

Coherent and incoherent antineutrino scattering on stable even-even isotopes of molybdenum detectors

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The recent observations of the coherent neutrino- and antineutrino-nucleus scattering have opened up a plethora of opportunities to probe physics within standard and non-standard theories of the electroweak interactions. In the present article, our goal is to explore the possibility of using the molybdenum material as detection medium for coherent and incoherent antineutrino- and neutrino-Mo scattering in the ongoing and future coherent elastic neutrino-nucleus scattering (CE ν NS) experiments by using relevant (anti-) neutrino beams as, e.g., stopped pion-decay neutrino beams, reactor antineutrino beams, astrophysical (solar or supernova) (anti)neutrino beams, etc. Our present coherent and incoherent scattering cross sections of Mo isotopes with neutrinos and antineutrinos are based on the deformed shell model (DSM) that has been previously employed for studying similar processes. On the other hand, in the past, CE ν NS events obtained with this model provided us with better fits to experimental data of COHERENT experimental data compared to phenomenological form factors.

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I. INTRODUCTION

The coherent elastic neutrino-nucleus scattering (CE ν NS) is a Standard Model rare event (due to the very weak neutrino-matter interaction) process, governed by the electroweak interaction, that was predicted long ago (1974) by Freedman [1]. In recent neutrino (ν) and antineutrino ($\bar{\nu}$) detection experiments, due to the very weak neutrino-nucleus interaction, a metric ton or multimetric ton mass scale of the detection medium is required. In the case of the CE ν NS experiments [2] extremely low-energy thresholds (of sub-keV scale) nuclear detectors are employed which drastically enhance the event rates. CE ν NS has proven to be a difficult process to detect due to the fact that the deposited recoil-energy in the nuclear detector is small (of the order of keV).

The first detection of CE ν NS was announced by the COHERENT collaboration [3,4] in 2017, i.e., forty three years after its first prediction by Freedman [1]. The first COHERENT experiment used sodium doped CsI detectors

with a decay at rest (DAR) pion neutrino-source while the second one used Ar target.

The COHERENT measurement of CE ν NS events motivated a steady increase in the experimental sensitivity through the XENONnT [5], LZ [6], and PandaX-4T [7] collaborations which reported their first observations of nuclear recoils from solar ^8B neutrino source via CE ν NS on Xenon and Germanium detectors. By assuming the non existence of new physics, these results provide a measurement of the component ^8B of solar neutrino flux and motivated investigations aiming at probing new physics via nonstandard neutrino interactions (NSI), light mediators, as well as the determination of the weak mixing angle at low momentum transfer [8–12].

Recently, the CONUS experiment [13,14] measured CE ν NS events by using a high-purity Ge detector with a total active mass (3.73 ± 0.02) kg, positioned at a distance 17.1 m from the core center of the nuclear power plant in Brokdorf, Germany, which provided a reactor antineutrino flux (up to $2.3 \times 10^{13} \text{ cm}^{-2} \text{ s}^{-1}$). The ionization energy threshold achieved in this Ge detector was very low (0.21 keV), while the combined likelihood fit performed for this experiment provided an upper limit of 143 antineutrino events (at 90% confidence level). According to CONUS collaboration, this number of events is less than a factor of 2 larger than the predicted events by the SM [13].

In the case of the CONUS + experiment [15,16], high purity Germanium crystals (with 160–180 eV threshold) were employed to measure reactor antineutrino-Ge

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scattering events for the first time. The reactor antineutrinos are of low-energy (up to about 8–10 MeV). The number of events measured by the CONUS + experiment was (395 ± 106) (during operational time of 119 days at the nuclear power plant in Leibstadt, Switzerland) with a statistical significance 3.7σ [15]. The predicted number of events by the Standard Model of the electroweak interactions is equal to (347 ± 59) which means that there is room for new physics beyond the Standard Model [10,17–19].

The possible detection of CE ν NS using both terrestrial and astrophysical neutrino-sources has inspired new constraints on the physics beyond the Standard Model (BSM) and other significant implications for particle physics, astrophysics, and nuclear physics; see for example [20–22].

Nowadays, in the intense experimental effort ongoing today around the world, the CE ν NS process is studied with a variety of neutrino sources and detector technologies by employing various detection materials [23–30] (research from a theoretical viewpoint is also increasing [31–48]). For example, new experiments have been designed, like the NUCLEUS experiment at TU Munich, Germany, aiming for the measurement of CE ν NS in a nuclear reactor using reactor neutrinos (there, the neutrino fluxes have high intensity) at the, relatively, very low energy range (up to 10 MeV) of the nuclear power plant in Chooz, France.

The NUCLEUS experiment is going to utilize dedicated cryogenic detectors with nuclear recoil energy thresholds estimated to be around 20 eV. This is the lowest used in the topic of CE ν NS. At present, the NUCLEUS setup is in the stage of construction (in the shallow underground laboratory at TUM) and commissioning, and it is expected that the experiment will be moved to France next year (a technical run is anticipated to take place in 2026). The scientific potential with several technological details of this experiment can be found, e.g., in [28,49].

Our aim in this paper is to share this effort by providing useful theoretical predictions for both the coherent and incoherent scattering of neutrinos and antineutrinos on the even-even molybdenum isotopes. Toward this aim, we utilize the well known Donnelly-Walecka multipole decomposition method based on the deformed shell model. The latter method, has been adopted previously to study rare event processes like the muon-to-electron conversion [50], direct dark matter detection [51] as well as detection of CE ν NS in several nuclear isotopes [52].

The Molybdenum (with atomic number $Z = 42$) is considered a promising detection medium for rare event processes as is the CE ν NS and has been recently employed by extremely sensitive experiments like the NEMO neutrinoless double beta decay [53] and the MOON direct dark matter detection [54] experiments. As it is known [55], molybdenum has seven stable isotopes with $N = 50, 52, 53, 54, 55, 56, 58$, and abundances as follows: ^{92}Mo (15.86%), ^{94}Mo (9.12%), ^{95}Mo (15.70%), ^{96}Mo (16.50%), ^{97}Mo (9.45%), ^{98}Mo (23.75%), and ^{100}Mo (9.62%).

In our present work, we have considered all the even-even isotopes $^{92,94,96,98,100}\text{Mo}$. Recently, the AMoRE collaboration has been searching for neutrinoless double beta decay of the ^{100}Mo using molybdate scintillating crystals (100 kg of enriched ^{100}Mo) via low temperature thermal calorimetric detection [56,57].

We note that CE ν NS detectors could also be used as dark matter (DM) detectors as has been pointed out in Ref. [58] where the anticipation of similar experimental challenges is discussed.

Within the context of the multipole expansion Donnelly-Walecka method [59,60], we calculate coherent elastic (anti)neutrino-Mo scattering cross sections, for ground state to ground state ($|gs\rangle \rightarrow |gs\rangle$) transitions as well as incoherent (inelastic) cross sections where the nucleus from the ground state transitions to an excited (higher energy) state ($|gs\rangle \rightarrow |f\rangle$), where $|f\rangle$ represents a low-lying (up to about 10 MeV) excited state. By using these cross sections total cross sections (total = coherent + incoherent) are readily obtained. Moreover, the important ratio of the coherent over the total cross section, $\rho = \text{coherent}/\text{total}$, for each Mo isotope is calculated. The latter quantity provides the portion of the total rate (unknown and appreciably difficult to measure) represented by the measured coherent rate in the CE ν NS experiments.

The required ground ($|gs\rangle$) and final ($|f\rangle$) states for the even-even Mo isotopes, $^{92,94,96,98,100}\text{Mo}$, are deduced through nuclear structure calculations by utilizing the deformed shell model (DSM) as is explained below. It should be noted that, the DSM eigenfunctions, with the corresponding eigenenergies for the ground and various excited states in $^{94,96,98,100}\text{Mo}$ have been determined in our previous publications [31–33]. Here, we will first present the DSM results for the remaining ^{92}Mo isotope and, then, we will proceed with a comparative analysis of these results through the computed ratio $\rho = \text{coherent}/\text{total}$ for all even-even Mo isotopes. We note that the results for neutrino scattering off nuclei differ from those of antineutrino scattering due to the cross-term “polar vector times axial vector” that appears in the interaction Hamiltonian [35].

In the rest of the paper the material is organized as follows. At first (Sec. II), the neutrino-nucleus scattering cross section formalism based on the Donnelly-Walecka method is briefly summarized. Also the deformed shell model (DSM) relying on Hartree-Fock deformed intrinsic states with angular momentum projection and band mixing, which provides the low-lying nuclear states required for our present calculations is briefly described. Following this, in Sec. III, initially, cross section results for ^{92}Mo are presented in detail. Then, in Sec. IV, a comparative analysis of the coherent and incoherent scattering cross section results of neutrinos and antineutrinos on the even-even molybdenum isotopes is presented and discussed in conjunction with the important quantity ρ . Finally, in Sec. V,

the conclusions extracted from the present study are briefly summarized.

II. FORMALISM OF NEUTRAL-CURRENT ν -NUCLEUS REACTIONS

A. Neutrino-nucleus scattering cross section within Donnelly-Walecka method

The relevant formulation that provides the neutral current ν -nucleus scattering differential cross sections as a function of the incoming neutrino energy ε_ν have been discussed in detail earlier in Refs. [31,59–61]. A few important steps are given below for the sake of completeness. In the low energy domain with the leptonic and hadronic currents represented by $\hat{j}_\mu^{\text{lept}}$ and $\hat{J}^\mu$, respectively, the Standard Model weak interaction neutrino-nucleus Hamiltonian $\hat{H}_I$ is written in the effective current-current form as

$$\hat{H}_I = -\frac{G}{\sqrt{2}} \int d^3\mathbf{x} \hat{j}_\mu^{\text{lept}}(\mathbf{x}) \hat{J}^\mu(\mathbf{x}). \quad (1)$$

In the above definition, G is the Fermi weak coupling constant. The double differential cross section for the scattering of low energy neutrinos from even-even Molybdenum isotopes for a transition from the initial state $|i\rangle$ represented by $|J_i, M_i\rangle$ to the final state $|f\rangle$ represented by $|J_f, M_f\rangle$ can be written as

$$\frac{d^2\sigma_{i\rightarrow f}}{d\Omega dw} = (2\pi)^4 \varepsilon_f^2 \sum_{s_f, s_i, M_f, M_i} \frac{|\langle f | \hat{H}_I | i \rangle|^2}{(2J_i + 1)} \quad (2)$$

where the sums are over the initial and final spin states (s_i, s_f) and magnetic quantum numbers (M_i, M_f), respectively. Applying a multipole analysis on the weak hadronic current following Donnelly-Walecka method [59,60], the neutrino-nucleus scattering cross section for the excitation energy ω of the target nucleus becomes

$$\frac{d^2\sigma_{i\rightarrow f}}{d\Omega d\omega}(\phi, \theta, \omega, \varepsilon_i) = \delta(E_f - E_i - \omega) \frac{2G^2 \varepsilon_f^2 \cos^2(\theta/2)}{\pi(2J_i + 1)} \times [C_V + C_A \mp C_{VA}] \quad (3)$$

Because of the energy conservation, ω is equal to the difference of the initial and final states of the nucleus. It is also equal to the difference of the incoming and outgoing energies ε_i and ε_f of the neutrinos, i.e., $\omega = E_f - E_i = \varepsilon_i - \varepsilon_f$. In Eq. (3), the $(-)$ sign corresponds to the scattering of the neutrinos and the $(+)$ to the scattering of the antineutrinos.

The terms C_V and C_A in Eq. (3) include a summation over the contributions coming from the polar-vector and axial-vector multipole operators, respectively and are given by

$$C_{V(A)} = \sum_{J=0}^{\infty} |\langle J_f | \hat{M}_J^{(5)}(q) + \frac{w}{q} \hat{L}_J^{(5)}(q) | J_i \rangle|^2 + \sum_{J=1}^{\infty} \left(-\frac{q_\mu^2}{2q^2} + \tan^2 \frac{\theta}{2} \right) |\langle J_f | \hat{T}_J^{\text{mag}(5)}(q) | J_i \rangle|^2 + |\langle J_f | \hat{T}_J^{\text{el}(5)}(q) | J_i \rangle|^2 \quad (4)$$

Similarly, the interference term C_{VA} in Eq. (3) contains the product of transverse polar-vector and transverse axial-vector matrix elements and is given by

$$C_{VA} = 2 \tan \frac{\theta}{2} \left(-\frac{q_\mu^2}{q^2} + \tan^2 \frac{\theta}{2} \right)^{1/2} \times \sum_{J=1}^{\infty} \text{Re} \langle J_i | \hat{T}_J^{\text{mag}}(q) | J_i \rangle \langle J_f | \hat{T}_J^{\text{el}}(q) | J_i \rangle^* \quad (5)$$

In the above equations, the superscript 5 refers to the axial vector components of the hadronic current and they are given in the Appendix. C_{VA} contains contributions of $\hat{T}_J^{\text{el}}$ and $\hat{T}_J^{\text{mag}5}$ operators for normal parity transitions, while for abnormal parity ones it contains matrix elements of $\hat{T}_J^{\text{mag}}$ and $\hat{T}_J^{\text{el}5}$.

The square of the four-momentum transfer and the three momentum transfer, the scattering angle θ which occur in the above equations can be expressed as

$$q_\mu^2 \equiv q_\mu q^\mu = -4\varepsilon_i(\varepsilon_i - \omega) \sin^2(\theta/2) \\ q \equiv |\mathbf{q}| = [\omega^2 + 4\varepsilon_i(\varepsilon_i - \omega) \sin^2(\theta/2)]^{1/2} \quad (6)$$

It is important to mention that in practice, we are interested in the differential cross section $d\sigma_{i\rightarrow f}/d\omega$ and this is obtained by integrating out ϕ and θ in Eq. (3). Note that the C_V , C_A , and C_{VA} also depend on θ . The reduced matrix elements that enter in the formulas for $C_{V(A)}$ and C_{VA} are evaluated within the deformed shell model (DSM) described ahead [62] for the nuclear structure part. Now we will briefly describe DSM.

B. Deformed shell model

The deformed shell model or DSM is based on Hartree-Fock (HF) deformed intrinsic states with angular momentum projection and band mixing. Over the years this model has been able to provide a good description of the properties of nuclei in the mass range $A = 60\text{--}90$ [62] like the spectroscopy of $N = Z$ odd-odd nuclei with isospin projection see for example [63], double beta decay half-lives [64,65], $\mu - e$ conversion in the field of the nucleus [50], etc. In recent years, the model has been applied to study neutrino-nucleus scattering, see for example our recent papers [31–33]. The details of this model have been described in many of our earlier publications, see for example [62].

Assuming axial symmetry, for a given nucleus, starting with a model space consisting of a given set of single particle (sp) orbitals and effective two-body Hamiltonian (TBME + spe), the lowest energy intrinsic states are obtained by solving the HF single particle equation self-consistently. Excited intrinsic configurations are obtained by making particle-hole excitations over the lowest intrinsic state. It is worth noting that intrinsic states $\chi_K(\eta)$ do not have definite angular momenta. Hence states of good angular momentum projected from an intrinsic state $\chi_K(\eta)$ can be written in the form

$$\psi_{MK}^J(\eta) = \frac{2J+1}{8\pi^2\sqrt{N_{JK}}} \int d\Omega D_{MK}^{J*}(\Omega) R(\Omega) |\chi_K(\eta)\rangle, \quad (7)$$

where N_{IK} is the normalization constant given by

$$N_{JK} = \frac{2J+1}{2} \int_0^\pi d\beta \sin\beta d_{KK}^J(\beta) \langle \chi_K(\eta) | e^{-i\beta J_y} | \chi_K(\eta) \rangle. \quad (8)$$

In Eq. (7), Ω represents the Euler angles (α, β, γ) , while $R(\Omega) = \exp(-i\alpha J_z) \exp(-i\beta J_y) \exp(-i\gamma J_z)$ represents the general rotation operator. The good angular momentum states projected from different intrinsic states are not in general orthogonal to each other. Hence they are orthonormalized and then band mixing calculations are performed. The resulting eigenfunctions are of the form

$$|\Phi_M^J(\eta)\rangle = \sum_{K,\alpha} S_{K\eta}^J(\alpha) |\psi_{MK}^J(\alpha)\rangle. \quad (9)$$

The reduced matrix elements occurring in Eqs. (4) and (5) are evaluated using the band mixed wave functions Φ_M^J defined in Eq. (9). Now, we will present the DSM results for ^{92}Mo .

III. RESULTS AND DISCUSSION

Initially, in our present calculations, we concentrated on the coherent and incoherent scattering cross sections (of neutrino and antineutrino) for the even-even molybdenum isotopes, focusing on the ^{92}Mo which had not been studied in detail previously [31,33]. Toward this purpose, we first derived and tested the DSM wave functions for the ground state and for the low-lying excited states (up to about $\omega = 15$ MeV) of this isotope. Afterward, by employing the wave functions deduced in our recent publications [31,33], we computed individual cross sections of antineutrino scattering on the $^{94,96,98,100}\text{Mo}$ isotopes, as described below.

A. The nuclear structure of ^{92}Mo isotope

The effective nucleon-nucleon interaction employed for the ^{92}Mo isotope was the *GWBG* with the ^{66}Ni as the inert core. The steps followed for its construction are described

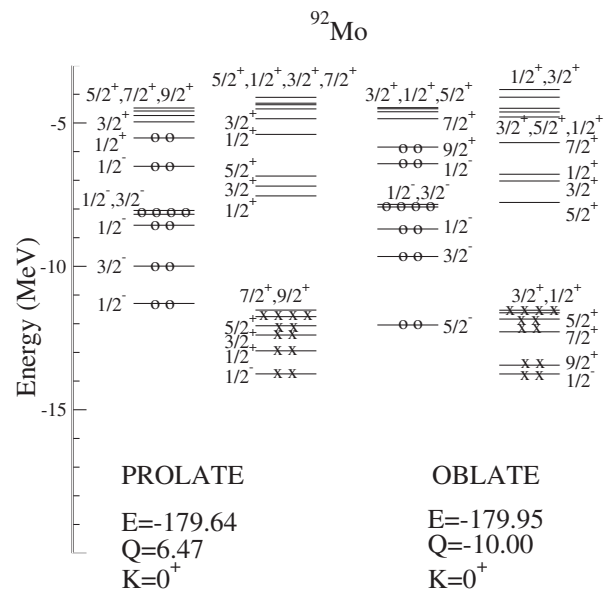


FIG. 1. HF single-particle spectra for ^{92}Mo corresponding to lowest energy prolate and oblate configurations. In the figure, circles represent protons and crosses represent neutrons. The HF energy E in MeV, mass quadrupole moment Q in units of the square of the oscillator length parameter and the total azimuthal quantum number K are given in the figure.

in detail in Ref. [66]. This interaction has been successfully utilized in our previous studies [31,33] by assuming that the active orbits for the protons are the $0f_{5/2}$, $1p_{3/2}$, $1p_{1/2}$, and $0g_{9/2}$, with single-particle energies of -5.322 , -6.144 , -3.941 , and -1.250 MeV, while for the neutrons the active orbits are $1p_{1/2}$, $0g_{9/2}$, $0g_{7/2}$, $1d_{5/2}$, $1d_{3/2}$, and $2s_{1/2}$. The single-particle energies for the first five neutron orbits are taken to be -0.696 , -2.597 , 5.159 , 1.830 , and 4.261 MeV, respectively. The $2s_{1/2}$ neutron orbit produces low-lying large deformed solutions even though molybdenum isotopes are known to be weakly deformed. Hence, the effect of this orbit is eliminated by taking the corresponding neutron single-particle energy at a high value, just as in Refs. [31,33].

As described in Sec. II, we first obtain the lowest HF configuration by performing an axially symmetric HF calculation. Then, various excited configurations are obtained by making particle-hole excitations over this lowest HF configuration. The lowest energy prolate and oblate HF energy spectrum is shown in Fig. 1. For both configurations, the intrinsic quadrupole moments are small indicating that this nucleus has small deformation. Both the prolate and oblate configurations are almost degenerate, with the oblate solution appearing lower by around 0.3 MeV. Hence, there is good mixing between prolate and oblate configurations.

Similar prolate-oblate mixing has been seen in the four stable even-even $^{94,96,98,100}\text{Mo}$ isotopes studied in Refs. [31–33] establishing the fact that the Mo isotopes

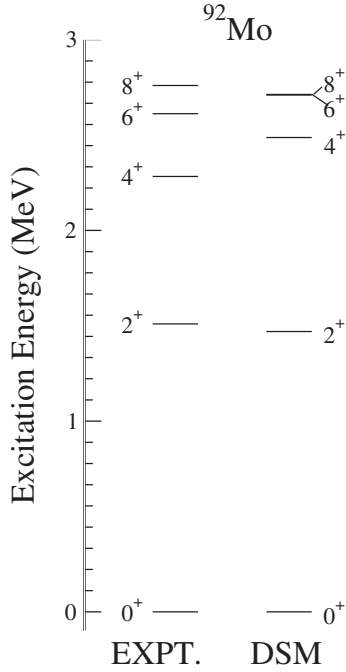


FIG. 2. The ground band observed for ^{92}Mo is compared with the DSM predicted values. The experimental data are taken from [67].

are transitional nuclei. By particle-hole excitations over the lowest energy configurations, we have generated 16 intrinsic states of prolate shape and 12 intrinsic states of oblate shape for ^{92}Mo . Good angular momentum states are projected from each of these intrinsic states, and then a band mixing calculation is performed as described before.

The calculated yrast-levels are compared with experimental data in Fig. 2. As can be seen from this figure the agreement is reasonable. Since ^{92}Mo is a neutron closed-shell nucleus with $N = 50$, the separation of ground state 0^+ to the first excited state 2^+ is large, a feature that is nicely reproduced in our calculation (see Fig. 2). Then, there are three close-lying levels 4^+ , 6^+ , and 8^+ , which are also reasonably well reproduced.

For further testing of the reliability of the wave functions generated by the DSM method, we have also calculated the magnetic moment of 2_1^+ level as well as the $B(E2)$ values for some of the low-lying excited states. Using the bare gyromagnetic ratios for protons and neutrons, the calculated value for the magnetic moment of the 2_1^+ state is $2.8 \mu_N$ which, compared to the experimental value $2.3 \mu_N$, shows that the agreement is reasonable. A slight adjustment of the values of gyromagnetic ratios would have yielded a much better agreement. Such adjustments have been made in many theoretical studies [68]. We note that, the magnetic moments for higher J -states have not been measured.

With respect to the $B(E2)$ values of ^{92}Mo , two transitions, namely, the $2_1^+ \rightarrow 0_1^+$ and the $4_1^+ \rightarrow 2_1^+$, with values of 8.4 Wu and <24 Wu, respectively, have been

experimentally observed. The calculated values for these two transitions are 8.5 Wu and 17.9 Wu, in good agreement with experimental data. We mention that, in the calculation of $B(E2)$ values, we have chosen the effective charges for protons and neutrons to be $e_p = 1.6e$ and $e_n = 1.0e$ as in our previous studies [31,62].

Relying on the DSM wave functions tested as described above, we proceeded further and calculated the coherent and incoherent neutrino- and antineutrino- ^{92}Mo scattering differential cross sections as discussed below. Moreover, by employing the DSM wave functions derived previously [31,33], the incoherent antineutrino scattering cross sections for the even $^{94,96,98,100}\text{Mo}$ isotopes were also computed. We remind that the corresponding neutrino scattering cross sections on the latter even Mo isotopes have been published in Refs. [31,33].

B. Incoherent cross section calculations for ^{92}Mo ν -detector

For the incoherent channel the calculations start from the double differential cross sections $d^2\sigma/d\Omega d\omega$ of Eq. (3) computed state-by-state within the DSM method. Subsequently, by integrating over the angles θ and ϕ , the cross sections $d\sigma/d\omega$ are evaluated.

Because cross sections for high- J^π multipole states (including also all the excited 0^+ states) are found to be negligibly small, the incoherent differential cross sections discussed in this work, originate from the low-lying multipole states for specific incoming neutrino energies ε_ν .

In Fig. 3, the results for the cross sections $d\sigma/d\omega$ are shown for the incoming neutrino energy of $\varepsilon_\nu = 15$ MeV (or antineutrino energy $\varepsilon_{\bar{\nu}} = 15$ MeV). We have chosen to illustrate the dependence of $d\sigma/d\omega(\omega)$ for the value of incoming neutrino energy $\varepsilon_\nu = 15$ MeV (or antineutrino $\varepsilon_{\bar{\nu}} = 15$ MeV), because it is slightly higher than the maximum energy of solar neutrinos (ν_e) and reactor antineutrinos ($\bar{\nu}_e$). We further mention that, in supernova (SN) neutrino simulations, the mean energies chosen for ν_e and $\bar{\nu}_e$ (low-energy SN neutrinos) are between 12 MeV and 15 MeV. This means that, roughly speaking, our results may be used for estimations in solar, reactor, and low-energy supernova neutrinos.

In the present study of the ^{92}Mo isotope, for $\varepsilon_\nu = 15$ MeV (or $\varepsilon_{\bar{\nu}} = 15$ MeV), we found (see Fig. 3) that in both reactions, $^{92}\text{Mo}(\nu, \nu')^{92}\text{Mo}^*$ (left panels) and $^{92}\text{Mo}(\bar{\nu}, \bar{\nu}')^{92}\text{Mo}^*$ (right panels), the pronounced peaks correspond to $J^\pi = 1^+, 2^+, 1^-,$ and 2^- . Note that, as has also been shown in Ref. [35], the dominance of the incoherent peaks changes rather significantly with the chosen incoming neutrino/antineutrino energy. This can be implied by comparing the upper panel of Fig. 3 (resulted with $\varepsilon_\nu = 15$ MeV or $\varepsilon_{\bar{\nu}} = 15$ MeV) and the lower panel (resulted with $\varepsilon_\nu = 20$ MeV or $\varepsilon_{\bar{\nu}} = 20$ MeV). As can be seen from Fig. 3, the incoherent cross section originates mostly from the 1^+ states and, specifically, from the two 1^+

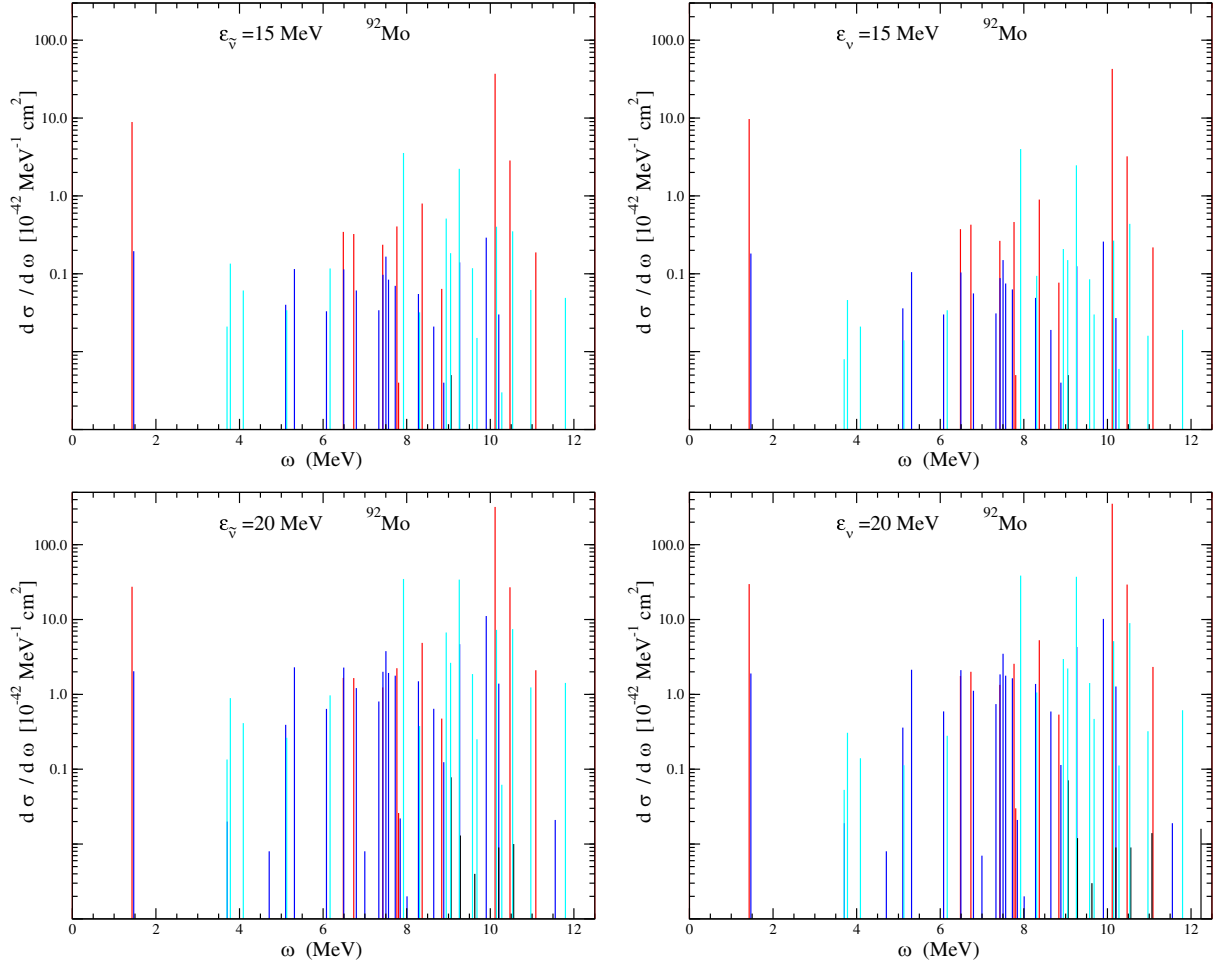


FIG. 3. The differential cross section as a function of the excitation energy ω for ^{92}Mo isotope: (a) Upper panel, at incoming antineutrino (left) energy $\varepsilon_{\bar{\nu}} = 15$ MeV and neutrino (right) energy $\varepsilon_{\nu} = 15$ MeV, and (b) Lower panel, at incoming antineutrino (left) energy $\varepsilon_{\bar{\nu}} = 20$ MeV and neutrino (right) energy $\varepsilon_{\nu} = 20$ MeV, for different excited states of the target nucleus. The contribution of the excitation to $J = 1^+$ state is represented in red, to $J = 2^+$ in blue, to $J = 1^-$ in black, and to $J = 2^-$ in cyan. The left side figures are for antineutrino scattering and the right side ones for neutrino scattering.

states at 1.43 MeV and at 10.47 MeV. These two excited states generate almost the total cross section due to the low-lying excitations of the ^{92}Mo isotope.

The results illustrated in Figs. 3 and 4 (also those published in Refs. [31–33]) show that the Mo-isotopes present rich responses in the excitation energy range $\omega \leq 15$ –20 MeV. This energy region includes transitions to the bound nuclear states and is relevant for solar neutrinos, reactor neutrinos, and geo-neutrinos, but also for the low-energy supernova neutrinos. Thus, the calculated (anti)neutrino-Mo isotopes cross sections are suitable for use also in astrophysical neutrino simulations (e.g., interpretation of neutrino oscillations, neutrino properties, etc.).

It is worth mentioning that in these calculations we used the quenching factor 0.35 (defined by Kay *et al.* [69]), which had also been utilized in our previous studies [31,33].

By utilizing the neutrino energy distributions of the important neutrino sources mentioned above, folded (differential) cross sections may be obtained that represent the response of the nuclear detector of interest to the specific neutrino sources (see, e.g., [35,36]).

IV. COMPARATIVE ANALYSIS OF INCOHERENT CROSS SECTIONS

The cross section results presented here for (anti)neutrino scattering on ^{92}Mo , may now be discussed in conjunction with the results we have reported earlier in Refs. [31,33] for the four stable even-even Mo isotopes with $A = 94, 96, 98$, and 100 . As is known, the natural abundances of the stable even-even molybdenum isotopes are: for the ^{92}Mo 15.86%, for the ^{94}Mo 9.12%, for the ^{96}Mo 16.50%, for the ^{98}Mo 23.75%, and for the ^{100}Mo 9.62%, as mentioned in Sec. I. This means that, for a pure (nonenriched) molybdenum neutrino-detector material,

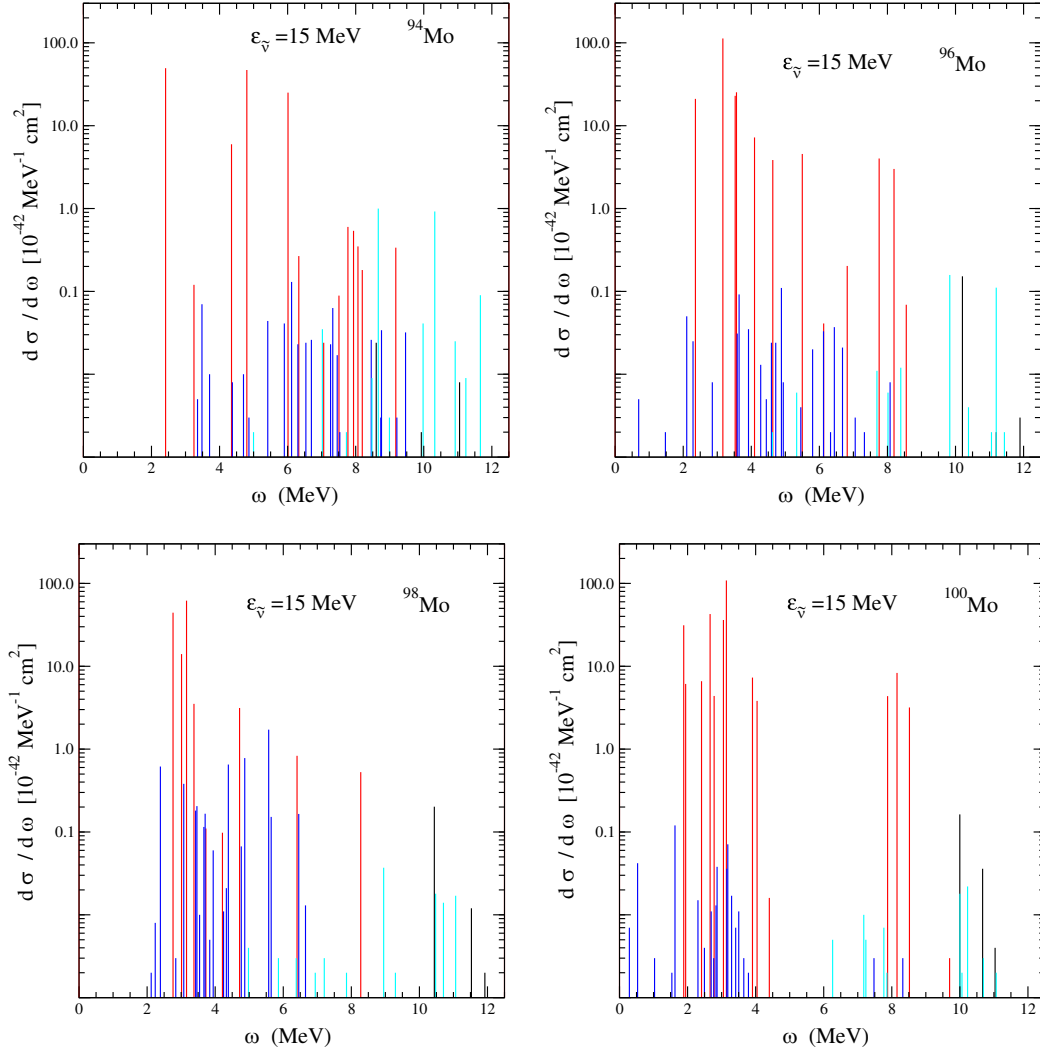


FIG. 4. The differential cross section as a function of the excitation energy ω for $^{94,96,98,100}\text{Mo}$ at incoming antineutrino energy $\varepsilon_{\bar{\nu}} = 15$ MeV for different excited states. The contribution of the excitation to $J = 1^+$ states is represented in red, to $J = 2^+$ in blue, to $J = 1^-$ in black, and to $J = 2^-$ in cyan.

the even-even isotopes represent nearly the 75% of the detector mass. On the other hand, the only stable odd Mo isotopes are the ^{95}Mo and ^{97}Mo , with natural abundances of 15.70% and 9.45%, respectively.

Even in the incoherent part, the difference is no more than 10% as seen from the numbers in Table I and also from Fig. 4.

We will now consider the incoherent part in more detail, as this is significant in the context of BSM, and consider only the antineutrino scattering cross sections.

The results illustrated in Fig. 3 for antineutrino scattering on ^{92}Mo are similar to those for $^{94,96,98,100}\text{Mo}$ shown in Fig. 4.

For the sake of comparison, in Fig. 5 we illustrate the scattering differential cross sections leading to the individual multipolarities 1^+ , 2^+ , 1^- , and 2^- . It is clear from Fig. 5 that most (>90%) of the cross section originates from the 1^+ excitations. This implies that, the neutrino scattering on

the natural Mo detector medium is expected to probe mostly the Gamow-Teller like 1^+ states.

Proceeding further, in Fig. 6, we show the cross section for the dominant 1^+ states as a function of their excitation

TABLE I. Coherent, incoherent and total differential cross sections (in units of $10^{-42} \text{ MeV}^{-1} \text{ cm}^2$) of neutrino and anti-neutrino scattering off the $^{92,94,96,98,100}\text{Mo}$ detector isotopes. The portion of the coherent into the total cross section is also listed.

Isotope	Coherent	Incoherent		Total Cross Section		Ratio ρ (%)	
		(ν)	$(\bar{\nu})$	(ν)	$(\bar{\nu})$	(ν)	$(\bar{\nu})$
^{92}Mo	998.6	67.7	63.1	1066.3	1061.7	93.65	94.05
^{94}Mo	1165.0	150.4	132.7	1315.4	1297.7	88.57	89.77
^{96}Mo	1338.5	230.0	206.1	1568.5	1544.6	85.34	86.66
^{98}Mo	1506.2	147.7	134.2	1653.9	1640.4	91.07	91.82
^{100}Mo	1692.9	290.0	263.5	1982.9	1956.4	85.37	86.53

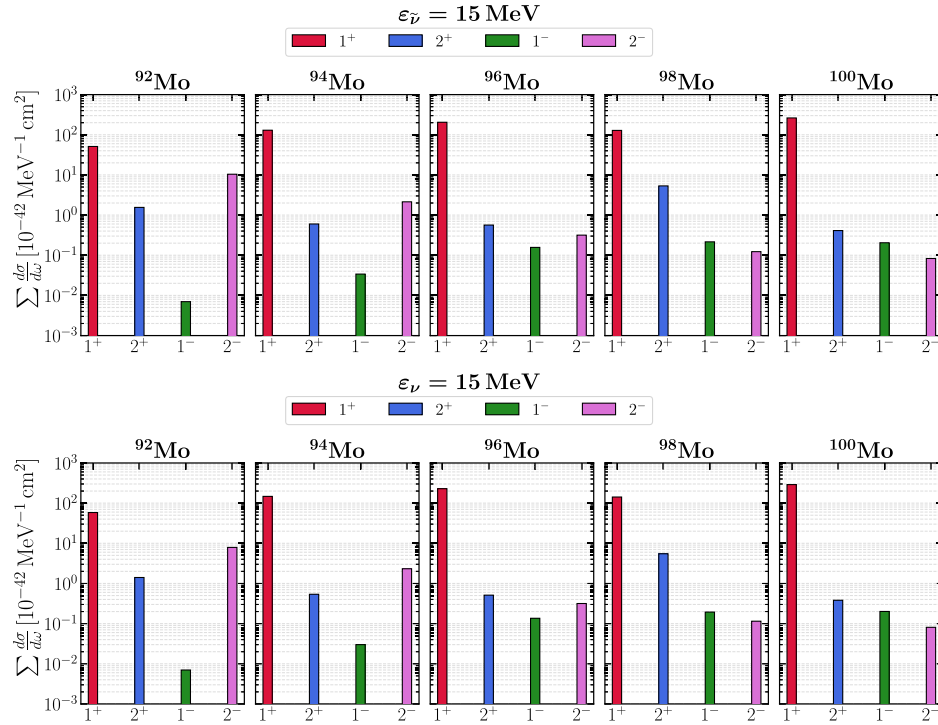


FIG. 5. Summed antineutrino-nucleus (upper panel) and neutrino-nucleus (lower panel) differential scattering cross sections for the 1^+ , 2^+ , 1^- , and 2^- states for the five isotopes $^{92,94,96,98,100}\text{Mo}$.

energies. By combining the results in Fig. 5 and Fig. 6, we conclude that, the cross section strength is mainly concentrated in three to four pronounced 1^+ levels and that this holds for all the five stable even-even Mo isotopes.

We note that, the cross sections for antineutrinos in the above figures differ by 10% (see Table I) from the

corresponding neutrino-Mo isotopes cross sections of Refs. [31,33].

V. COHERENT, TOTAL CROSS SECTIONS AND THE RATIO ρ

In Table I, we list the coherent cross section, the sum of the cross section contributions coming from nearly all the excited states, called incoherent cross sections, and the total (= coherent + incoherent) cross section for both cases: neutrino-Mo and $\bar{\nu}$ -Mo isotopes neutral current scattering reactions. By exploiting the results of Table I we compare the coherent, incoherent, and total cross sections for the five stable even-even Mo-isotopes and calculate the ratio of the coherent divided by the total cross section, i.e., the quantity

$$\rho = \frac{\text{Coherent}}{\text{Total}}. \quad (10)$$

This is interesting from an experimental point of view since, in general, the quantity ρ represents the portion of the total cross section that is measured in the CE ν NS experiments.

It is important to note that the coherent cross sections for neutrino-nucleus and antineutrino-nucleus scattering are identical (only in the incoherent cross sections are different). Furthermore, as can be readily found the coherent cross sections show a good correlation to the square of the neutron number ($\sim N^2$) of each isotope.

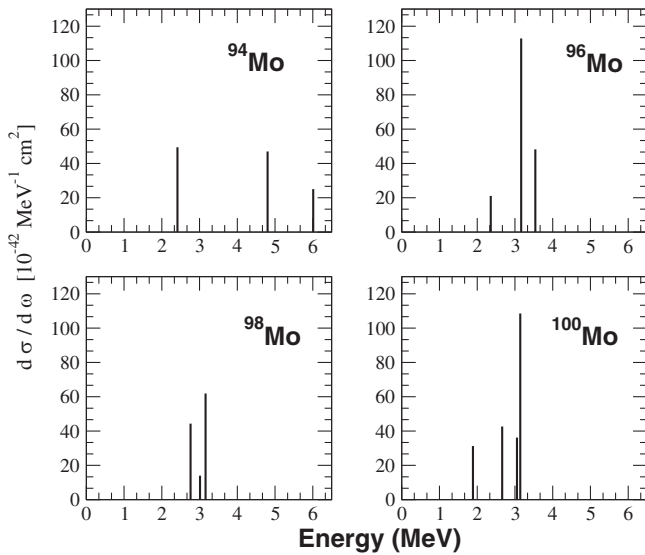


FIG. 6. The antineutrino nucleus differential scattering cross section for the dominant 1^+ states as a function of their excitation energy for the four isotopes $^{94,96,98,100}\text{Mo}$.

It should be stressed that the results listed in Table I for Mo-(anti)neutrino scattering show Standard Model predictions. As has been observed by the CEνNS experiments, the number of the measured events is larger (both for the neutrino and antineutrino sources) than the SM predictions. Thus, by invoking BSM physics one may choose several scenarios, such as neutrino magnetic moments, neutrino millicharges, light neutrino bosons, etc. Then, in view of future measurements of CEνNS events on Mo detectors one may compare our SM predictions with the SM + BSM ones on the one side and the experimental data on the other.

In the case of the Ge detector medium, such calculations have been published and the reader is referred to Refs. [10,13,19,22]. It is worth mentioning that, in general, even though the neutral-current neutrino-nucleus scattering cross sections evaluated here, are substantially smaller than those of the charged-current neutrino-nucleus scattering ones, the former may provide important information that is attributed to the following reasons. In the low-energy range of our calculations, $\epsilon_\nu \lesssim 15$ MeV (also for energies $\epsilon_\nu \lesssim 100$ MeV), neutrinos are not able to produce the massive leptons (μ , τ) in the detector, which means that in low-energy neutrino detectors the portion of neutrinos that will participate in charged-current scattering is limited.

Moreover, charged-current antineutrino scattering, in medium heavy nuclei (as is the case of Mo-isotopes) as well as heavy nuclei, is suppressed due to Pauli blocking effects. Hence, only electron-neutrino charged-current reactions are significant for such detectors. On the other hand, in neutral-current scattering of (anti)neutrinos on, e.g., Mo detectors studied in this work, specifically the heavy flavor neutrinos (ν_x , with $x = \mu, \tau$) can also be detected.

Regarding the main uncertainties in neutral current neutrino-nucleus scattering cross sections or event rates, in the case of the measured CEνNS channel (where the cross sections are essentially proportional to N^2) they generally stem from the weak neutron form factor, the neutrino flux, and the effects related to the specific detector employed. For intense neutrino fluxes required for CEνNS experiments, the number of neutrinos produced by the source is a major uncertainty and is expected to decrease as detectors like those used at the Spallation Neutron Source (SNS) at Oak Ridge [2,3] improve their measurement capabilities (future experiments aim for around a 2%–3% improvement). The source of uncertainty related to the effects of the specific detector employed is that, the relationship between the detected nuclear recoil energy and the true recoil energy, is not perfectly known. On the other hand, uncertainties in the detector's quenching factor for nuclear recoils require improved measurements of these factors, which are crucial for achieving high precision. Currently, detector-related uncertainties are reduced through careful calibration, which is essential for precise measurements [21,41].

Focusing on the theoretical cross-section predictions, the main uncertainties for the coherent calculations originate from the nuclear model (random phase approximation, shell model, relativistic Fermi gas model, etc.) utilized to compute the weak nuclear form factor for the neutron (mainly) and proton distributions within the nucleus, which is a significant source of error. The calculations relying on various nuclear models usually lead to somewhat different predictions, especially when they are based on effective interactions or adjustments in single-particle parameters. Predictions based on the shell model or the deformed shell model employed in the present work may be uncertain due to factors like the uncertainties in the input values for magnetic transitions and the accuracy of the predicted shell-model states themselves. Advanced nuclear structure calculations usually employ shell models to better constrain the form factor uncertainties.

It is worth mentioning that the axial-vector coupling parameter g_A used in nuclear physics is considered uncertain largely due to effects from the nuclear medium (in the complex nuclear environment, g_A is an effective coupling parameter) that quench the value from its free-nucleon value $g_A = 1.27641$, which is highly precise. In many neutrino interaction models, the axial-vector coupling parameter is not precisely known and influences the cross-section calculations [68]. Nuclear many-body effects and interactions with the nuclear medium lead to a quenching of the axial-vector coupling; i.e., the effective value is smaller than the free-nucleon value. We also mention that at higher neutrino energies (e.g., high-energy supernova neutrinos), the uncertainties of theoretical calculations can be more complex due to partial coherence. Finally, when including BSM physics, the challenge of distinguishing a signal of physics beyond the Standard Model from Standard Model theoretical uncertainties is one of the main hurdles for precision experiments; thus, the effort focuses on improving experimental precision while simultaneously reducing theoretical uncertainties.

VI. SUMMARY AND CONCLUSIONS

In this work, relying initially on nuclear structure calculations performed within the deformed shell model (DSM), we computed coherent and incoherent (anti)neutrino- ^{92}Mo scattering cross sections. For natural (nonenriched) detection media used in the CEνNS experiments, is however, one should compute the CEνNS cross sections (events) coming from each of the stable Mo isotopes. Then, through the abundances of the stable isotopes in the natural molybdenum, one may obtain the (theoretical) number of events separately for the coherent channel (to be compared with those measured by the Mo detector) and the incoherent channel. Theoretically, as shown in Table I, one may estimate the portion of the neutral-current (anti)neutrino-nucleus scattering events (expected to be measured by

CEvNS experiments). In this work, we found that this ratio is in between 86% and 94%.

From an experimental viewpoint, the measured CEvNS events with accelerator pion decay-at-rest (DAR) neutrino beam on a Ge detector (at the SNS source) have been comparatively analyzed in conjunction with those measured by using a reactor antineutrino source on a common Ge detector recently. These complementary results offer the possibility of a combined analysis of CEvNS results on a germanium target and a determination of important nuclear physics parameters [22].

The present cross sections of Mo-(anti)neutrino scattering are Standard Model predictions. In case of the existence of room for physics beyond the Standard Model (BSM), as shown from other detector media, one may choose various BSM theories (neutrino magnetic moments, the light neutrino bosons, the neutrino millicharges, etc.), for the interpretation of the CEvNS data [46]. Then, whenever measurements of CEvNS events on Mo detectors will come, one may compare our SM predictions with the SM + BSM ones with the experimental data. Furthermore, these results may be discussed in conjunction with the published Ge detector results [10,13,19,22].

In conclusion, the potential of CEvNS may be further exploited by the next generation detectors with improved sensitivity, broader physics goals, and varied target materials. For example, the molybdenum detection medium studied in this work. Thus, future measurements of CEvNS events may significantly improve the precision of the relevant observables such as the neutron root-mean-square (rms) radius of the target Mo-isotopes, and other fundamental Standard Model quantities like the neutrino charge radius and the weak mixing angle originating from the Mo detector.

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DATA AVAILABILITY

Detailed calculational data are available, upon reasonable request, from the author R. Sahu.

APPENDIX

The definitions of the eight multipole operators $\hat{M}_J$, $\hat{L}_J$, $\hat{T}_J^{el}$, $\hat{T}_J^{mag}$, $\hat{M}_J^5$, $\hat{L}_J^5$, $\hat{T}_J^{el5}$, and $\hat{T}_J^{mag5}$, where the subscript 5 refers to the axial vector components of the hadronic current, are as follows:

$$\begin{aligned}\hat{M}_{JM}^{Coul}(q\mathbf{r}) &= F_1^Z \hat{M}_M^J(q\mathbf{r}), & \hat{L}_{JM}(q\mathbf{r}) &= \frac{q_0}{q} \hat{M}_{JM}^{Coul}(q\mathbf{r}), \\ \hat{T}_{JM}^{el}(q\mathbf{r}) &= \frac{q}{M} \left[F_1^Z \Delta_M^{JJ}(q\mathbf{r}) + \frac{1}{2} (F_1^Z + 2MF_2^Z) \Sigma_M^J(q\mathbf{r}) \right], \\ i\hat{T}_{JM}^{mag}(q\mathbf{r}) &= \frac{q}{M} \left[F_1^Z \Delta_M^J(q\mathbf{r}) - \frac{1}{2} (F_1^Z + 2MF_2^Z) \Sigma_M^{JJ}(q\mathbf{r}) \right], \\ i\hat{M}_{JM}^5(q\mathbf{r}) &= \frac{q}{M} \left[F_A^Z \Omega_M^J(q\mathbf{r}) + \frac{1}{2} F_A^Z \Sigma_M^{JJ}(q\mathbf{r}) \right], \\ -i\hat{L}_{JM}^5(q\mathbf{r}) &= F_A^Z \Sigma_M^{JJ}(q\mathbf{r}), & -i\hat{T}_{JM}^{el5}(q\mathbf{r}) &= F_A^Z \Sigma_M^{JJ}(q\mathbf{r}), \\ \hat{T}_{JM}^{mag5}(q\mathbf{r}) &= F_A^Z \Sigma_M^J(q\mathbf{r}).\end{aligned}\tag{A1}$$

where we neglect the pseudoscalar form factor. In Eq. (A1), the first three and last multipole operators have normal parity, $\pi = (-)^J$ while others have abnormal parity, $\pi = (-)^{J+1}$.

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