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Unstable relics as a source of galactic positrons

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Abstract

We calculate the fluxes of 511 keV photons from the galactic bulge caused by positrons produced in the decays of relic particles with masses less than 100 MeV. In particular, we tighten the constraints on sterile neutrinos over a large domain of the mass-mixing angle parameter space, where the resulting photon flux would significantly exceed the experimental data. At the same time, the observed photon fluxes can be easily caused by decaying sterile neutrinos in the mass range $1 < m_{\text{sterile}} \lesssim 50$ MeV with the cosmological abundance typically within $10^{-9} \lesssim \Omega_{\text{sterile}} \lesssim 10^{-5}$, assuming that Ω_{sterile} comes entirely from the conversion of active neutrinos in the early universe. Other candidates for decaying relics such as neutral (pseudo)scalar particles coupled to leptons with the gravitational strength can be compatible with the photon flux, and can constitute the main component of cold dark matter.

1. Introduction

Recent observations of 511 keV photons from the galactic center produced by positron annihilation performed by the SPI spectrometer on the space observatory INTEGRAL [1] confirm previously reported photon fluxes [2] and improve the angular resolution. New data are inconsistent with a single point source, although more data with better angular resolution are needed to exclude or resolve multiple point sources. An apparent delocalization of the source around the galactic center combined with the absence of well-recognized astrophysical sources of positrons that can fill the galactic bulge have prompted speculations about a non-Standard Model origin of galactic positrons.

An interesting explanation was suggested recently in Ref. [3], where galactic positrons were attributed to annihilating dark matter. In order to have photon fluxes at the observable level, one needs rather large number densities of dark matter particles near the center that requires smaller masses than the 1 GeV–1 TeV range preferred by most models of weakly interacting massive particles. Combined with the additional requirement that the produced positrons slow down

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to non-relativistic velocities before annihilation, this puts dark matter particles in a mass range of 1 to 100 MeV [3]. According to the scenario of Ref. [3], the galactic positrons can be effectively trapped by the magnetic field and dissipate their energy by multiple collisions with the gas. The stopping distance is estimated to be comparable or shorter than the "free path" before the annihilation for positrons with energies less than about 100 MeV. An interesting spin-off of this scenario is the possibility to probe the steepness of the dark matter halo profile near the center of the galaxy. Using the power-like parametrization (see, e.g., [4]) of dark matter energy density, $\rho_{\rm DM} \sim r^{-\gamma}$, near the center, the Oxford–Paris group [3] concluded that $0.4 \lesssim \gamma \lesssim 0.8$ is preferred. With more data coming in the future with better angular resolution, and with the better understanding of conventional (dark-matter unrelated) sources of non-relativistic positrons, the flux of 511 keV photons can provide more detailed information on the distribution of dark matter near the center of our galaxy. Even though the explanation of the 511 keV line from the bulge by annihilating dark matter certainly deserves further considerations, one can easily see that this requires very unusual properties of dark matter. Smaller mass and relatively large annihilation cross sections to electron–positron pairs would require either extreme fine tuning in the models of dark matter (see, e.g., [5]) or introduction of new mediating forces that would enable more efficient annihilation [6].

In this Letter, we argue that the decay of quasi-stable relic particles can provide sufficient amount of positrons to explain the SPI/INTEGRAL data. These relics may be the main component of the dark matter, or a sub-dominant fraction of it. From the point of view of properties of dark matter, this would not be as restrictive as the scenario with annihilating dark matter. At the same time, it is clear that if the decaying dark matter can provide sufficient numbers of positrons, the resulting photon angular distribution would imply *steeper* profiles of dark matter near the center than in the scenario with annihilating dark matter.

To substantiate our claims we consider two types of relics that can easily provide a sufficient amount of positrons. Sterile neutrinos in the mass range of few keV – few MeV have been extensively studied in a variety of cosmological and astrophysical settings. Their cosmological abundance is thoroughly investigated by now as a function of their mass m_s, mixing angle θ with active neutrinos [7], as well as their lepton asymmetry and flavor [9,10]. We calculate the decay width of sterile neutrinos and their branching to positrons which is always larger than the branching ratio to gamma quanta if the decay into e^+e^- pair is kinematically allowed, $m_s > 1$ MeV. We demonstrate that in the large part of the $\theta - m_s$ parameter space the resulting flux of photons is *too large* to be compatible with observations, and therefore is excluded. Thus, the observed fluxes of 511 keV photons [2] firmly rule out sterile neutrino dark matter with $m_s > 1$ MeV. On the fringe of this parameter space lies an allowed area compatible with the SPI/INTEGRAL observations. This area is large enough to span a wide range of the parameter space, compatible with sterile neutrino abundances $10^{-9} \lesssim \Omega_{\rm sterile} \lesssim 10^{-5}$. We also show that the constraints on the parameter space of the model from the positron annihilation in the galactic center are complimentary to those coming from the excessive cosmological background of gamma quanta produced in the decay of sterile neutrinos.

Other models of unstable relics can be more flexible, and accommodate dark matter and observable fluxes of photons at the same time. Here we consider models with scalar and/or pseudoscalar particles with MeV range masses and $O(M_{\rm Pl}^{-1})$ couplings to matter. Such scalar particles can be motivated by models with extra space-like dimensions, where radions may have similar properties. We argue the electron–positron channel can be a dominant decay mode for these scalars and that they can be the main component of cold dark matter.

2. Sterile neutrinos as a source of positrons

Sterile neutrinos are the simplest and the most credible extension of the Standard Model. Right-handed neutrinos with a relatively heavy Majorana mass m_R are often used to explain the apparent smallness of the mass range for the left-handed neutrinos hinted by various neutrino experiments. A priori, the right-handed mass m_R can be in a rather large mass range, from the GUT scale all the way down to the keV scale.

In this Letter, we consider a model of a sterile neutrino that is characterized by two main parameters, the mass m_s and the mixing angle θ to the left-handed neutrino:

$$N = \cos\theta N_R + \sin\theta \nu_L,$$
$$\nu = \cos\theta \nu_L - \sin\theta N_R, \tag{2.1}$$

where N_R refers to a pure SM singlet. We do not confine our discussion to the see-saw model and allow θ and m_s to vary independently, without imposing a $m_\nu \sim \theta^2 m_s$ relation. Such models were studied extensively in the mass range of few keV–few MeV [7–11], where sterile neutrinos could constitute entire cold dark matter or at least an appreciable portion of it. Note that the left-handed neutrino in (2.1) need not be a mass eigenstate, and therefore the sterile neutrino N can have interaction with all three charged leptons. Thus, the decay of N to positrons will have an additional parameter V, $|V| < 1$. With this parameter, the interaction of sterile neutrinos with a light active neutrino and electrons takes the following form:

$$\mathcal{L}_{\rm int} = \frac{\sin 2\theta}{2}\left[\frac{eZ_\mu}{2\sin\theta_W\cos\theta_W}\bar N\gamma_\mu\nu + \frac{g_w V W_\mu^+}{\sqrt{2}}\bar N\gamma_\mu e + \frac{g_w V^* W_\mu^-}{\sqrt{2}}\bar e\gamma_\mu N\right] + \text{Higgs terms}. \tag{2.2}$$

Note that the existence of interactions with Higgs boson(s) may be important for the cosmological production of N only if $m_s \gtrsim 1$ GeV, which is outside of the range of 1–100 MeV that we are interested in here. Thus, we neglect Higgs interactions in any further analysis.

In the mass range of 1–100 MeV, the sterile neutrino N can decay into three light neutrinos, light neutrino and a e^+e^- pair, and a neutrino and γ. The decay widths depend on the nature of N, which can be both Dirac or Majorana. For a Majorana particle, the decay width is twice the decay rate for a Dirac particle. This is simply because two decay channels, $N \to \nu\nu\bar\nu$ and $N \to \nu\bar\nu\bar\nu$, are allowed for a Majorana neutrino. For the Dirac sterile neutrino N the decay width into three light active neutrinos (all three flavors) is given by

$$\Gamma_{N\to\nu\nu\bar\nu} = \Gamma_0\frac{\sin^2 2\theta}{4}\left(\frac{m_s}{m_\mu}\right)^5, \tag{2.3}$$

where Γ_0 is the muon decay width, $\Gamma_0 = (192\pi^3)^{-1} \times G_F^2 m_\mu^5$, introduced for convenience.

The decay of N into one light neutrino and the electron–positron pair may occur due to charged and neutral currents. To simplify our formulae, we assume to a few per cent accuracy $1 - 4\sin^2\theta_W \simeq 0$, or $\sin^2\theta_W \simeq 0.25$. In this approximation, and in the limit of $m_e/m_s \ll 1$, the decay width to νe^+e^- summed over all neutrino flavors can be expressed as

$$\Gamma_{N\to\nu e\bar e} = \Gamma_0\frac{\sin^2 2\theta}{4}\left(\frac{m_s}{m_\mu}\right)^5\left[\frac{|V|^2}{2} + \frac{1}{8}\right]. \tag{2.4}$$

More accurate formulae that take into account the phase space suppression for $m_s \sim 2m_e$ and more accurate values of $\sin^2\theta_W$ can be easily derived but not needed for the present discussion. The complete result can be found in Ref. [12]. In Eq. (2.4) we take into account Z-exchange and its interference with W-exchange for an arbitrary V. The existence of both charged and neutral currents is important, because even for $V = 0$, $\Gamma_{N\to\nu e\bar e}$ is different from zero, which ensures a non-zero positron width even for “muon” or “tau” sterile neutrinos.

The decay width into $\nu\gamma$ [12,13] contributes less than 1% to the total decay width,

$$\Gamma_{N\to\nu\gamma} = \Gamma_0\frac{\sin^2 2\theta}{4}\left(\frac{m_s}{m_\mu}\right)^5\frac{27\alpha}{8\pi}, \tag{2.5}$$

so that the total decay width is given by the sum of νe^+e^- and $\nu\nu\bar\nu$ channels, $\Gamma_N \simeq \Gamma_{N\to\nu\nu\bar\nu} + \Gamma_{N\to\nu e\bar e}$. The branching ratio to e^+e^- has a very simple form, and depends only on the size of the mixing matrix element V:

$$\mathrm{Br}_{e^+e^-} = \frac{\Gamma_{N\to\nu e\bar e}}{\Gamma_N} = \frac{1+4|V|^2}{9+4|V|^2}. \tag{2.6}$$

When $|V|^2 = 1$ the branching ratio to e^+e^- reaches the maximum value of $5/13$, while $V = 0$ minimizes it, $\mathrm{Br}^{\rm min}_{e^+e^-} = 1/9$. Thus, the intrinsic model dependence of $\mathrm{Br}_{e^+e^-}$ is a factor of 3.5.

Having determined the total decay width of N, we are ready to approximate the cosmological abundance of sterile neutrinos by combining the exponential decay factor with the power-like pre-exponent reflecting the production of sterile neutrinos in the early universe. Assuming a lepton asymmetry less than 10^{-4}, and using the results of Refs. [7,10], we arrive at

$$\Omega_s h^2 = \Omega_{0s}\{-\Gamma_N \tau\} \sim 0.3\left(\frac{\sin^2 2\theta}{10^{-14}}\right)\left(\frac{m_s}{10\ \text{MeV}}\right)^2 \exp\{-\Gamma_N \tau\}, \tag{2.7}$$

where τ is the age of the universe. In this formula, Ω_{0s} is a fictitious energy density in neglection of neutrino decay, introduced here for future convenience. It coincides with an actual energy density as long as $\Gamma_N \tau$ is small. Note that the pre-exponent in Eq. (2.7) is very approximate. In practice, there will be additional dependence on the lepton asymmetry of the universe and on elements of the V matrix [10]. There is an additional limitation of (2.7), namely it works as long as m_s is smaller than $T \sim 130$ MeV, that is the temperature in the early universe where most of the sterile neutrino production takes place.

In the mass range between 1 and 100 MeV sterile neutrinos behave as cold dark matter. Neutrinos N will be distributed inside a galaxy with the very same profile as the rest of the cold dark matter. Thus, the number density of sterile neutrinos as the function of coordinates can be determined as

$$n_s(r) = \frac{\Omega_s}{\Omega_{\text{DM}}} \frac{\rho_{\text{DM}}(r)}{m_s}, \tag{2.8}$$

where $\Omega_{\text{DM}} = 0.23$ is the "concordance" value for the cold dark matter energy density, and $\rho_{\text{DM}}(r)$ is the cold dark matter energy density in our galaxy.

Using the above equations, we are ready to estimate the emissivity of photons that originate from the galactic center. Assuming that all positrons quickly slow down in matter and annihilate at rest, we estimate the number of emitted 511 keV photons per volume per time,

$$N_{511\gamma}(r) = 2n_s(r)\Gamma_N \text{Br}_{e^+e^-}, \tag{2.9}$$

from which we can estimate the flux that reaches the Solar system,

$$\Phi_{511\gamma}(\theta) = \frac{1}{4\pi} \int_{\text{los}} N_{511\gamma}\big(r(s)\big)\, ds, \tag{2.10}$$

where θ is the angle from the galactic center and the integration is along the line of sight. The total flux is obtained by integrating over the acceptance of the spectrometer. Combining the above equations, one arrives at the following expression:

$$\begin{aligned}\Phi \simeq\ & 8.7 \times 10^{-3}\left(\frac{\sin^2 2\theta}{10^{-15}}\right)^2\left(\frac{m_s}{\text{MeV}}\right)^6 \left(4|V|^2+1\right) \\ & \times \exp\left\{-5\times 10^{-4}\left(\frac{\sin^2 2\theta}{10^{-15}}\right)\left(\frac{m_s}{\text{MeV}}\right)^5 \times \left(4|V|^2+9\right)\right\} \\ & \times \frac{1}{2\pi}\int d\Omega \\ & \times \int_{\text{los}} \left(\frac{\rho_{\text{DM}}}{\text{MeV}\,\text{cm}^{-3}}\right) d\left(\frac{s}{\text{kpc}}\right)\ \text{cm}^{-2}\,\text{s}^{-1}.\end{aligned} \tag{2.11}$$

Integrals involving the dark matter density are commonly performed with a NFW density function [4] which can be simplified to $\rho(r) \propto 1/r^\gamma$ for the inner region of the galaxy in which we are interested. We have done our calculations for different values of γ. However, a better functional form for the NFW density has been developed recently [14] which avoids the singularity and thus does a more careful job of dealing with the radial dependence near the center. This density function is given by

$$\rho_{\text{DM}}(r) = \rho_o \exp\left\{-\frac{2}{\alpha}\left[\left(\frac{r}{r_o}\right)^\alpha - 1\right]\right\}, \tag{2.12}$$

where $r_o = 20h^{-1}$ kpc. To fix an overall normalization ρ_o we take $\rho = 0.3\ \text{GeV}\,\text{cm}^{-3}$ at $r = 8.5$ kpc, the most commonly used value for the dark matter energy density near the Solar system. The values of α which make this expression consistent with previous density distributions for larger values of r range between $\alpha = 0.1$ and $\alpha = 0.2$.

The spectrum of the 511 keV line emission from the galactic center region was obtained by the spectrometer SPI on the INTEGRAL gamma-ray observatory. The data suggest an azimuthally symmetric galactic bulge component with FWHM of $\sim 9^\circ$ with a 2σ uncertainty range covering 6°–18°. The flux of 511 keV photons was measured as $\Phi_{\text{exp}} = 9.9^{+4.7}_{-2.1} \times 10^{-4}\ \text{ph}\,\text{cm}^{-2}\,\text{s}^{-1}$. We have obtained angular distributions and total fluxes expected from the decay of sterile neutrinos using the density function (2.12) in the expression for the flux (2.11) and averaging over the 2° angular resolution of the spectrometer. By comparing our calculated flux to the experimental value

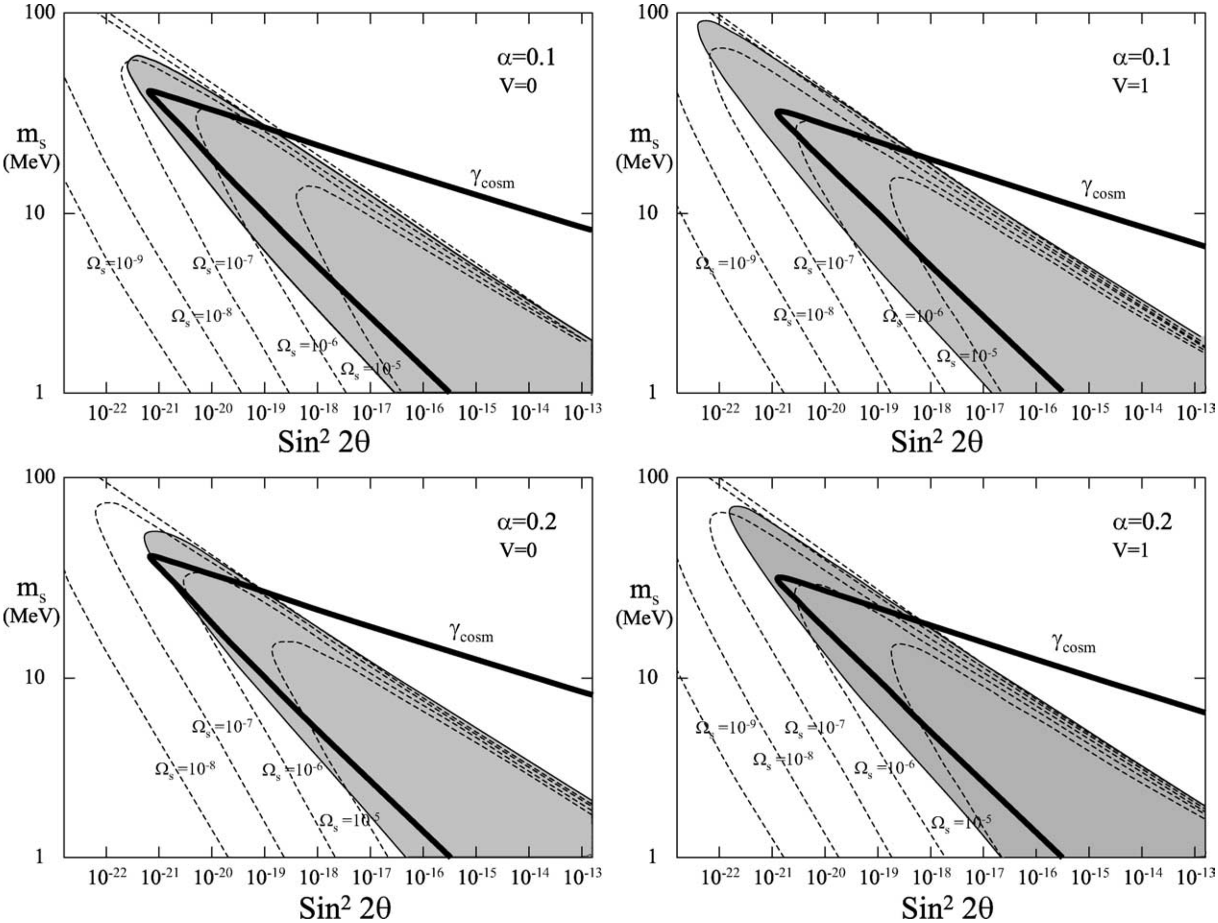

Fig. 1. Mixing angle–sterile neutrino mass parameter space with the contours of sterile neutrino abundance for different halo profiles, $\alpha = 0.1$ and 0.2, and for different mixing angle $|V| = 0$ and 1. The dark shaded region covers the combination of parameters that creates *too large* a flux of 511 keV photons. The boundary of the shaded region gives the same flux of photons, as observed by INTEGRAL/SPI. The area bounded by a thick solid curve is excluded by the diffuse cosmological background of gamma photons.

we find the effective constraints on $\sin^2 2\theta$ and m_s. Fig. 1 shows the results of our calculations for $\alpha = 0.1$ and 0.2, for the extreme values of $|V| = 0, 1$, with the shaded area providing too large a flux compared to the experimental value. Therefore, we can exclude a significant portion of the parameter space in the mass interval 1–100 MeV *regardless* what the source of the galactic positrons really is, conventional astrophysics or exotic scenarios. We also include in these graphs contours of the different values of Ω_s, given by Eq. (2.7). One can readily see that the flux of 511 keV photons is a very powerful tool for probing sterile neutrinos, as even minuscule abundances of sterile neutrinos, $\Omega_s \sim 10^{-9}$, can give observable fluxes.

To produce a model-independent exclusion plot, we assume the conservative value of $\alpha = 0.2$ since this value of α yields the lowest count, and we allow $|V|$ to vary. Then the conservatively excluded region will be the overlap of the two $\alpha = 0.2$ plots for the extreme values of $|V| = 0, 1$, shown in Fig. 2. We can also disconnect the result from a particular functional form of the dark-matter density near the center by fixing its value at $r = 3$ kpc and keeping it constant inside that radius while keeping the same parametrization (2.12) with $\alpha = 0.2$ outside of this region. Clearly, this will yield the most conservative outcome, and we show it also in Fig. 2. Notice that the assumption about the Majorana nature of sterile neutrinos would increase all

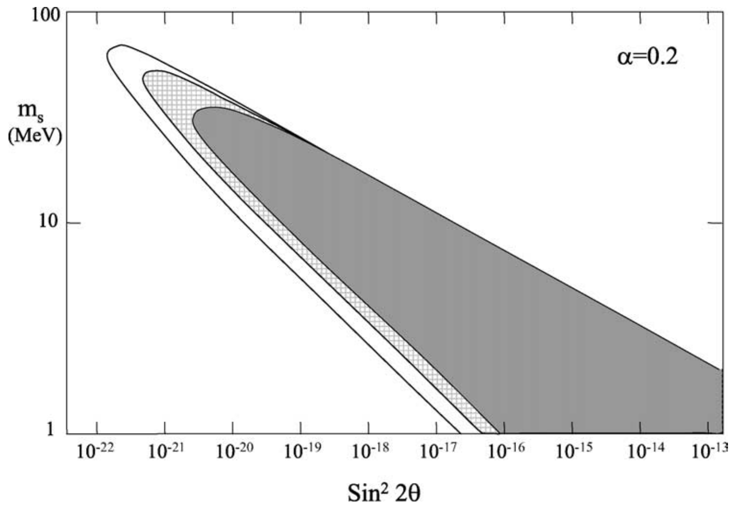

Fig. 2. The exclusion plot for sterile neutrino parameters based on conservative assumptions about the dark matter profile near the center of the galaxy. The unshaded area corresponds to $|V| = 1$ and the hatched area to $|V| = 0$. The solid area is the most conservative model-independent result, corresponding to a dark matter density which is constant inside 3 kpc.

rates by a factor of 2, and thus lead to a larger excluded region.

Returning to Fig. 1, we observe that the combinations of m_s and θ that lie at the boundary of the shaded region give the photon flux that can explain the INTEGRAL/SPI signal. We see that the observed flux can be accommodated with densities in a wide range of sterile neutrino abundances, $10^{-9} < \Omega < 10^{-5}$. Each of these curves contains two branches: for a specified value of the mixing angle θ there are two values of m_s consistent with the photon flux. The branch with lower m_s corresponds to the sterile neutrino lifetimes that are larger than the age of the universe, $\Gamma_N \tau < 1$, and low values of Ω_s in this case are caused by rather inefficient production of sterile neutrinos in the hot universe. Such low Ω_s and mixing angles makes this scenario immune to other possible constraints. At the same time, for the upper branch of the curve, $\Gamma_N \tau > 1$ and small Ω_s can also be caused by the exponential suppression in (2.7). Consequently, for $m_s \sim$ few MeV along this branch the energy density of sterile neutrinos at higher redshifts could be of order one compared to the rest of matter, and thus is subject to other cosmological constraints such as diffuse photon background, (see, e.g., [10,15]).

If we assume that the flux of photons seen by INTEGRAL/SPI comes from the decay of sterile neutrinos, we can use the angular distribution to determine the preferred cold dark matter density profile in the inner galactic region in a similar fashion as in Ref. [3]. Averaging over the angular resolution of the spectrometer and using the common flux normalization at the zero angle result in the angular profiles shown in Fig. 3 for $\alpha = 0.1$ and $\alpha = 0.2$. The distribution for $\alpha = 0.2$ falls well within the 2σ confidence interval, and thus is clearly favored by the data over smaller values of α. The angular distribution that we calculated with the power-like density function $\rho(r) \propto 1/r^\gamma$ for $\gamma = 1$ comes very close to the one for $\alpha = 0.2$.

The photon width $\Gamma_{N\to\nu\gamma}$, although smaller than $\Gamma_{N\to\nu e\bar{e}}$, should lead to the additional photon line at $E_\gamma = m_s/2$. The intensity of $m_s/2$ line, can nevertheless come close to that of 511 keV if m_s is large because of $m_s/(2m_e)$ enhancement factor. Unfortunately, the prediction of photon fluxes at $m_s/2$ relative to 511 keV,

$$\frac{\Phi_{0.5m_s\gamma}}{\Phi_{511\gamma}} = \frac{\Gamma_{N\to\nu\gamma}}{2\Gamma_{N\to\nu e\bar{e}}} = \frac{27\alpha}{\pi(4|V|^2+1)} = \frac{0.031}{4|V|^2+1}, \tag{2.13}$$

varies by a factor of 5 depending on $|V|$.

It is important to compare the limits from positron production in the decays of the sterile neutrinos with the limits derived from a comparison to the cosmo-

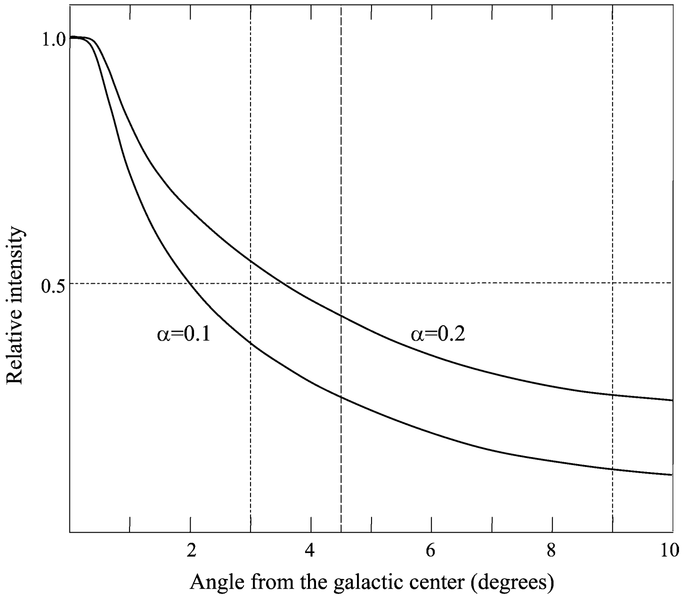

Fig. 3. The angular distribution of 511 keV photons for two different inner halo profiles, $\alpha = 0.1$ and 0.2, averaged over the 2° resolution of the SPI spectrometer. The experimental data were fit with a Gaussian of full width at half maximum at a 9° diameter circle (shown as a vertical dashed line), with a 2σ confidence interval of 6°–18° (shown as vertical dashed-dotted lines). As can be seen, the distribution for $\alpha = 0.2$ falls well within the 2σ confidence interval.

logical diffuse gamma background. Ref. [9] estimates the diffuse photon background produced by the radiative sterile neutrinos with an additional assumption of sterile neutrinos constituting the dominant part of dark matter. As a first approach, we can effectively compare their results by replacing their assumed neutrino abundance of 0.3 by our calculated $10^{-9} \lesssim \Omega_{\text{sterile}} \lesssim 10^{-5}$, which relaxes their constraints on $\sin^2\theta$ by many orders of magnitude. For the diffuse photon background this meant, effectively, a bound on the lifetime $\tau > 4 \times 10^{15}$ years, but as can be seen from our Eq. (3.1) our sensitivity to τ is at the level of 5×10^{17} years assuming the same abundance, 1 MeV neutrinos and a minimum branching ratio to e^+e^- of 0.1. Therefore, it appears that 511 keV photon flux provides more stringent bounds than the diffuse photon background by one to two orders of magnitude. It is nevertheless important to consider this issue in more detail, as the experimental data on photon background used in Ref. [9] are clearly outdated.

The spectral density of the photon flux can be expressed in terms of the sterile neutrino mass m_s, the branching ratio to photons, the neutrino lifetime, the "initial" neutrino abundance Ω_{0s}, the age of the universe τ and its energy density today ρ_0,

$$\frac{d\Phi_\gamma}{dE} = \frac{1}{4\pi}\Omega_{0s}\rho_0 \mathrm{Br}_\gamma \frac{3}{m}\left(\frac{2E}{m}\right)^{1/2} \times \Gamma_N\tau \exp\left\{-\Gamma_N\tau(2E/m)^{3/2}\right\}. \tag{2.14}$$

The maximum in this spectrum corresponds to photon energies

$$E_{\max} = \begin{cases} m/2 & \text{for } 3\Gamma_n\tau < 1, \\ (3\Gamma_n\tau)^{-1}m/2 & \text{for } 3\Gamma_n\tau > 1. \end{cases} \tag{2.15}$$

On the experimental side, [16], we assume a conservative bound on the spectrum of the cosmological diffuse gamma ray background,

$$\left(\frac{d\Phi_\gamma}{dE}\right)_{\exp} \lesssim \frac{10^{-2}}{E^2}\,\mathrm{MeV\,str^{-1}\,cm^{-2}\,s^{-1}}. \tag{2.16}$$

Plugging in the expression for $E_{\max}$ (2.15) to the photon spectrum (2.14), we require $(d\Phi_\gamma/dE)_{\max}$ to be smaller than $(d\Phi_\gamma/dE)_{\exp}$. This results in the following constraint on the parameters of the model,

$$3.5 \times 10^{-7} > \begin{cases} \frac{13}{9+4|V|^2}\Omega_{0s}h^2 x \exp\{-\frac{1}{3}x\} \\ \quad \text{for } x < 1, \\ \frac{13}{9+4|V|^2}\Omega_{0s}h^2 x^{-3/2} \exp\{-\frac{1}{3}x^{-1/2}\} \\ \quad \text{for } x > 1, \end{cases} \tag{2.17}$$

where $x = 3\Gamma_N\tau$ is introduced for concision. This constraint simplifies when $x \ll 1$, in which case we obtain a very powerful limit on the combination of masses and mixing angles,

$$\left(\frac{\sin^2 2\theta}{10^{-19}}\right)^2 \left(\frac{m_s}{1\ \text{MeV}}\right)^5 \lesssim 0.8 \tag{2.18}$$

where V is taken to be one.

Plotting this constraint in Fig. 1, we superimpose it on the limits obtained from the 511 keV line. We observe that the 511 keV line typically provides better sensitivity to parameters of the model when $\Gamma_N\tau < 1$, while the cosmological background of gamma can provide more restrictive bounds if $\Gamma_N\tau > 1$. Unlike the limit from 511 keV photons, the cosmological background constraint is not sensitive to the details of the dark matter distribution in the galactic center.

3. Positron flux from decaying dark matter

In this section, we generalize the analysis of the previous section to an arbitrary decaying relic, without immediately specifying its cosmological origin and keeping its abundance Ω_s as a free parameter, $\Omega_s \leqslant \Omega_{\text{DM}}$. The equality would mean that this relic is the main component of the cold dark matter. The flux of the galactic 511 keV photons, normalized to the measured value, can be written as a function of the mass m_s, the decay width into e^+e^- and Ω_s:

$$\frac{\Phi(m_s, \Gamma_{e^+e^-}, \Omega_s)}{\Phi_{\text{exp}}} \simeq \left(5 \times 10^{17}\ \text{yr} \times \Gamma_{e^+e^-}\right) \frac{10\ \text{MeV}}{m_s} \frac{\Omega_s}{0.23}. \tag{3.1}$$

In this equation, $\alpha = 0.2$ was used. We observe that even if dark matter particles have much longer lifetimes than the age of the universe, they are still capable of providing significant numbers of positrons.

Sterile neutrinos in the mass range of 1–50 MeV can account for the galactic positrons but cannot constitute any appreciable fraction of the dark matter if they are produced thermally from active neutrino species under the assumption of a small lepton asymmetry. To this end, it is interesting to relax the assumptions that led to Eq. (2.7) and using Eqs. (3.1) and (2.4) estimate the resulting photon flux,

$$\frac{\Phi(m_s, \theta, \Omega_s)}{\Phi_{\text{exp}}} \simeq 1.7\left(1 + 4|V|^2\right) \frac{\sin^2 2\theta}{10^{-24}} \left(\frac{m_s}{10\ \text{MeV}}\right)^4 \frac{\Omega_s}{0.23}. \tag{3.2}$$

Combining the analysis of Ref. [10] at $|V| = 1$ and the lepton asymmetry $L = 0.1$ with Eq. (3.2), we find that all points on the $\theta - m_s$ parameter space that correspond to $\Omega_s \geqslant 10^{-3}$ are firmly ruled out by the excessive photon flux. Therefore, even for an arbitrary lepton asymmetry, sterile neutrinos with $m_s > 1$ MeV cannot constitute a significant fraction of dark matter. The combination of both features, the sterile neutrino dark matter ($m_s > 1$ MeV) and acceptable positron numbers from its decay, may only come at the price of abandoning the conventional mechanism of thermal transfer of active neutrinos into sterile, and assuming some initial sterile neutrino energy density due to either non-thermal production after inflation and/or decay of some heavy non-SM particles into N before the electroweak epoch.

Another example of an unstable relic that could produce positrons is an extremely weakly interacting scalar or pseudoscalar particle ϕ in the same MeV mass range. Quite generically, we can write its interaction with operators composed from the Standard Model fields as

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{1}{2}m_\phi^2\phi^2 + \frac{\phi}{M_*}\sum_i c_i O_i^{(4)}, \tag{3.3}$$

where the leading order dimension 4 operators that induce the decay of ϕ are included. In models with extra space-like dimensions, a scalar particle called radion would couple to the trace of the stress-energy tensor of visible matter. If ϕ represents a radion, a scalar mode corresponding to the fluctuation of the volume of extra dimension, the coupling constant M_* can vary within a large range, from M_{pl} or even larger scale down to M_W with the latter being the case in the Randall–Sundrum model [17].

For $1 < m_\phi \lesssim 100$ MeV, the leading interaction with fermions is given by $\phi M_*^{-1} m_e \bar{e} e$ for scalars and $\phi M_*^{-1} m_e \bar{e} i\gamma_5 e$ for pseudoscalars. In both cases, the width to e^+e^- is given by

$$\Gamma_{\phi\to e^+e^-} = \frac{m_e^2 m_\phi}{8\pi M_*^2} \simeq \left(2\times 10^{17}\ \text{yr}\right)^{-1} \frac{m_\phi}{1\ \text{MeV}} \left(\frac{10^{19}\ \text{GeV}}{M_*}\right)^2, \tag{3.4}$$

where the threshold effects are neglected, $2m_e < m_\phi$. Whether Eq. (3.4) provides the dominant contribution to the total decay width depends on whether or not the anomalous interaction with photons, $\phi F_{\mu\nu} F^{\mu\nu}$ or $\phi F_{\mu\nu} \tilde{F}^{\mu\nu}$, is generated at one loop. This depends on the UV completion of the *effective* Lagrangian (3.3), and both outcomes, zero or non-zero couplings to photons, are possible (for the scalar case see, e.g., [18] and references therein). In particular, the one-loop anomalous coupling of scalars and photons is not generated in the models where the interaction term of the scalar ϕ with matter is obtained from a conformal rescaling of Brans–Dicke theory with the mass term for ϕ. If $\phi F_{\mu\nu}^2$ arises at two loops or not at all, then the total decay width is dominated by $\Gamma_{\phi\to e^+e^-}$. For the pseudoscalar case, one could also avoid couplings to photons by postulating a derivative interaction with the SM fields, i.e., $\partial_\mu \phi \bar{\psi} \gamma_\mu \gamma_5 \psi$. Taking the point of view that e^+e^- represents the main decay channel, we combine Eqs. (3.1) and (3.4) to provide an estimate of the photon flux from the Galactic bulge,

$$\frac{\Phi(M_*, \Omega_\phi)}{\Phi_{\text{exp}}} \simeq 25 \left(\frac{10^{19}\ \text{GeV}}{M_*}\right)^2 \frac{\Omega_\phi}{0.23}, \tag{3.5}$$

which turns out to be independent of m_ϕ. It is remarkable that M_* as large as the Planck scale,

$$M_* \sim 5\times 10^{19}\ \text{GeV} \sqrt{\Omega_\phi/0.23}, \tag{3.6}$$

can be probed this way for maximal Ω_ϕ!

Is it realistic to have $\Omega_\phi \sim 0.23$? There are many possible scenarios of non-thermal production of scalars and pseudoscalars. In fact, the scalar field energy density can be over-produced easily leading to the closure of the universe (moduli problem). This is, of course, very much model-dependent, and can be easily avoided in some scenarios. Thermal production can also be quite efficient if in addition to very weak decay terms the scalar ϕ has appreciable bilinear couplings to the Standard Model fields. For example, the Lagrangian [5,19,20],

$$\mathcal{L}_\phi = \frac{1}{2}\partial_\mu \phi \partial^\mu \phi - \frac{1}{2}\left(m_0^2 \phi^2 - \lambda \phi^2 H^\dagger H\right) + \frac{\phi}{M_*} m_e \bar{e} e, \tag{3.7}$$

where H is the Standard Model Higgs doublet, has the potential of providing a dark matter candidate ϕ. To ensure the right abundance of ϕ one can choose λ to be large, in which case ϕ is thermalized above m_ϕ and depletes its abundance via annihilation at the freeze-out. In this case however, there is a direct experimental bound on m_ϕ that requires it to be heavier than 400 MeV [21], thus making it unsuitable for explaining the positron flux. Another possibility for obtaining a cosmologically required value of Ω_ϕ is to assume a small value of λ, initial absence of ϕ scalars, and their thermal production at $T \sim M_W$ [20]. This option is indeed viable, and $\lambda \sim 10^{-7}$ would lead to $\Omega_\phi \sim 0.2$. Thus, we see that the scalar field model (3.7) can account for the existence of cold dark matter and lead to the observable fluxes if condition (3.6) is satisfied. In the original formulation of this model without the decay term the stability of the scalar was ascribed to either Z_2 symmetry [5] or global $U(1)$ symmetry if ϕ is a complex scalar [19]. The appearance of the decay terms $\sim M_{\text{Pl}}^{-1} \phi m_e \bar{e} e$ might be speculated to arise due to the violation of these symmetries in string theory/quantum gravity at the Planck scale.

4. Discussion and conclusions

We have shown that the decaying relics that may or may not constitute the main component of dark matter can easily account for the flux of galactic photons observed by the INTEGRAL/SPI instrument [1]. The lifetimes of these relics of order 10^{17} yr, i.e., much larger than the age of the universe, can be probed this way. This is also longer than the lifetime of unstable $\sim$ MeV-scale relics that can be probed with diffuse cosmic gamma-ray background [9]. This sensitivity motivates further theoretical studies of the flux of 511 keV photons from the bulge.

We have studied in detail the model of sterile neutrinos produced in the early universe by conversion from active neutrino species. We have shown that if the decay to e^+e^- pairs is kinematically allowed, the branching ratio in this mode is never smaller than 10% even if the sterile neutrinos have no charged currents with electrons ($V = 0$). Using previous results of abundance calculation [7,10] and the total decay width of the sterile neutrino and its branching to positrons obtained in this work, we were able to calculate the flux of the 511 keV photons as a function of the mixing angle θ with active species and the sterile neutrino mass m_s. Over a large domain of the parameter space for $1 < m_s \lesssim 50$ MeV this flux turns out to be significantly larger than the observed value. This enables us to exclude a significant portion of the parameter space of sterile neutrinos regardless of whether the source of the galactic positrons is astrophysical or it is due to exotic physics. In particular, we are able to rule out any cosmologically significant quantities of sterile neutrinos, $\Omega_s < 10^{-3}$ with masses larger than 1 MeV.

The boundary of the excluded region corresponds to fluxes that are comparable to the observed value. If indeed the experimental results can be interpreted as the consequence of positrons from the decaying sterile neutrinos, the photon flux points to a sterile neutrino abundance in the range of $10^{-9} \lesssim \Omega_{\mathrm{sterile}} \lesssim 10^{-5}$. It is also interesting to note that the upper value for the mass of sterile neutrinos (~ 50 MeV) compatible with the observed flux of photons comes naturally into this analysis, and not as an external requirement as in the case of annihilating dark matter [3]. The radial distribution of the decaying relics inside the galactic halo should trace the distribution of the dark matter. Thus, if indeed future data support the link of galactic positrons to exotic physics (e.g., by firmly establishing their diffuse origin as opposed to point-like sources), the two scenarios, decaying relics and annihilating dark matter, can be distinguished by the angular distribution of the signal. The signal from decaying relics is proportional to the number density of the dark matter particles $n(r)$ while the annihilating dark matter has $n^2(r)$ dependence. Thus, the decaying relics would suggest steeper Galaxy profiles than the scenario with the annihilation of dark matter. In both scenarios, decaying relics and annihilating dark matter, one would generally expect the additional sharp line due to the annihilation or decay to photons. In the case of decaying sterile neutrinos, the flux of photons with energies of $0.5m_s$ is predicted to be at the level of few percent of the flux of 511 keV neutrinos. Finally, we would like to comment that the scenario with annihilating dark matter can be searched for experimentally using the existing data on e^+e^- colliders. Since the results of Ref. [3] suggest an appreciable low-energy dark matter annihilation cross section into e^+e^-, one would naturally look for the missing energy signal (plus a hard photon) in the e^+e^- collisions at higher energies where the SM model cross sections become smaller and while the cross section of the MeV-mass dark matter production would become larger.

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References

[1] P. Jean, et al., Astron. Astrophys. 407 (2003) L55; J. Knodlseder, et al., astro-ph/0309442.
[2] M. Leventhal, et al., Astrophys. J. 25 (1993) 405; D.M. Smith, et al., Astrophys. J. 414 (1993) 165; M.J. Harris, et al., Astrophys. J. 55 (1998) 501.
[3] C. Boehm, D. Hooper, J. Silk, M. Casse, astro-ph/0309686.
[4] J.F. Navarro, C.S. Frenk, S.D.M. White, Astrophys. J. 462 (1996) 563.
[5] C.P. Burgess, M. Pospelov, T. ter Veldhuis, Nucl. Phys. B 619 (2001) 709.
[6] C. Boehm, P. Fayet, J. Silk, hep-ph/0311143.
[7] S. Dodelson, L.M. Widrow, Phys. Rev. Lett. 72 (1994) 17.
[8] A.D. Dolgov, S.H. Hansen, G. Raffelt, D.V. Semikoz, Nucl. Phys. B 590 (2000) 562.
[9] A.D. Dolgov, S.H. Hansen, Astropart. Phys. 16 (2001) 339.
[10] K. Abazajian, G.M. Fuller, M. Patel, Phys. Rev. D 64 (2001) 023501.
[11] G.M. Fuller, A. Kusenko, I. Mocioiu, S. Pascoli, Phys. Rev. D 68 (2003) 103002.
[12] V.D. Barger, R.J.N. Phillips, S. Sarkar, Phys. Lett. B 352 (1995) 365; V.D. Barger, R.J.N. Phillips, S. Sarkar, Phys. Lett. B 356 (1995) 617, Erratum.
[13] P.B. Pal, L. Wolfenstein, Phys. Rev. D 25 (1982) 766.
[14] J.F. Navarro, et al., astro-ph/0311231.

[15] A.D. Dolgov, Y.B. Zeldovich, Rev. Mod. Phys. 53 (1981) 1.
[16] EGRET Collaboration, P. Sreekumar, et al., Astrophys. J. 494 (1998) 523.
[17] L. Randall, R. Sundrum, Phys. Rev. Lett. 83 (1999) 3370.
[18] K.A. Olive, M. Pospelov, Phys. Rev. D 65 (2002) 085044.
[19] J. McDonald, Phys. Rev. D 50 (1994) 3637.
[20] J. McDonald, Phys. Rev. Lett. 88 (2002) 091304.
[21] C. Bird, P. Jackson, R. Kowalewski, M. Pospelov, hep-ph/0401195, Phys. Rev. Lett., in press.