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Searching for dark matter sterile neutrinos in the laboratory

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If the dark matter of the Universe is made of sterile neutrinos with mass in the keV region, it can be searched for with the help of x-ray satellites. We discuss the prospects of *laboratory* experiments that can be competitive and complementary to space missions. We argue that the detailed study of β decays of tritium and other nuclei with the help of cold target recoil ion momentum spectroscopy can potentially enter into an interesting parameter range and even supersede the current astronomical bounds on the properties of the dark matter sterile neutrinos.

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I. INTRODUCTION

The nature of dark matter (DM) in the Universe is a puzzle. Many different hypothetical particles coming from physics beyond the standard model (SM) were proposed to play the role of a dark matter particle; none of them have been discovered yet. In this paper we will discuss the possibilities for a laboratory search of one of the dark matter candidates: sterile neutrinos with mass in the keV region.

In short, here is the case for a sterile neutrino as a dark matter particle. There are not that many experimental facts in particle physics which cannot be described by the standard model. These are neutrino oscillations (neutrinos of the SM are exactly massless and do not oscillate), dark matter (the SM does not have any stable neutral massive particle), and baryon asymmetry of the Universe (substantial deviations from thermal equilibrium, needed for baryogenesis, are absent for experimentally allowed mass of the Higgs boson; in addition, it is a challenge to use *CP* violation in Cabibbo-Kobayashi-Maskawa mixing of quarks to produce baryon asymmetry in the SM). This calls for an extension of the SM. Perhaps, the most economical one that can describe all these phenomena in a unified way is the νMSM of [1,2]. In this model, three leptonic singlets (other names for them are right-handed, Majorana, or sterile neutrinos) are added, making the structure of quark and lepton sectors of the theory similar, up to the introduction of Majorana mass for the new leptonic states. The Majorana nature of singlet fermions leads to nonzero masses for active neutrinos and, therefore, to neutrino oscillations, solving in this way one of the SM problems. The lightest of these new particles with mass in the keV region can have a lifetime greater than that of the Universe [3] and thus can play a role of (warm) dark matter [4]. The preference for the keV mass scale is coming from the cosmological structure formation arguments related to the missing satellites problem [5,6] and to cuspy DM distributions in cold dark matter cosmologies [7,8]. For other astrophysical applications of keV sterile neutrinos, see [9–13]. The presence of two other heavier fermions with mass in the $\mathcal{O}(1)$ GeV region leads to the generation of baryon asymmetry of the Universe [2] via resonant sterile neutrino oscillations [14] and electroweak sphalerons [15]. These fermions can be searched for in particle physics experiments with high intensity proton beams [16,17].

The only way considered, until now, to detect the dark matter sterile neutrino N is through astrophysical x-ray observations [3,18]. The interaction of N with intermediate vector bosons W and Z and ordinary charged leptons and neutrinos is suppressed by the so-called mixing angle $\theta = m_D/M_M$, where m_D and M_M are, respectively, the Dirac and Majorana masses of sterile neutrinos. The main (but undetectable) decay modes of sterile neutrinos are $N \to 2\nu + \bar{\nu}$, $N \to \nu + 2\bar{\nu}$. In addition, sterile neutrinos has a radiative decay channel $N \to \nu + \gamma$, $N \to \bar{\nu} + \gamma$, producing a narrow line in the x-ray spectrum coming from dark matter halos of different astronomical objects. Recently, a number of constraints on the mixing angle θ became available, coming from the analysis of the x-ray data of Chandra and XMM-Newton satellites [19–27]. It is expected that the best results will come from the analysis of dwarf satellite galaxies in the Milky way halo, as having the largest mass-to-light ratios and weakest x-ray background [21].

In addition to x-ray constraints, there are bounds on the mass and momentum of dark matter sterile neutrinos coming from the analysis of Lyman-α forest clouds and structure formation [28–30]. They depend, however, on the specific mechanism of cosmological production of sterile neutrinos [31]. The most conservative limit on the mass, $M_N > 0.3$ keV, is coming from the analysis of rotational curves of dwarf spheroidal galaxies [32–34] (Tremaine-Gunn bound).

Yet another constraint comes from the requirement that the amount of sterile neutrinos produced in the early universe due to the mixing with ordinary neutrinos must be smaller than the amount of the dark matter observed. In the absence of entropy production due to decays of heavier singlet fermions [31] and assuming that the standard big bang theory is valid at temperatures below a few hundreds of MeV, these bounds are stronger than those coming from x-ray observations for $M_N < 3.5$ keV [35,36]. However, relaxing the above-mentioned assumptions can eliminate these constraints [31,37].

To summarize, the most conservative constraints on the dark matter sterile neutrinos are coming from x-ray observations and from rotational curves of dwarf galaxies; only those will be used in what follows. In Fig. 1 we present the main x-ray bounds taken from [21,25] in the mass range allowed by the Tremaine-Gunn bound, for comparison with the proposed laboratory experiment sensitivity. We stress that these constraints are purely observational and do not depend on any theoretical bias; the only assumption is that the sterile neutrinos constitute 100% of the dark matter in the universe. If only a fraction p of dark matter is in sterile neutrinos, the x-ray constraints are weaker by a factor of p. For p considerably smaller than 1, the Tremaine-Gunn bound is also not applicable.

Imagine now that some day an unidentified narrow line will be found in x-ray observations. Though there are a number of tests that could help to distinguish the line coming from DM decays from the lines associated with atomic transitions in interstellar medium, how can we be sure that the dark matter particle is indeed discovered? Clearly, a laboratory experiment, if possible at all, would play a key role. Current bounds in the interesting mass region were mostly based on a kink search in beta decay, inspired by the possible discovery of 17 keV neutrinos. The present bounds [38] are given in Fig. 1 and are much

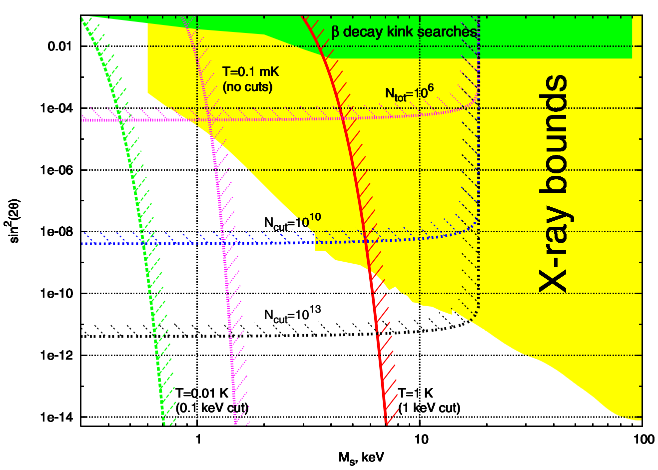

FIG. 1 (color online). Constraints on the mixing angle θ of a sterile neutrino with an active neutrino from x-ray observations of large Magellanic clouds and the Milky Way by XMM-Newton and the Milky Way by HEAO-1 satellites and from kink searches in beta decay. The x-ray bound is in the assumption that sterile neutrinos constitute $p = 100\%$ of the DM; for smaller p the bound is relaxed accordingly. The boundaries of the parameter space accessible to β-decay experiments are shown for various values of the source temperature, kinematic cuts, and collected statistics.

weaker than required to compete with x-ray observations. Figure 1 demonstrates that the search for DM sterile neutrinos in terrestrial experiments is very challenging, as the strength of interaction of DM sterile neutrinos with the matter is roughly θ^2 times weaker than that of ordinary neutrinos. In this paper we analyze different processes where DM sterile neutrinos can be searched for in the laboratory. We argue that the only potential possibility is provided by a precise study of the kinematics of beta decays. We note, in particular, that this kind of study is not impossible in view of a novel momentum space imaging technique [cold target recoil ion momentum spectroscopy (COLTRIMS)], proposed and developed at the end of the 1990s (for reviews see [39,40]).

The paper is organized as follows. In Sec. II we will discuss different reactions in which the DM sterile neutrinos can potentially manifest themselves. In Sec. III we consider the requirements to the β-decay recoil experiments that can enter into the interesting region of a mixing angle for DM sterile neutrinos and provide arguments that they could be potentially feasible. The last section is conclusions.

II. DM STERILE NEUTRINOS IN THE LABORATORY

In the νMSM, interaction of DM sterile neutrino N with fermions of the SM can be derived from the standard 4-fermion weak interaction by replacements $\nu_\alpha \to \nu_\alpha + \theta_\alpha N$, where $\alpha = e, \mu, \tau$, $\theta^2 = \sum \theta_\alpha^2$, and we work in the lowest order in mixing angles θ_α. Thus, sterile neutrinos participate in all reactions the ordinary neutrinos do with a probability suppressed by θ^2. Additionally, the fact that they are Majorana particles, $N = N^c$ (c is the sign of charge conjugation), leads to lepton number nonconservation.

From very general grounds the possible experiments for the search of sterile neutrinos can be divided in three groups.

(i) Sterile neutrinos are *created* and subsequently *detected* in the laboratory. The number of events that can be associated with sterile neutrinos in this case is suppressed by θ^4 in comparison to similar processes with ordinary neutrinos. The smallness of the mixing angle, as required by x-ray observations, makes experiments of this type hopeless. For example, for sterile neutrino mass $m_s = 5$ keV, the suppression in comparison with neutrino reactions is at least of the order of 10^{-19}.

(ii) Sterile neutrinos are created somewhere else in large amounts and then *detected* in the laboratory. The x-ray space experiments are exactly of this type: the number density of sterile neutrinos is fixed by the DM mass density, and the limits on the x-ray flux directly give the limit on θ^2 rather than θ^4 as in the previous case. Another potential possibility is to look for sterile neutrinos coming from the Sun. The flux of sterile neutrinos from, say, pp reactions is

$F_N \sim 6 \times 10^{10} \theta_e^2/\mathrm{cm}^2$ s. The only way to distinguish sterile neutrinos from this source from electronic neutrinos is the kinematics of the reactions $\nu_e n \to pe$ and $Nn \to pe$, which looks hopeless. For higher energy sources, such as ^{8}B neutrinos, the emission of sterile N would imitate the antineutrinos from the Sun due to the reaction $Np \to ne^+$, which is allowed since N is a Majorana particle. However, this process is contaminated by irremovable background from atmospheric antineutrinos. Even if all other sources of background can be eliminated, an experiment like KamLAND would be able to place a limit of the order of $\theta^4 < 3 \times 10^{-7}$, which is weaker than the x-ray limit for all possible sterile neutrino masses obeying the Tremaine-Gunn bound [32–34]. The current KamLAND limit can be extracted from [41] and reads $\theta^4 < 2.8 \times 10^{-4}$. The sterile neutrino can also be emitted in supernovae (SNe) explosions in amounts that could be potentially much larger than $\theta^2 F_\nu$, where F_ν is the total number of active neutrinos coming from SNe. The reason is that the sterile neutrinos interact much weaker than ordinary ν and thus can be emitted from the volume of the star rather than from the neutrino sphere. Using the results of [3], the flux of SNe sterile neutrinos due to $\nu_e - N$ mixing is $F_N \simeq 5 \times 10^3 \theta_e^2 (m_s/\mathrm{keV})^4 F_\nu$. In spite of this enhancement, we do not see any experimental way to distinguish the N and $(\nu, \bar{\nu})$ induced events in the laboratory.

(iii) The process of sterile neutrino *creation* is studied in the laboratory. In this case, one can distinguish between two possibilities. In the first one, we have a reaction which would be exactly forbidden if sterile neutrinos are absent. We were able to find just one process of this type, namely $S \to$ invisible, where S is any scalar boson. Indeed, in the SM the process $S \to \nu\bar{\nu}$ is not allowed due to chirality conservation, and $S \to \nu\nu$ is forbidden by the lepton number conservation. With sterile neutrinos, the process $S \to \nu N$ may take place. However, a simple estimate shows that the branching ratios for these modes for available scalar bosons such as π^0 or K^0 are incredibly small for admitted (by x-ray constraints) mixing angles. So, only one option is left out: the detailed study of kinematics of different β decays.

An obvious possibility would be the main pion decay mode $\pi^- \to \mu\nu$ with creation of sterile neutrino N instead of the active one. This is a two body decay, so the energy muon spectrum is a line with the kinetic energy $(m_\pi - m_\mu)^2/2m_\pi = 4.1$ MeV for decay with active neutrino and $((m_\pi - m_\mu)^2 - m_s^2)/2m_\pi$ for decay with massive sterile neutrino. Thus, for m_s of keV order, one needs the pion beam with energy spread less than 0.01 eV to distinguish the line for sterile neutrino, which seems to be impossible to get with current experimental techniques.

In the case of beta decay there are two distinct possibilities. One is to analyze the electron spectrum only. In this case the admixture of sterile neutrinos leads to the kink in the spectrum at the distance m_s from the end point. However, distinguishing a small kink of the order of θ^2 on top of the electron spectra is very challenging from the point of view of statistically large physical background and nontrivial uncertainties in electron spectrum calculations.

The case of full kinematic reconstruction of beta decay of radioactive nucleus is more promising and will be analyzed in the next section.

III. COLTRIMS AND β-DECAYS

The idea of using beta decay for sterile neutrino detection is quite simple: measuring the full kinematic information for the initial isotope, recoil ion, and electron, one can calculate the neutrino invariant mass on an event by event basis. In an ideal setup of exact measurement of all of these three momenta, such an experiment provides a background-free measurement where a single registered anomalous event will lead to the positive discovery of heavy sterile neutrinos. This idea was already exploited at the time of neutrino discovery and testing of the Fermi theory of β decay [42]. It was also proposed to use full kinematic reconstruction to verify the evidence for the 17 keV neutrino found in the kink searches [43,44] to get rid of possible systematics deforming the beta spectrum. Recently, bounds on sterile neutrino mixing were achieved by full kinematic reconstruction of the ^{38m}K isotope confined in a magneto-optic trap [45] but for a neutrino in the mass range 0.7–3.5 MeV, which is much heavier than considered here. For the 370–670 keV mass range, a similar measurement was performed in electron capture decay of ^{37}Ar [46]. We will discuss below a possible setup for a dedicated experiment for a search of keV scale DM sterile neutrinos.

Let us consider an idealized experiment in which a cloud of β-unstable nuclei, cooled to temperature T, is observed. For example, for ^{3}H the normal beta decay is

$$^3\mathrm{H} \to {}^3\mathrm{H} + e + \bar{\nu}_e,$$

while in the presence of sterile neutrinos the fraction θ^2 of the events (up to the kinematic factor) proceeds as

$$^3\mathrm{H} \to {}^3\mathrm{He} + e + N,$$

where N is a sterile neutrino in the mostly right-handed helicity state. Suppose that it is possible to register the recoil momentum of the daughter ion and of the electron with high enough accuracy. Indeed, existing COLTRIMS experiments are able to measure very small ion recoil [39,40]. They are utilized for investigation of the dynamics of ionization transitions in atoms and molecules. The ion momenta is determined by time-of-flight measurement. A small electric field is applied to the decay region to extract charged ions into the drift region. After the drift region, the ions are detected by a position sensitive detector, which allows one to determine both the direction of the momenta and the time of flight. Characteristic energies of the recoil ion in beta decay is of the order of the recoil momenta measured by existing COLTRIMS in ion-atom collisions. Precisions currently achieved with such apparatus are of the order of 0.2 keV for the ion momentum [39,47–50].

Electron detection is more difficult, as far as the interesting energy range is of the order of 10 keV for ^{3}H decay (or greater for most other isotopes). This is much higher than typical energies obtained in atomic studies. One possible solution would be to use the similar time-of-flight technique as for the recoil ions, but with adding magnetic field parallel to the extraction electric field, thus allowing to collect electrons from a wider polar angle. In existing applications, such a method was used for electrons with energies of only 0.1 keV [51,52]. In [53] a retarding field was added in the electron drift region allowing one to work with electrons of up to 0.5 keV energies. Alternatively, one may try to use electrostatic spectrometers for electron energy measurement, as it was proposed in [43,44]. On the one hand, the latter method allows one to use the electron itself to detect the decay moment for recoil time-of-flight measurement. On the other hand, it is hard to reach high polar angle acceptance with this method, thus losing statistics.

The decay moment needed for the time-of-flight measurement can be tagged by registering the Lyman photon emission of the excited ion or by the electron detection, if electron energy is determined by a dedicated spectrometer. Note that, for the $^3H_2$ case, a Lyman photon is emitted only in about 25% of the events [44], so the photon trigger also induces some statistics loss.

According to [39] it is possible to achieve sensitivity for measuring normal active neutrino masses of 10 eV for each single event; the accuracy needed for the case of sterile neutrinos is considerably less than that as the mass of N is expected to be in the keV region. Moreover, the measurement in the latter case is a relative measurement, which is much simpler than absolute measurement of the peak position required for active neutrino mass determination.

Let us estimate the required precision of momentum measurements and source temperature. Suppose that an initial molecule with mass M decays *at rest* into recoil ion, electron, and neutrino with momenta $\mathbf{p}$, $\mathbf{k}$, and $\mathbf{q} = \mathbf{p} + \mathbf{k}$, respectively. The energy release will be denoted by $\mathcal{Q}$. Then the neutrino mass can be defined from ion and electron momenta as

$$m_\nu^2 = (\mathcal{Q} - E_e - E_p)^2 - (\mathbf{p} + \mathbf{k})^2,$$

where $E_e = \sqrt{m_e^2 + \mathbf{k}^2} - m_e$ and $E_p = \sqrt{M^2 + \mathbf{p}^2} - M$ are electron and recoil ion kinetic energies. It is immediately seen that, for measuring keV neutrino mass, precision of 0.5 keV in momenta measurement would be sufficient. For example, for ^{3}H decay this means 0.5×10^{-2} precision in momentum measurement.

Let us turn now to the question of the temperature of the cloud. Nonzero thermal velocity $\mathbf{v}$ of the initial molecule spoils the measurement. The measured values of the momenta would be $\tilde{\mathbf{p}} = \mathbf{p} + M\mathbf{v}$ and $\tilde{\mathbf{k}} = \mathbf{k} + m_e\mathbf{v}$ so the measured mass m_ν^{eff} is now

$$m_\nu^{\text{eff}^2} \equiv (\mathcal{Q} - \tilde{E}_e)^2 - (\tilde{\mathbf{p}} + \tilde{\mathbf{k}})^2 \simeq m_\nu^2 + M^2\mathbf{v}^2 - 2M\tilde{\mathbf{q}}\mathbf{v},$$

where $\tilde{\mathbf{q}} \equiv \tilde{\mathbf{p}} + \tilde{\mathbf{k}}$ and we have neglected terms of the order of m_e, $\mathcal{Q}$, and $|\mathbf{k}|$ compared to the ion mass M. The average squared thermal velocity is $M^2\langle\mathbf{v}^2\rangle = 3MT$, while the average momenta (in the rest frame) is $|\mathbf{q}| \lesssim \mathcal{Q}$. It is immediately seen that for reasonably low temperatures the last term leads to the dominant error. Assuming an isotropic thermal probability distribution of decaying atoms $P(\mathbf{v}) \propto \exp(-M\mathbf{v}^2/2T)$, it is easy to find (with exponential accuracy) the probability to get nonzero value of $m_\nu^{\text{eff}} \equiv m_s$ from an event originating from a β decay to massless ordinary neutrino,

$$P(m_s) \propto \exp\left(-\frac{m_s^2}{2MT} f^2\left(\frac{|\mathbf{q}|}{m_s}\right)\right), \tag{1}$$

where $f(z) = 1/(\sqrt{1 + z^2} + z)$.

Comparing this background with the number of sterile neutrino events, which is proportional to θ^2, we get that the required temperature must satisfy

$$\frac{m_s^2}{2MT} f^2\left(\frac{|\mathbf{q}|}{m_s}\right) \gtrsim \log(1/\theta^2).$$

Now, if all β decay events are considered, the maximal value of $|\mathbf{q}|$ is $\mathcal{Q}$, and the bound on the temperature is rather rigid and reads approximately

$$T \lesssim \frac{0.7 \times 10^{-3}}{\log(1/\theta^2)} \left(\frac{m_s}{1\ \text{keV}}\right)^4 \left(\frac{6\ \text{GeV}}{M}\right) \left(\frac{18.6\ \text{keV}}{\mathcal{Q}}\right)^2 (1\ \text{K}). \tag{2}$$

However, one can loosen this bound considerably at the cost of the effective source intensity by imposing a kinematic cut on the momenta $(\tilde{\mathbf{q}})^2 \lesssim 3MT$, which leads to $|\mathbf{q}|/m_s \simeq 0$ in (1) and to much higher acceptable temperatures

$$T \lesssim \frac{1}{\log(1/\theta^2)} \left(\frac{m_s}{1\ \text{keV}}\right)^2 \left(\frac{6\ \text{GeV}}{M}\right) (1\ \text{K}). \tag{3}$$

Another bound on the experimental sensitivity to the mixing angle is provided by the requirement that at least one decay with sterile neutrino happens during the observation. The number of signal events is estimated using the differential decay width for massive neutrino $d\Gamma \sim \theta^2 q^2 E_e dq$, where $E_e \simeq \sqrt{2m_e(\mathcal{Q} - \sqrt{q^2 + m_\nu^2})}$ is the electron kinetic energy and $q = |\mathbf{q}|$ is the neutrino momentum. If the kinematic cut on the momenta is much smaller than $\mathcal{Q}$ and m_s, the number of sterile neutrino events can be estimated as

$$N_{\rm events} \simeq \theta^2\sqrt{1-(m_s/Q)}N_{\rm cut},$$

with $N_{\rm cut}$ being the number of active neutrino events satisfying $|\mathbf{q}| < C$, which is related to the total number of decays $N_{\rm tot}$ by

$$\frac{N_{\rm cut}}{N_{\rm tot}} \simeq \frac{35}{16}\left(\frac{C}{Q}\right)^3.$$

For the case with generic cut numerical integration of the differential width can be performed. One should also note that the cut on the momentum suppresses active neutrino events more than sterile neutrino events.

After this general discussion let us turn to specific numbers. As we see, the best beta decay sources should have small mass and small energy release. Then the contribution from the thermal motion of the decaying atoms will be small and created sterile neutrinos will be not too relativistic, thus making it easier to measure its mass. Here the exceptional opportunity is provided with tritium ^{3}H, which is much lighter than all other radioactive elements and has a short enough lifetime of 12.3 years. However, if the neutrino mass is higher than the energy release $Q = 18.591$ keV in tritium decay other isotopes should be used. The requirements on the temperature for the case of the $^3H_2$ molecule can be read directly from (2) and (3) by substituting M by 6 GeV. The bounds (3) and $N_{\rm events} > 10$ (3σ with zero background) for tritium are given in Fig. 1 for several different temperatures and exposures.

One can see that, for $T \sim 0.01$ K and $N_{\rm cut} \sim 10^{13}$, the experiment is better than astrophysical x-ray bounds for all achievable for tritium experiment masses. For a more modest number, like $N_{\rm cut} \sim 10^{10}$ (i.e. year of observation with 1000 decay counts per second, corresponding to a typical recoil ion time of flight in a 1 m long spectrometer), the sensitivity is smaller, but still competes with x-ray experiments for a vast range of masses. Currently [54], one can create supersonic gas jets with particle densities of about 10^{11}–10^{12} cm^{-3} and temperatures of the order of 0.1 K and prepare sources with magneto-optical traps with densities of 10^{10} cm^{-3} and temperature ~0.1 mK. These methods provide 10^6–10^8 beta decays per year for the source size of about 1 mm^3. The line corresponding to the statistics $N_{\rm tot} \sim 10^6$ is shown in Fig. 1 together with the line for 0.1 mK without making any cuts; it indicates that the current magneto-optical trap technologies for low temperature sources allow one to obtain bounds better than existing kink searches and to enter in the interesting parameter region provided that neutrinos make less than 100% of DM.

An important point should be emphasized. All the estimates above have been done with the thermal distribution of the source, which is not necessarily the case. As far as the number of signal events for keV neutrino is expected to be extremely small, the tails of the distribution of thermal velocities are important. Large non-Gaussian tails will not spoil too much the precision of the mass measurement, but will significantly penalize the sensitivity to the mixing angle. This may be a problem for some cooling techniques. For example, the supersonic jet cooling has some non-Gaussian tails in the velocity distribution [39,55].

IV. CONCLUSIONS

In this paper we argued that the detailed study of kinematics of β decays with the help of COLTRIMS may enter into an interesting parameter region for the search of dark matter sterile neutrinos which is complementary to cosmic x-ray missions and indispensable for revealing the nature of DM in the Universe. Even with currently existing technologies, entering in the interesting parameter region is possible for light sterile neutrinos. Extending the study to compete with x-ray bounds for higher masses is a challenging but valuable experimental task.

For a detailed feasibility study of β-decay experiments to search for DM sterile neutrinos, a number of extra points, including existence of possible backgrounds, should be clarified. One obvious background appears from the fact that after the β-decay a fraction (15%) of the $(^3\mathrm{He}^3\mathrm{H})^+$ ions dissociates, leading to an error in the determination of momentum of the detected recoil ions. However, the momentum after dissociation is high [44], so that these events do not look like the heavy neutrino ones. In addition, one should take into account the scattering of ions as a source for possible background. Also, a careful analysis may lead to the choice of another optimal isotope, which has higher decay energy release but is short lived, providing thus larger statistics. A very hard problem is the low density of cold atoms (serving as a source of beta-decays), available at present.

A similar experiment can also be made with isotopes decaying by electron capture. In case one has a single-electron ion of such an isotope, the final state would contain just two particles, and to find the invisible particle mass (ν or N) it is sufficient to measure only the energy of recoil ion, which is a simpler task than the full momentum reconstruction. However, if the initial isotope is not highly ionized, then, generally, one or several Auger electrons are emitted after the decay, carrying away considerable momentum. All these Auger electrons should be detected and taken into account (see [46] for ^{37}Ar case). We leave the comparative study of this type of experiment for future discussion.

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[1] T. Asaka, S. Blanchet, and M. Shaposhnikov, Phys. Lett. B **631**, 151 (2005).
[2] T. Asaka and M. Shaposhnikov, Phys. Lett. B **620**, 17 (2005).
[3] A. D. Dolgov and S. H. Hansen, Astropart. Phys. **16**, 339 (2002).
[4] S. Dodelson and L. M. Widrow, Phys. Rev. Lett. **72**, 17 (1994).
[5] B. Moore *et al.*, Astrophys. J. **524**, L19 (1999).
[6] P. Bode, J. P. Ostriker, and N. Turok, Astrophys. J. **556**, 93 (2001).
[7] T. Goerdt, B. Moore, J. I. Read, J. Stadel, and M. Zemp, Mon. Not. R. Astron. Soc. **368**, 1073 (2006).
[8] G. Gilmore *et al.*, astro-ph/0608528.
[9] A. Kusenko and G. Segre, Phys. Lett. B **396**, 197 (1997).
[10] P. L. Biermann and A. Kusenko, Phys. Rev. Lett. **96**, 091301 (2006).
[11] M. Mapelli, A. Ferrara, and E. Pierpaoli, Mon. Not. R. Astron. Soc. **369**, 1719 (2006).
[12] J. Stasielak, P. L. Biermann, and A. Kusenko, Astrophys. J. **654**, 290 (2007).
[13] J. Hidaka and G. M. Fuller, Phys. Rev. D **74**, 125015 (2006).
[14] E. K. Akhmedov, V. A. Rubakov, and A. Y. Smirnov, Phys. Rev. Lett. **81**, 1359 (1998).
[15] V. A. Kuzmin, V. A. Rubakov, and M. E. Shaposhnikov, Phys. Lett. **155B**, 36 (1985).
[16] M. Shaposhnikov, Nucl. Phys. **B763**, 49 (2007).
[17] D. Gorbunov and M. Shaposhnikov (unpublished).
[18] K. Abazajian, G. M. Fuller, and W. H. Tucker, Astrophys. J. **562**, 593 (2001).
[19] A. Boyarsky, A. Neronov, O. Ruchayskiy, and M. Shaposhnikov, Mon. Not. R. Astron. Soc. **370**, 213 (2006).
[20] A. Boyarsky, A. Neronov, O. Ruchayskiy, and M. Shaposhnikov, Phys. Rev. D **74**, 103506 (2006).
[21] A. Boyarsky, A. Neronov, O. Ruchayskiy, M. Shaposhnikov, and I. Tkachev, Phys. Rev. Lett. **97**, 261302 (2006).
[22] S. Riemer-Sorensen, S. H. Hansen, and K. Pedersen, Astrophys. J. **644**, L33 (2006).
[23] C. R. Watson, J. F. Beacom, H. Yuksel, and T. P. Walker, Phys. Rev. D **74**, 033009 (2006).
[24] S. Riemer-Sorensen, K. Pedersen, S. H. Hansen, and H. Dahle, astro-ph/0610034.
[25] A. Boyarsky, J. Nevalainen, and O. Ruchayskiy, astro-ph/0610961.
[26] A. Boyarsky, O. Ruchayskiy, and M. Markevitch, astro-ph/0611168.
[27] K. N. Abazajian, M. Markevitch, S. M. Koushiappas, and R. C. Hickox, astro-ph/0611144.
[28] S. H. Hansen, J. Lesgourgues, S. Pastor, and J. Silk, Mon. Not. R. Astron. Soc. **333**, 544 (2002).
[29] U. Seljak, A. Makarov, P. McDonald, and H. Trac, Phys. Rev. Lett. **97**, 191303 (2006).
[30] M. Viel, J. Lesgourgues, M. G. Haehnelt, S. Matarrese, and A. Riotto, Phys. Rev. Lett. **97**, 071301 (2006).
[31] T. Asaka, M. Shaposhnikov, and A. Kusenko, Phys. Lett. B **638**, 401 (2006).
[32] S. Tremaine and J. E. Gunn, Phys. Rev. Lett. **42**, 407 (1979).
[33] D. N. C. Lin and S. M. Faber, Astrophys. J. **266**, L21 (1983).
[34] J. J. Dalcanton and C. J. Hogan, Astrophys. J. **561**, 35 (2001).
[35] T. Asaka, M. Laine, and M. Shaposhnikov, J. High Energy Phys. 01 (2007) 091.
[36] T. Asaka, M. Laine, and M. Shaposhnikov, J. High Energy Phys. 06 (2006) 053.
[37] G. Gelmini, S. Palomares-Ruiz, and S. Pascoli, Phys. Rev. Lett. **93**, 081302 (2004).
[38] W. M. Yao *et al.* (Particle Data Group), J. Phys. G **33**, 1 (2006).
[39] R. Dorner *et al.*, Phys. Rep. **330**, 95 (2000).
[40] J. Ullrich *et al.*, J. Phys. B **30**, 2917 (1997).
[41] K. Eguchi *et al.* (KamLAND), Phys. Rev. Lett. **92**, 071301 (2004).
[42] B. Pontecorvo, Rep. Prog. Phys. **11**, 32 (1947).
[43] S. Cook, M. Fink, S. D. Thomas, and H. Wellenstein, Phys. Rev. D **46**, R6 (1992).
[44] G. Finocchiaro and R. E. Shrock, Phys. Rev. D **46**, R888 (1992).
[45] M. Trinczek *et al.*, Phys. Rev. Lett. **90**, 012501 (2003).
[46] M. M. Hindi *et al.*, Phys. Rev. C **58**, 2512 (1998).
[47] R. Dörner *et al.*, Nucl. Instrum. Methods Phys. Res., Sect. B **124**, 225 (1997).
[48] V. Mergel, Ph.D. thesis, Universität Frankfurt, Shaker, Aachen, 1996.
[49] V. Mergel *et al.*, Phys. Rev. Lett. **79**, 387 (1997).
[50] R. Dörner *et al.*, Phys. Rev. A **57**, 1074 (1998).
[51] R. Moshammer, J. Ullrich, M. Unverzagt, W. Schmitt, and B. Schmidt-Bocking, Nucl. Instrum. Methods Phys. Res., Sect. B **108**, 425 (1996).
[52] H. Kollmus, W. Schmitt, R. Moshammer, M. Unverzagt, and J. Ulrich, Nucl. Instrum. Methods Phys. Res., Sect. B **124**, 377 (1997).
[53] T. Jahnke *et al.*, J. Electron Spectrosc. Relat. Phenom. **141**, 229 (2004).
[54] J. Ullrich *et al.*, Rep. Prog. Phys. **66**, 1463 (2003).
[55] D. Miller, *Atomic and Molecular Beam Methods* (Oxford University Press, New York, 1988), Vol. 14.