{p}_3|\cos\theta$ or for massless case $dq^2 = -2|\vec{p}_1||\vec{p}_3|d(\cos\theta)$. So that

$$\frac{d\sigma}{dq^2} = -\frac{d\phi}{2|\vec{p}_1||\vec{p}_3|} \frac{\langle |\mathcal{M}^2| \rangle}{64\pi^2 E_{CM}^2} = -\frac{1}{|\vec{p}_1||\vec{p}_3|} \frac{\langle |\mathcal{M}^2| \rangle}{64\pi E_{CM}^2}. \quad (\text{A.21})$$

Since $q^2 = -2m_X T \rightarrow dq^2 = -2m_X dT$, then

$$\frac{d\sigma}{dT} \frac{dT}{dq^2} = -\frac{1}{2m_X} \frac{d\sigma}{dT}, \quad (\text{A.22})$$

then, since in the center of mass frame $|\vec{p}_1| = |\vec{p}_3|$, we obtain

$$\frac{d\sigma}{dT} = \frac{m_X \langle |\mathcal{M}^2| \rangle}{32\pi E_{CM}^2 |\vec{p}_1|^2}. \quad (\text{A.23})$$

Next, we may square the incident flux of Eq.(A.16) to obtain $(E_2|\vec{p}_1| - E_1|\vec{p}_2|)^2 = E_2^2|\vec{p}_1|^2 + E_1^2|\vec{p}_2|^2 - 2E_2|\vec{p}_1|E_1|\vec{p}_2|$, where the last term can be replaced from relation $(p_1 \cdot p_2)^2 = E_1^2 E_2^2 + |\vec{p}_1|^2 |\vec{p}_2|^2 - 2E_1 E_2 |\vec{p}_1| |\vec{p}_2|$, hence

$$\begin{aligned} (E_2|\vec{p}_1| - E_1|\vec{p}_2|)^2 &= E_2^2|\vec{p}_1|^2 + E_1^2|\vec{p}_2|^2 + (p_1 \cdot p_2)^2 - E_1^2 E_2^2 - |\vec{p}_1|^2 |\vec{p}_2|^2 \\ &= (p_1 \cdot p_2)^2 - (E_1^2 - |\vec{p}_1|^2)(E_2^2 - |\vec{p}_2|^2) \\ &= (p_1 \cdot p_2)^2 - m_1 m_2. \end{aligned} \quad (\text{A.24})$$

Since from the kinematic $p_1 \cdot p_2 = m_X E_\nu$ and set $m_1 = m_\nu \approx 0$, $m_2 = m_X$, in the center of mass frame $\vec{p}_1 = -\vec{p}_2$ we obtain the following relation

$$E_{CM}^2 |\vec{p}_1|^2 = m_X^2 E_\nu^2. \quad (\text{A.25})$$

Hence, the general form of the differential cross-section becomes

$$\frac{d\sigma}{dT} = \frac{\langle |\mathcal{M}^2| \rangle}{32\pi m_X E_\nu^2}, \quad (\text{A.26})$$

where m_X can be electron, m_e , for $\text{E}\nu\text{ES}$ and nucleus, m_N , for $\text{CE}\nu\text{NS}$ process.

FIERZ TRANSFORMATION

Fierz identity is a mathematical tool used to simplify expressions involving products of two Dirac bilinears by rewriting them as a sum of other bilinears. It's particularly useful when dealing with four-fermion interactions, where two pairs of fermions interact at a single point. The general form of the Fierz transformation is given by the expansion

$$\bar{u}_a u_b = \frac{1}{4} \sum_i [\Gamma_i]_{ba} \bar{u} \Gamma_i u. \quad (\text{A.27})$$

The operator Γ_i is the full set of Dirac matrices $\{1, \gamma^\mu, \gamma^5, i\gamma^5 \gamma^\mu, \sigma^{\mu\nu}\}$. For definiteness, we note that we obtain a unity matrix when we applying $(\Gamma_j)_{ab}$ to both sides and noting that $\text{Tr}[\Gamma_i \Gamma_j] = 2g_{ij}$.

In the case of electron-neutrino interaction, the charge current has the following form

$$[\bar{e} \gamma^\mu (1 - \gamma^5) \nu] [\bar{\nu} \gamma_\mu (1 - \gamma^5) e]. \quad (\text{A.28})$$

By explicitly giving indices to these form, we have

$$\bar{e}_a [\gamma^\mu (1 - \gamma^5)]_{ab} \nu_b \bar{\nu}_c [\gamma_\mu (1 - \gamma^5)]_{cd} e_d. \quad (\text{A.29})$$

Hence, using the Fierz transformation

$$\bar{\nu}_c \nu_b = \frac{1}{4} \sum_i (\Gamma_i)_{bc} \bar{\nu} \Gamma_i \nu, \quad \text{and} \quad \bar{e}_a e_d = \frac{1}{4} \sum_j (\Gamma_j)_{da} \bar{e} \Gamma_j e. \quad (\text{A.30})$$

Inserting this to the current leads to

$$\frac{1}{16} \sum_i \sum_j [\gamma^5 (1 - \gamma^5)]_{ab} [(\Gamma_i)_{bc} \bar{\nu} \Gamma_i \nu] [\gamma_\mu (1 - \gamma^5)]_{cd} [(\Gamma_j)_{da} \bar{e} \Gamma_j e] \quad (\text{A.31})$$

Following the same indices, we can then simplify this to be

$$\frac{1}{16} \sum_i \sum_j \text{Tr}[\gamma^\mu (1 - \gamma^5) \Gamma_i \gamma_\mu (1 - \gamma^5) \Gamma_j] [\bar{\nu} \Gamma_i \nu] [\bar{e} \Gamma_j e] \quad (\text{A.32})$$

The next task is to compute the subsequent 16 possible terms, originated from the form of the operator Γ_i . Since the overall term is determined by result of the Tr, we can perform calculation for each term. Most of the terms are zero:

- $\Gamma_i = 1, \Gamma_j = 1$:

$$\text{Tr}[\gamma^\mu (1 - \gamma^5) \gamma_\mu (1 - \gamma^5)] = \text{Tr}[\gamma^\mu (1 - \gamma^5) (1 + \gamma^5) \gamma_\mu] = 0. \quad (\text{A.33})$$

- $\Gamma_i = 1, \Gamma_j = \gamma^\mu$:

$$\text{Tr}[\gamma^\mu (1 - \gamma^5) \gamma_\mu (1 - \gamma^5) \gamma^\rho] = \text{Tr}[\gamma^\mu (1 - \gamma^5) (1 + \gamma^5) \gamma_\mu \gamma^\rho] = 0. \quad (\text{A.34})$$

- $\Gamma_i = i\gamma^5\gamma^\nu, \Gamma_j = \gamma^5$:

$$\begin{aligned}\text{Tr}[\gamma^\mu(1-\gamma^5)i\gamma^5\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^5] &= -i\text{Tr}[\gamma^\mu(1-\gamma^5)(1-\gamma^5)\gamma^\nu\gamma_\mu] \\ &= -2i\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu] = 0.\end{aligned}\quad (\text{A.35})$$

- $\Gamma_i = i\gamma^5\gamma^\nu, \Gamma_j = \sigma^{\rho\sigma}$:

$$\begin{aligned}\text{Tr}[\gamma^\mu(1-\gamma^5)i\gamma^5\gamma^\nu\gamma_\mu(1-\gamma^5)\sigma^{\rho\sigma}] &= -i\text{Tr}[\gamma^\mu(1-\gamma^5)(1-\gamma^5)\gamma^\nu\gamma_\mu\sigma^{\rho\sigma}] \\ &= -2i\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu\sigma^{\rho\sigma}] = 0.\end{aligned}\quad (\text{A.36})$$

- $\Gamma_i = \sigma^{\nu\rho}, \Gamma_j = \sigma^{\sigma\lambda}$:

$$\text{Tr}[\gamma^\mu(1-\gamma^5)\sigma^{\nu\rho}\gamma_\mu(1-\gamma^5)\sigma^{\sigma\lambda}] = i\text{Tr}[\gamma^\mu(1-\gamma^5)(1+\gamma^5)\sigma^{\nu\rho}\gamma_\mu\sigma^{\sigma\lambda}] = 0. \quad (\text{A.37})$$

The non zero terms are only come from the following terms

- $\Gamma_i = \gamma^\nu, \Gamma_j = \gamma^\rho$:

$$\begin{aligned}\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^\rho] &= \text{Tr}[\gamma^\mu(1-\gamma^5)(1-\gamma^5)\gamma^\nu\gamma_\mu\gamma^\rho] \\ &= 2\text{Tr}[(1+\gamma^5)(2g^{\mu\nu} - \gamma^\nu\gamma^\mu)\gamma_\mu\gamma^\rho] \\ &= 4g^{\mu\nu}\text{Tr}[(1+\gamma^5)\gamma_\mu\gamma^\rho] - 2\text{Tr}[(1+\gamma^5)\gamma^\nu\gamma^\mu\gamma_\mu\gamma^\rho] \\ &= -4\text{Tr}[\gamma^\nu\gamma^\rho] = -16g^{\nu\rho}.\end{aligned}\quad (\text{A.38})$$

$$\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu(1-\gamma^5)i\gamma^5\gamma^\rho] = -i\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^\rho] \quad (\text{A.39})$$

$$= +i16g^{\nu\rho}. \quad (\text{A.40})$$

- $\Gamma_i = i\gamma^5\gamma^\nu, \Gamma_j = \gamma^\rho$:

$$\begin{aligned}\text{Tr}[\gamma^\mu(1-\gamma^5)i\gamma^5\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^\rho] &= -i\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^\rho] \\ &= -i16g^{\nu\rho}.\end{aligned}\quad (\text{A.41})$$

- $\Gamma_i = i\gamma^5\gamma^\nu, \Gamma_j = i\gamma^5\gamma^\rho$:

$$\begin{aligned}\text{Tr}[\gamma^\mu(1-\gamma^5)i\gamma^5\gamma^\nu\gamma_\mu(1-\gamma^5)i\gamma^5\gamma^\rho] &= -\text{Tr}[\gamma^\mu(1-\gamma^5)\gamma^\nu\gamma_\mu(1-\gamma^5)\gamma^\rho] \\ &= -16g^{\nu\rho}\end{aligned}\quad (\text{A.42})$$

Collecting the non-zero results, we have

$$\begin{aligned}\frac{1}{16} \Bigg[&-16g^{\nu\rho}[\bar{\nu}\gamma^\nu\nu][\bar{e}\gamma^\rho e] + i16g^{\nu\rho}[\bar{\nu}\gamma^\nu\nu][\bar{e}i\gamma^5\gamma^\rho e] \\ &- i16g^{\nu\rho}[\bar{\nu}i\gamma^5\gamma^\nu\nu][\bar{e}\gamma^\rho e] - 16g^{\nu\rho}[\bar{\nu}i\gamma^5\gamma^\nu\nu][\bar{e}i\gamma^5\gamma^\rho e] \Bigg],\end{aligned}\quad (\text{A.43})$$

and after simplification

$$[\bar{e}\gamma^\nu e][\bar{\nu}\gamma^\nu(1-\gamma^5)\nu] - [\bar{e}\gamma^5\gamma^\nu e][\bar{\nu}(1-\gamma^5\gamma_\nu)\nu], \quad (\text{A.44})$$

and then we obtain

$$[\bar{e}\gamma^\nu(1-\gamma^5)e][\bar{\nu}\gamma_\nu(1-\gamma^5)\nu]. \quad (\text{A.45})$$

This form can be directly added to the neutral current contribution of the amplitude.

We can also use the transformation for the bilinear from charged Higgs contribution to $E\nu\text{ES}$. We have the form

$$[\bar{\nu}^c P_L e][\bar{e}^c P_L \nu]. \quad (\text{A.46})$$

The charged-conjugate fermions satisfies

$$\nu^c = C\bar{\nu}^T, \quad \bar{\nu}^c = -\nu^T C^{-1}. \quad (\text{A.47})$$

With the C operator, we also have $C^{-1}\gamma^\mu C = -(\gamma^\mu)^T$ and $P_L^T = P_L$. Writing indices to the bilinear and expand the form we have

$$\bar{\nu}_a^c [P_L]_{ab} e_b \bar{e}_c^c [P_L]_{cd} \nu_d = [-\nu_{a'}^T (C^{-1})_{a'a}] [P_L]_{ab} e_b [-e_{c'}^T C^{-1}] [P_L]_{cd} \nu_d, \quad (\text{A.48})$$

and then we can simplify it to

$$\nu_a^T [C^{-1} P_L]_{ab} e_b e_c^T [C^{-1} P_L]_{cd} \nu_d = [(C^{-1} P_L)_{ab} (C^{-1} P_L)_{cd}] [\nu_a^T \nu_d] [e_b e_c^T]. \quad (\text{A.49})$$

The first term is a special Fierz identity, that obeys

$$(C^{-1} P_L)_{ab} (C^{-1} P_L)_{cd} = \frac{1}{2} (\gamma^\mu P_L)_{ad} (\gamma_\mu P_L)_{cb}. \quad (\text{A.50})$$

Inserting it back and regrouping leads to

$$\frac{1}{2} [\nu_a^T (\gamma^\mu P_L)_{ad} \nu_d] [e_c^T (\gamma_\mu P_L)_{cb} e_b]. \quad (\text{A.51})$$

Using $\nu^T A \nu = -\bar{\nu}^c C^{-1} A^T C \nu^c$, we then have

$$\frac{1}{2} (\bar{\nu}^c \gamma^\mu P_L \nu^c) (\bar{e}^c \gamma_\mu P_L e^c). \quad (\text{A.52})$$

And noting that $\bar{\nu}^c \gamma_\mu P_L \nu^c = -\bar{\nu} \gamma_\mu P_L \nu$, we obtain

$$\frac{1}{2} (\bar{\nu} \gamma^\mu P_L \nu) (\bar{e} \gamma_\mu P_L e). \quad (\text{A.53})$$

This for is relevant with the SM amplitudes.

CURRICULUM VITAE

He completed undergraduate program at the Department of Physics, Hasanuddin University in 2017. After graduation, he spent one year as an assistant at the same university. He finished his master program at the Department of Physics, Karadeniz Technical University in 2021. This was completed with the scholarship awarded from the Turkish Government through the YTB program. At the same year, he started his doctoral program at the same university.

During his academic journey, the author had attended various scientific activities. He already contributed in eleven SCI publications. In the scope of this dissertation, the list as follows:

- **M. F. Mustamin**, M. Demirci, I.A. Qattan, Nucl. Phys. B 1018 (2025), 117091.
- M. Demirci, **M. F. Mustamin**. Phys.Rev.D 111 (2025) 5, 055032.
- M. Demirci, **M. F. Mustamin**. Eur.Phys.J.C 85 (2025) 1.
- M. Demirci, **M. F. Mustamin**. Phys.Rev.D 109 (2024) 1, 015021.
- **M. F. Mustamin**, M. Demirci. Braz.J.Phys. 51 (2021) 3, 813-819.

He also contributed in thirteen proceedings, some of which have been presented in the following conference in the scope of this dissertation:

- Beyond Standard Model: From Theory to Experiment (BSM-2025), Sabancı University, Türkiye, 6-9 October 2025, Presenter.
- The 22nd Lomonosov Conference on Elementary Particle Physics, Moscow, Russia, 21 - 27 August 2025, Presenter.
- International Conference on Neutrinos and Dark Matter (NuDM-2024), Cairo, Egypt, 11 - 14 December 2024, Presenter, Published proceeding.
- Beyond Standard Model: From Theory to Experiment (BSM- 2023), Al-Ghardaqah, Egypt, 06 - 09 November 2023, Presenter, Published proceeding.
- Turkish Physical Society 39th International Physics Congress, Muğla, Türkiye, 31 August - 04 September 2023, Presenter.

He knows Turkish, English, Bahasa Indonesia, and Bahasa Bugis.