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Hadronically decaying heavy dark matter and high–energy neutrino limits

M. Yu. Kuznetsov*

Institute for Nuclear Research of the Russian Academy of Sciences,

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Abstract

We consider dark matter consisting of long–living particles with masses $10^7 \lesssim M \lesssim 10^{16}$ GeV decaying through hadronic channel as a source of high energy neutrinos. Using recent data on high energy neutrino from IceCube and Pierre Auger experiments we derive the upper-limits on neutrino flux from dark matter decay and constraints on dark matter parameter space. The constraints derived are weaker that these obtained for the same dark matter models using the high energy gamma-ray limits.

Keywords: heavy dark matter, neutrino.

1 Introduction

The idea that the dark matter consists of heavy long–lived particles arose in the context of inflationary cosmology where there are several mechanisms of production of such particles [1, 2, 5–13] (although the heavy dark matter was proposed in other context, see e.g. Refs. [3, 4]). It was also realised that these particles can be sources of ultra high energy cosmic rays (UHECRs) that evade the GZK cutoff [1, 2]. Although the absence of the GZK cutoff is not confirmed by modern cosmic–ray experiments [35, 36] the heavy dark matter is still being considered as a possible source of high energy cosmic rays, in particular photons and neutrino. The special features of dark matter decay fluxes, such as anisotropy and peculiar relation between photon and neutrino components allows one to distinguish between dark matter and astrophysical scenarios of cosmic ray origin.

There are several sources of constraints for heavy dark–matter parameters: the mass M_X is subject to cosmological constraints [8, 9, 14–17], for a given mass the lifetime τ of the dark–matter particles can be effectively constrained using the observed fluxes of various particles or the limits on these fluxes. For example, in Ref. [37] the constraints have been put using the shape of charged cosmic–ray spectra. However, with the modern

*mkuzn@inr.ac.ru

cosmic ray data this method bounds τ not so good as gamma–ray and neutrino flux limits. In this study we use neutrino limits to constraint the dark matter lifetime, conservatively considering the mass range $10^7 \lesssim M_X \lesssim 10^{16}$ GeV which corresponds to the energy range accessible in modern high energy neutrino experiments. The case of absolutely stable heavy dark matter particles is not so interesting from the experimental point of view — its annihilation cross section is bound by unitarity [39]: $\sigma_{ann.} \sim 1/M_X^2$, making their indirect detection impossible for the today experiments.

The detection of high energy neutrino events by IceCube experiment [19, 20] has attracted significant attention. There were many works interpreting these events as an astrophysical neutrino signal [31–33] as well as the dark matter decay signal [24–27]. At the same time, the constraints on various models of neutrino origin have been placed [28]. There were also pre-IceCube studies where neutrino limits were employed to constrain heavy dark matter parameters [29, 30].

This study is mainly inspired by the publication of the new refined sample of IceCube high energy neutrino data along with the updated exposure of this experiment [18]. In that analysis, stringent cuts were employed to eliminate the atmospheric neutrino background. The resulting data set contains only two events with PeV order energy, both consistent with the astrophysical neutrino Monte–Carlo. This fact together with the non-observation of higher energy events allows the IceCube collaboration to place limits on the astrophysical neutrino flux and to constrain several models of astrophysical neutrino origin. We will use the same data sample to place limits on dark matter decay neutrino flux and to constrain parameters of decaying dark matter. For comparison we also derive the constrains using Pierre Auger Observatory data [22] reporting non-detection of neutrino with energies $E_\nu \gtrsim 10^{17}$ eV. This study complements our previous research of constraining heavy decaying dark matter with high energy gamma–ray limits [23].

2 Neutrino flux from dark matter decay

In this study we consider dark matter consisting of heavy particles X decaying through the hadronic channel, i.e. via the chain of processes

$$X \to q\bar{q}\ (gg) \to \text{hadrons} \to \nu_i\ (\bar{\nu}_i) \tag{1}$$

assuming hadronisation and decay of unstable hadrons. The method of calculation of the final state stable particles spectra for this type of decay was reviewed in our previous work [23] and mainly follows the Refs. [38,40]. In this study we consider the neutrino flux. The main contribution to the flux comes from the decay of charged pions via processes

$$\pi \to \mu\nu_\mu, \quad \mu \to e\nu_\mu\nu_e\,. \tag{2}$$

There are also contributions from kaons as well as from charmed mesons but they are an order of magnitude smaller that the contribution of pions. Moreover, the uncertainty of the pion flux which is dominated by the uncertainty of the pion fragmentation functions on the initial energy scale is of the same order as the contributions of other mesons to the neutrino flux [41]. Therefore we assume that the neutrino production is saturated by the pion decays. Using the results of Ref. [53] we are also make sure of negligibility of electro–weak corrections to the decay spectra.

We consider the pion spectrum $\frac{dN_\pi}{dx}$ where $x = \frac{2E_\pi}{M_X}$, to obtain it on M_X scale we evolve pion fragmentation functions from some initial scale using DGLAP equations [42, 43]:

$$\frac{\partial D_i^\pi(x,s)}{\partial \ln s} = \sum_j \frac{\alpha_s(s)}{2\pi} P_{ij}(x, \alpha_s(s)) \otimes D_j^\pi(x,s)\,, \tag{3}$$

where $D_i^\pi(x,s)$ is the fragmentation function of pion from the parton i, s is the factorization scale, $\otimes$ denotes the convolution $f(x)\otimes g(x) \equiv \int_x^1 dz/z f(z) g(x/z) = \int_x^1 dz/z f(x/z) g(z)$ and $P_{ij}(x,s)$ is the splitting function for the parton branching $i \to j$. We use the same assumptions about DGLAP evolution and fragmentation functions as in our previous work [23], namely we assume all quark flavors are coupled to gluon similarly and consider the mixing of gluon fragmentation function with the quark singlet fragmentation function. As in our previous work we use the code of Ref. [38] to solve DGLAP equations numerically in the leading order of $\alpha(s)$. We take the initial fragmentation functions parametrized on the scale of 1 GeV from Ref. [41] and extrapolate them in the interval $10^{-5} \leq x \leq 1$.

The neutrino spectrum from pions decay is given by

$$\frac{dN_{\pi\to\nu_\mu}}{dx} = 2\,R \int_{xR}^{1} \frac{dy}{y}\,\frac{dN_\pi}{dy}\,, \tag{4}$$

while the neutrino spectrum from the decay of secondary muons is

$$\frac{dN_{\mu\to\nu_i}}{dx} = 2 \int_x^1 \frac{dz}{z} f_{\nu_i}\left(\frac{y}{z}\right)\,\frac{dN_\pi}{dz}\,; \tag{5}$$

where $r = (m_\mu/m_\pi)^2 \simeq 0.573$, $R = \frac{1}{1-r}$ and the functions $f_{\nu_i}(x)$ are taken from Ref. [44]:

$$f_{\nu_i}(x) = g_{\nu_i}(x)\,\Theta(x-r) + (h_{\nu_i}^{(1)}(x) + h_{\nu_i}^{(2)}(x))\,\Theta(r-x)\,,$$

$$g_{\nu_\mu}(x) = \frac{3-2r}{9(1-r)^2}\,\left(9x^2 - 6\ln x - 4x^3 - 5\right),$$

$$h_{\nu_\mu}^{(1)}(x) = \frac{3-2r}{9(1-r)^2}\,\left(9r^2 - 6\ln r - 4r^3 - 5\right),$$

$$h_{\nu_\mu}^{(2)}(x) = \frac{(1+2r)(r-x)}{9r^2}\,\left[9(r+x) - 4(r^2+rx+x^2)\right]\,,$$

$$g_{\nu_e}(x) = \frac{2}{3(1-r)^2}\,\left[(1-x)\,\left(6(1-x)^2 + r(5+5x-4x^2)\right) + 6r\ln x\right],$$

$$h_{\nu_e}^{(1)}(x) = \frac{2}{3(1-r)^2}\,\left[(1-r)\,\left(6 - 7r + 11r^2 - 4r^3\right) + 6r\ln r\right],$$

$$h_{\nu_e}^{(2)}(x) = \frac{2(r-x)}{3r^2}\,\left(7r^2 - 4r^3 + 7xr - 4xr^2 - 2x^2 - 4x^2 r\right)\,.$$

The examples of neutrino spectra from the decay of X particles with different masses are shown in Fig. 1.

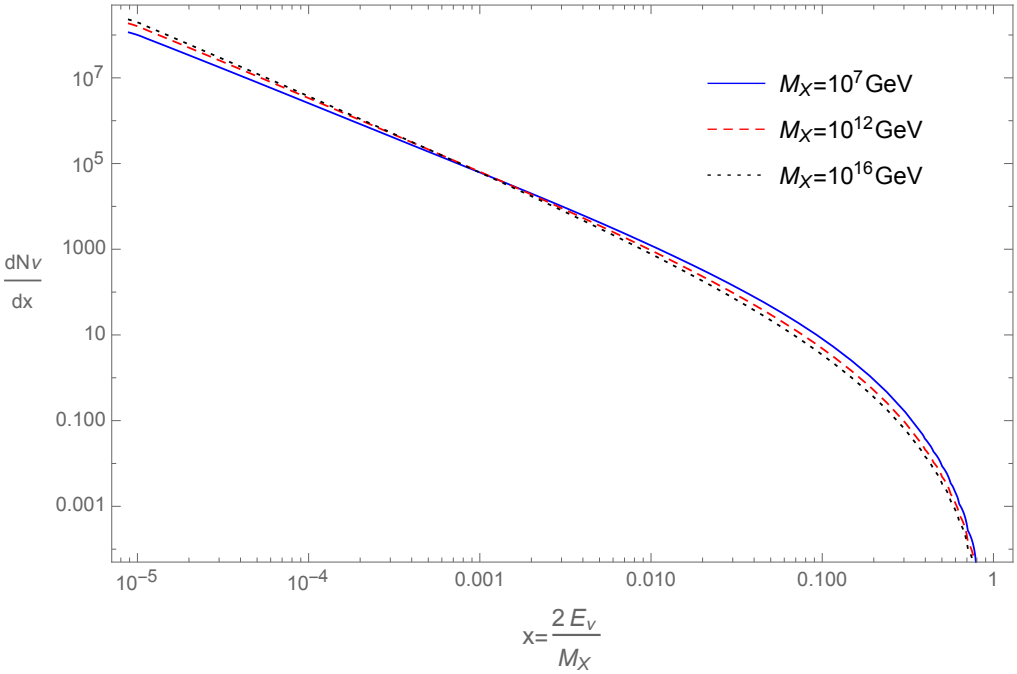

Figure 1: The total spectra of neutrino and antineutrino from X particle decay for three different values of M_X.

Neutrinos propagate cosmological distances unattenuated. The resulting flux reaching the Earth consists of the galactic and extragalactic parts. The initial flavor composition of the pion decay products is modified by the neutrino oscillations during the propagation. We assume the flux reaching the Earth is completely mixed, i.e. the flavor ratio $\nu_e : \nu_\mu : \nu_\tau$ is $1 : 1 : 1$. We also assume that neutrinos are radiated isotropically in the decay of an X particle. For the galactic neutrino flux calculation, we use the Navarro-Frenk-White dark matter distribution [45,46] with the parametrization for the Milky Way from Ref. [53]. Being strongly anisotropic, the galactic signal has to be convolved with the particular experiment's exposure to obtain the perceived flux (see next section). Contrary, the extragalactic flux is isotropic and takes into account the cosmological redshifting of neutrinos

$$\frac{dN_\nu^{\rm EG}}{dE_\nu}(E_\nu) = \frac{1}{4\pi M_X \tau} \int_0^\infty \frac{\rho_0\, c/H_0}{\sqrt{\Omega_m(1+z)^3 + (1-\Omega_m)}} \frac{dN_\nu}{dE_\nu}(E'_\nu) \tag{6}$$

where $c/H_0 = 1.37 \cdot 10^{28}$ cm is the Hubble length, $\rho_0 = 1.15 \cdot 10^{-6}$ Gev/cm^3 is the average cosmological dark matter density for today, $\Omega_m = 0.27$ and the injected spectrum $\frac{dN_\nu}{dE_\nu}$ is taken as a function of neutrino energy at redshift z: $E'_\nu = E_\nu(1+z)$.

3 Analysis and discussion

The method of constraining the dark–matter parameters with neutrino limits slightly differs from that using with the gamma–ray limits. The exposure of neutrino observatories depends on the neutrino energy, therefore flux limits depend on neutrino spectrum. Below we briefly describe the method. The quantity we need to compare with the observation

is the total number of neutrino events that would be detected in the given experiment assuming the given neutrino spectrum. The method of calculation of this quantity was given in Ref. [34]. Below all the quantities are related to neutrino so we are omitting the index ν. For the galactic neutrino flux we have

$$N_{\rm G} = \frac{1}{4\pi M_X \tau} \int\limits_{\Delta E} \int\limits_{V} \rho\left[R(r,\delta,\alpha)\right]\, \varepsilon(E,\delta,\alpha)\, \frac{dN}{dE}(E)\, \cos(\delta)\, dr\, d\delta\, d\alpha\, dE\ ; \tag{7}$$

where $\rho[R]$ is a dark matter density as a function of distance from galactic center R, r is distance from Earth, ε is the exposure of the given observatory as a function of the neutrino energy E and equatorial coordinates $\{\delta, \alpha\}$. The integration takes over all volume of the dark–matter halo ($R < 260$ kpc) and over the entire range ΔE of the neutrino energies accessible for a given observatory. In practice, the exposure is given for several bands of zenith angle, averaged over each band. For IceCube we adopt the exposure as a function of declination (which uniquely translates to zenith angle in the case of IceCube) and energy as it is given in Ref. [21] and normalize it to the actual IceCube exposure of Ref. [18]. For Pierre Auger we use the exposure given in Ref. [22] together with the formula of the effective exposure of extensive air shower observatory [51, 52]:

$$\omega(a_0, \delta, \theta_{\rm max}) \sim (\cos a_0\, \cos\delta\, \sin\alpha_m + \alpha_m \sin a_0\, \sin\delta), \tag{8}$$

where a_0 is the latitude of the given observatory, $\theta_{\rm max}$ is the maximal zenith angle accessible for fully efficient observations in this observatory and α_m is given by

$$\alpha_m = \begin{cases} 0 & ; \xi > 1, \\ \pi & ; \xi < -1, \\ \arccos\xi & ; -1 < \xi < 1\,; \end{cases} \tag{9}$$

$$\xi = \frac{(\cos\theta_{\rm max} - \sin a_0\, \sin\delta)}{\cos a_0\, \cos\delta}. \tag{10}$$

The number of events from the extragalactic flux is

$$N_{\rm EG} = \int\limits_{\Delta E} \varepsilon(E)\, \frac{dN_{\rm EG}}{dE}(E)\, dE\ ; \tag{11}$$

where the exposure $\varepsilon(E)$ is integrated over the celestial sphere. Thus the total number of events predicted by the theory is

$$N_{\rm th} = N_{\rm G} + N_{\rm EG}\,. \tag{12}$$

The example of $N_{\rm G}$ and $N_{\rm EG}$ for fixed τ and various masses M_X is shown in Fig. 2. We can see that the neutrino signal from our Galaxy is generally two times higher than the extragalactic one.

The parameter we want to constrain for each dark matter mass M_X is its lifetime τ. The standard technique of Ref. [54] implies that we vary τ until the predicted number of events $N_{\rm th}$ reaches from below the number $N_{\rm limit}$ specified for a given number of observed events $N_{\rm obs}$, number of background events $N_{\rm bg}$ and given confidence level. We may

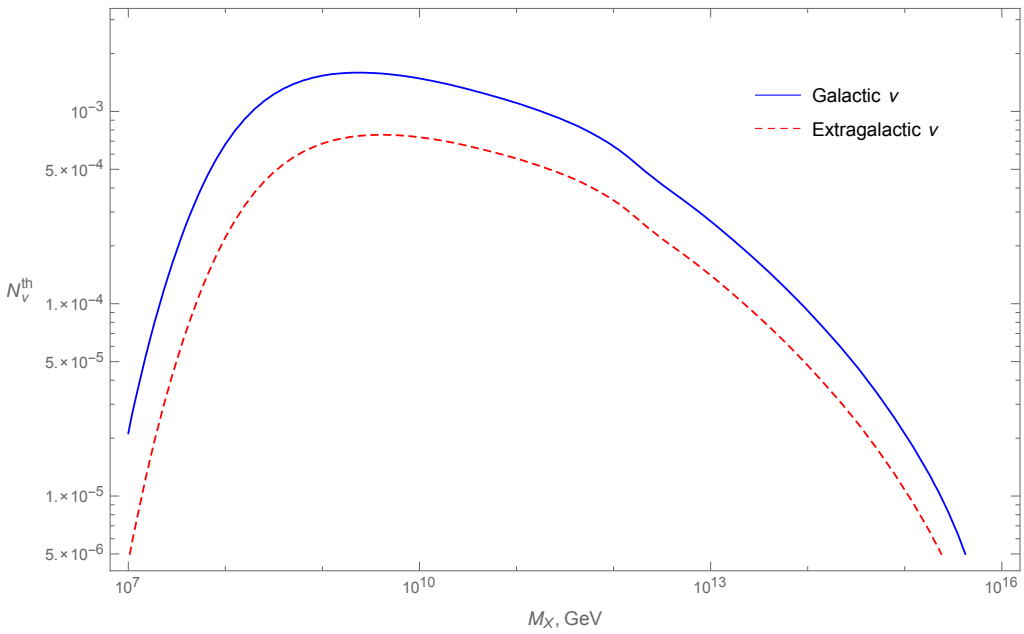

Figure 2: Total number of galactic (solid line) and extragalactic (dashed line) neutrinos from decays of dark matter particles with various masses M_X and $\tau = 10^{24}$ yr as it could be received by the IceCube experiment.

calculate $N_{\rm th}$ over full range of accessible energies or in separate energy intervals. It is obvious that in the latter case the constraints on parameter τ proved to be weaker, since the number $N_{\rm limit}$ does not depend on the length of the energy interval. In the case when $N_{\rm bg} = 0$ and $N_{\rm obs} \neq 0$, the other method is more appropriate. We split the full energy range in separate intervals ΔE^i with certain $N^i_{\rm obs}$ in each one and generate Monte–Carlo set which places in the i-th interval the number of events complying the Poisson distribution with the mean $\lambda^i = N^i_{\rm th}$, the theoretical number of events calculated in the respective energy interval ΔE^i. For the particular value of the parameter τ we generate some large number of these Monte–Carlo realisations. Then we vary the parameter τ until the fraction of realisations that exceeds the distribution of $N^i_{\rm obs}$ reaches $\beta =$ C.L.. In the case of all $N^i_{\rm obs} = 0$ this method yields the same results as the Feldman–Cousins technique. While for $N_{\rm obs} > 0$ the constraints of the Monte–Carlo method appears somewhat stronger. In the IceCube dataset we neglect the background of $0.064^{+0.023}_{-0.039}$ atmospheric neutrino events and therefore can apply the described method.

The constraints on the parameter space $\{M_X, \tau\}$ are presented in Fig. 5 together with the constraints of works [30] as well as the gamma–ray constraints obtained in our previous work [23]. We can see that the gamma–ray constraints overlap the neutrino ones in almost all dark–matter mass range. Nevertheless neutrino observation remains a crucial tool for dark matter indirect detection. The ratio of neutrino and photon fluxes from dark–matter particles decay have the certain value, which can not be reproduced by the photon and neutrino fluxes of other origin.

Some examples of neutrino fluxes from the decay of the dark matter with the marginally allowed lifetime are shown in Figs. 3—4 together with some competing astrophysical and cosmogenic neutrino fluxes. We can see that while the all–sky averaged fluxes of dark

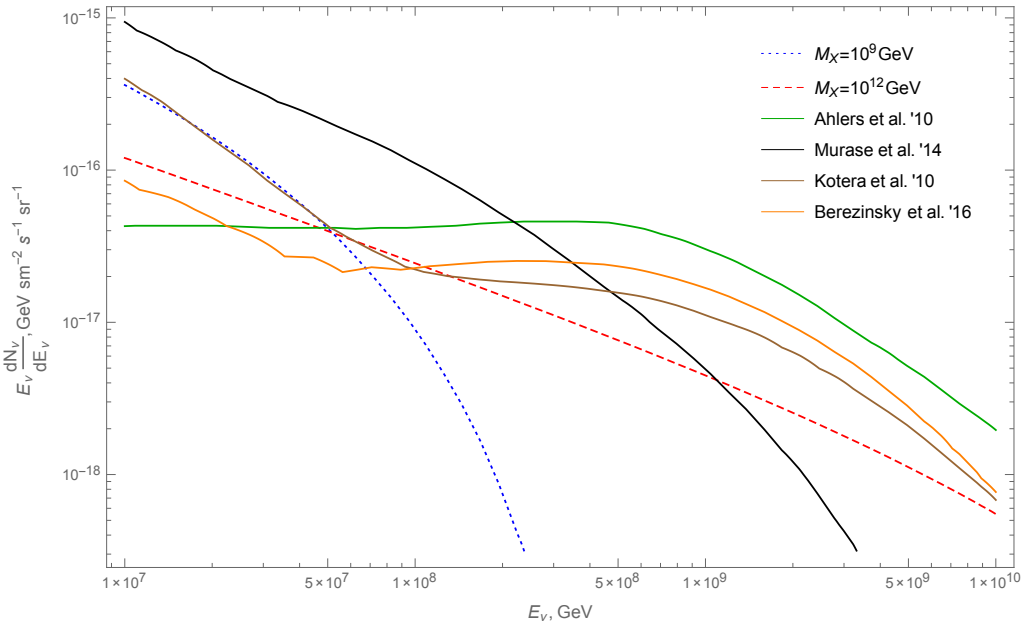

Figure 3: All–sky averaged neutrino fluxes from decays of dark–matter particles with masses $M_X = 10^9$ and $M_X = 10^{12}$ GeV and marginally allowed lifetime ($\tau = 6.6 \cdot 10^{20}$ and $\tau = 3.5 \cdot 10^{20}$ yr respectively) compared with various models of astrophysical [50] (solid black) and cosmogenic [47–49] neutrino fluxes (the sum of neutrino and antineutrino of all flavours).

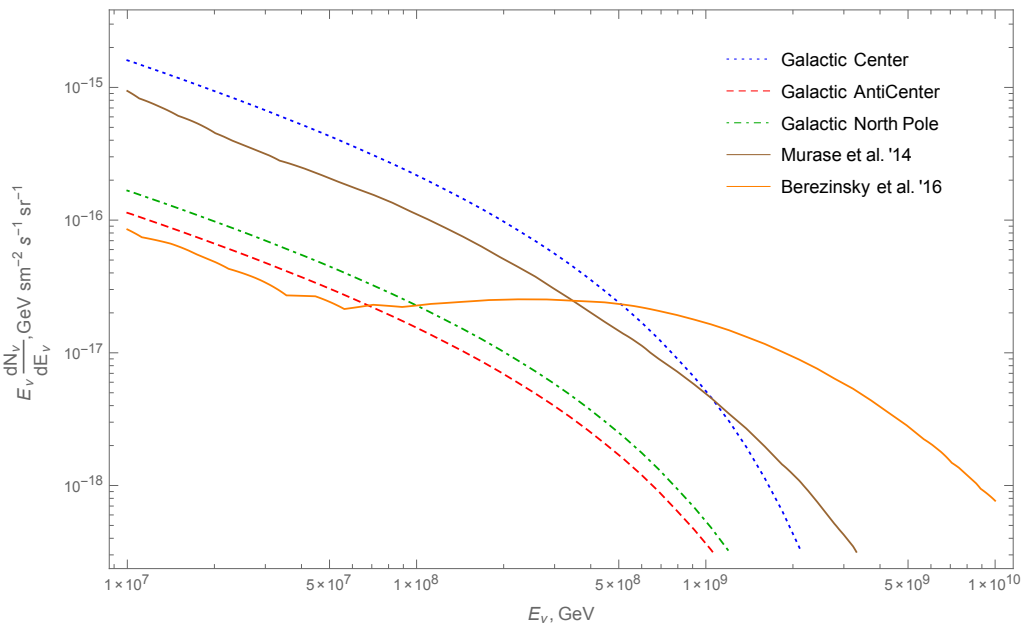

Figure 4: Comparison of neutrino fluxes from decays of dark–matter particles with mass $M_X = 10^{10}$ GeV and marginally allowed lifetime $\tau = 7.75 \cdot 10^{20}$ yr, coming from several directions with one model of astrophysical neutrino flux [50] (solid brown) and one model of cosmogenic one [47] (solid orange) (the sum of neutrino and antineutrino of all flavours).

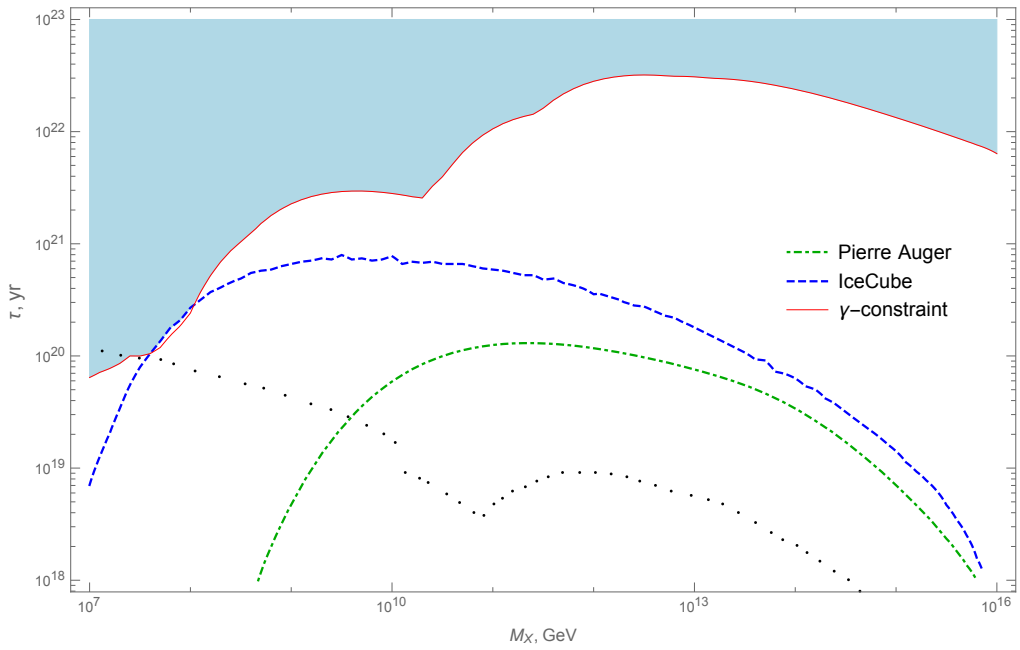

Figure 5: 90% C.L. exclusion plot for mass M_X and lifetime τ of dark–matter particles. White area is excluded. For comparison we present the constraints obtained with photon limits [23] (solid thin red line). We also show the constraint obtained in the dark matter model with $X \to \nu\bar{\nu}$ decay channel [30] (black dots).

matter decay are hardly distinguishable from astrophysical and cosmogenic ones, the signals from the Galactic center direction are not, manifesting the importance of analysis of the Galactic anisotropy in the indirect dark matter search.

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