ULTRA-HIGH ENERGY NEUTRINO ABSORPTION

BY NEUTRINO DARK MATTER

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Abstract

If the dark matter in the Universe consists of massive neutrinos ($m_{\nu_i} \simeq 30-100$ eV, $i = \mu, \tau$), cosmic ray $\nu_i$ of ultra-high energies may be absorbed at large redshifts, or even while crossing galactic haloes, due to their annihilation with the dark matter neutrinos. These processes can become very important for $E_\nu \simeq M_Z^2/2m_\nu \simeq 5 \times 10^{19}$ eV, for which the $Z$-exchange cross section is resonant, and may lead to a dip in the cosmic ray neutrino spectrum. I compute all the relevant contributions to the $\nu$ cross section and use them to evaluate the absorption redshift at ultra-high energies and to analyze the possible absorption in haloes of galaxies. If the determination of the neutrino spectrum at those energies were to become experimentally feasible, the observation of the features discussed here would provide a clear indication of the nature of the dark matter.

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I Introduction

Many sources of ultra-high energy (UHE) neutrinos, $E_{\nu} > 10^{16}$ eV, have been proposed in the past. Neutrinos can be produced in astrophysical objects by the decay of pions or kaons produced in the interaction of accelerated protons and nuclei with ambient protons and photons. The acceleration sources could be shocks of supernova explosions, pulsars, the active nuclei of Seyfert galaxies and quasars (AGN), a bright phase of galaxy formation, etc. [1]. The UHE neutrinos could also result from the decay of very heavy exotic particles, such as long-lived relics from the Big Bang [2] or, for instance, by ultra-heavy particle emission from saturated superconducting cosmic strings [3] (see however [4]), by mini-black hole evaporation [5], by cusp radiation from ordinary cosmic strings [6], etc.

The hadronic component of cosmic rays (CR) has been observed in air shower arrays up to energies $\sim 10^{20}$ eV, and experiments are being done and planned to continue the exploration of the spectrum in the range $10^{20}$–$10^{21}$ eV. Although the CR composition is not completely established, one expects that if there is a proton component up to the maximum observed CR energies the associated neutrino spectrum due to photopion production at the source will extend up to $E_{\nu} \sim E_{p,\text{max}}/20$, i.e. at least up to $5 \times 10^{18}$ eV (the $\pi$ has typically $\langle E_{\pi} \rangle \approx 0.2 E_p$ and the mean neutrino energy is $E_{\nu}/4 \approx E_p/20$). Furthermore, if the cosmic ray protons have a component originating from distant sources beyond the Local Cluster of galaxies, their interactions with the cosmic microwave photons would not only produce a secondary neutrino flux [7] but also give rise to the Greisen-Zatsepin-Kuzmin [8] cutoff in the proton spectrum at $E_p \sim 10^{20}$ eV (starting at $3 \times 10^{19}$ eV, the photopion production threshold for 3 K background photons). Hence, even if no CR nucleons were observed much above $10^{20}$ eV, it may still be that the CR-induced neutrino flux extends significantly beyond $10^{19}$ eV and arises from extragalactic cosmic rays. The other possibilities mentioned above, involving (more exotic) non-astrophysical sources, can produce important fluxes of neutrinos extending even up to $10^{25}$–$10^{28}$ eV.

Another very important aspect of neutrinos is the existence of a cosmic neutrino background, resulting from early decoupling just before the time of nucleosynthesis, which is expected to have an average density at present of $n_{e} + n_{\nu} \approx 108$ cm$^{-3}$ for each flavour. This has led to the nice possibility of explaining the Dark Matter (DM) problem just by invoking a small neutrino mass $m_{\nu_i} \approx 92 \Omega_{c} \hbar^{2}$ eV ($h$ is Hubble’s constant
in units of $100 \text{ km/s Mpc}, \ 0.4 < h < 1$. If this is actually the case, large enhancements in the $\nu_\ell$ (and $\bar{\nu}_\ell$) density are expected around galaxies to account for the dark haloes responsible for the observed flatness of galaxy rotation curves. For example, if we take for the halo density the usual parametrization

$$\rho_h(r) = \frac{\rho_0}{1 + (r/r_c)^\alpha},$$

where for our galaxy the estimated central density is $\rho_0 \simeq 0.014 M_\odot/\text{pc}^3$ and the core radius of the halo is $r_c \simeq 7.9 \text{ kpc}$ [9], then the neutrino number density inside the core of our galaxy will be $n_\nu = n_\sigma \simeq 4 \times 10^6/m_{\nu_5} \text{ cm}^{-3}$, i.e. a factor $\sim 10^5$ larger than the average relic $\nu$ density ($m_{\nu_5} \equiv m_\nu/50 \text{ eV}$). Hereafter $\Omega_\nu, \Omega_\sigma \simeq 1$ will be taken, assuming that at present 10–20 % of the DM neutrinos are in galactic haloes while the rest are distributed more smoothly, as suggested by numerical simulations [10].

I consider here the possible attenuation that the DM neutrinos may produce in the cosmic neutrino fluxes. As was shown many years ago by Weiler [11], the smooth DM neutrino background can attenuate neutrinos produced at large redshifts. I reanalyze those processes taking into account all the contributions to the cross section responsible for the CR $\nu$ absorption, and also show that the DM neutrinos clustered around galaxies may attenuate the CR neutrinos crossing dark haloes. Since the ‘target’ neutrinos are massive, many of the electroweak channels for neutrino-antineutrino annihilation become sizeable for $E_\nu \gtrsim 10^{19} \text{ eV}$ ($\sigma_{\nu\bar{\nu}} \gtrsim 10^{-34} \text{ cm}^2$) and there is a special enhancement at the pole of the $Z$ resonance, i.e. at an energy $E_{\nu\nu, \text{res}} = M_Z^2/2m_\nu \simeq 8 \times 10^{19} \text{ eV}/m_{\nu_5}$, at which $\sigma_{\nu\nu} \simeq 5 \times 10^{-31} \text{ cm}^2$. These processes could cause strong absorption of the neutrinos with energies close to $E_{\nu\nu, \text{res}}$, a value that may still be inside the cosmic ray neutrino continuum. Unlike the case considered here, if the background neutrinos were massless the effects would be much less important since the annihilation would become resonant only at energies $E_\nu \simeq 10^{25} \text{ eV}$ [11,12].

If the neutrino fluxes are of astrophysical origin, i.e. arising from light meson decays, one expects fluxes of $\nu_\mu$ and $\nu_\tau$ (and the antineutrinos) in the ratio 2:1, with no sizeable flux of $\nu_e$’s. In this case the most important effects would be expected if the dark matter is composed of muon neutrinos ($m_{\nu_\mu} > m_{\nu_\tau}$). Although this is not the outcome of see–saw models with a common heavy Majorana mass, a heavier $\nu_\mu$ can naturally result in scenarios of radiatively generated neutrino masses [13]. If the UHE neutrino fluxes
are instead due to exotic particle decays, it is plausible that there will be comparable fluxes of the different flavours, and the ones most affected by absorption will be those of the same flavour as the DM neutrinos.

In Section II the relevant cross sections for neutrino annihilation are computed. Section III discusses the attenuation of cosmological neutrinos at high redshifts, while Section IV considers their possible attenuation in dark haloes. Further discussion is given in Section V.

II Neutrino annihilation cross section

Since DM neutrinos are non-relativistic, the total centre-of-mass (c.m.) energy squared for the annihilation of a CR neutrino of energy $E_\nu$ in the lab frame, the Earth, in which the DM neutrinos (of mass $m_\nu$) are almost at rest, is just $s \simeq 2m_\nu E_\nu$. If the annihilations were instead with massless thermal neutrinos ($T = 1.9$ K), the total c.m. energy squared would be 5 orders of magnitude smaller and, since the cross section is correspondingly reduced (below the $Z$ peak, $\sigma_{Z} \propto s$), these other processes are in general negligible when compared with the annihilations with the massive ones.

Figure 1 shows the total neutrino absorption cross section as a function of the energy of the CR neutrino impinging on the DM one, together with the contribution of the different channels. The most important process is $Z$-exchange in the $s$-channel (long-dashed line), which has resonant behaviour. The corresponding cross section is

$$\sigma_Z(\nu_i\bar{\nu}_i \rightarrow f\bar{f}) = \sum_{f} \frac{2G_F}{3\pi} n_f s_P Z \frac{t^2_{3}(f) - 2t_{3}(f)Q_{f}s_W^2 + 2Q_{f}^2s_W^4}{2},$$

(2.1)

where the sum is over all fermions with $m_f < \sqrt{s}/2$, with charge $Q_{f}$, isospin $t_{3}(f)$, and $n_f = 1(3)$ for leptons (quarks) and their masses were neglected. In eq. (2.1) $P_{Z} \equiv M_{Z}^2/((s - M_{Z}^2)^2 + \Gamma_{Z}^2 M_{Z}^2)$ and $s_{W}^2 \equiv \sin^2 \theta_{W} \simeq 0.23$. Note that the annihilation (2.1) also includes the production of neutrinos, with a 20% branching ratio, and although this is not exactly an absorption process, the secondary $\nu$ fluxes it produces are depleted in energy with respect to the primary flux so that, for a typical power law initial spectrum, the associated ‘pile up’ effect may be neglected and the process will also appear as an absorption. This approximation is also especially appropriate for the study of the attenuation of the neutrino fluxes with energies in the resonance region, to which we pay particular attention below.
In addition to the $s$-channel $Z$ exchange, the $t$-channel $Z$ exchange contributes to $\nu_i \bar{\nu}_j \rightarrow \nu_i \bar{\nu}_j$ (short-long dashes), with

$$\sigma^t_Z (\nu_i \bar{\nu}_j \rightarrow \nu_i \bar{\nu}_j) = \frac{G_F^2}{2\pi} sF_1 \left( \frac{s}{M_Z^2} \right),$$

(2.2)

where $F_1(y) = \left[ y^2 + 2y - 2(1 + y) \ln(1 + y) \right]/y^3$, while the interference of $s$ and $t$ channels (for $i = j$) is

$$\sigma^{st}_Z (\nu_i \bar{\nu}_i \rightarrow \nu_i \bar{\nu}_i) = \frac{G_F^2}{2\pi} P_2(s - M_Z^2) sF_2 \left( \frac{s}{M_Z^2} \right),$$

(2.3)

with $F_2(y) = \left[ 3y^2 + 2y - 2(1 + y)^2 \ln(1 + y) \right]/y^3$.

Another relevant process, which becomes very important at energies beyond the $Z$-peak, is the $t$-channel $W$ exchange (dotted line), with

$$\sigma_W (\nu_i \bar{\nu}_j \rightarrow \ell_i \ell_j) = \frac{2G_F^2}{\pi} sF_1 \left( \frac{s}{M_W^2} \right).$$

(2.4)

Note that this interaction will be the most important process for the attenuation of the UHE neutrinos with a flavour different from that of the DM ones. If they are both of the same flavour, $i = j$, one must also include the interference of the processes (2.1) and (2.4), that is

$$\sigma_{WZ} (\nu_i \bar{\nu}_i \rightarrow \ell_i \ell_j) = \frac{2G_F^2}{\pi} \left( \frac{s}{M_W^2} - \frac{1}{2} \right) P_2(s - M_Z^2) sF_2 \left( \frac{s}{M_W^2} \right).$$

(2.5)

For $E_\nu > 2M_W^2/m_\nu$, the annihilation into $W^+ W^-$ pairs is kinematically allowed, proceeding both through $t$-channel lepton exchange and $s$-channel $Z$ exchange; it gives (dot-long-dashed line)

$$\sigma(\nu_i \bar{\nu}_i \rightarrow W^+ W^-) = \frac{G_F^2 s^2}{12\pi} \left[ \frac{\beta^2 M_W^4}{(s - M_Z^2)^4} (12 + 20y + y^2) + \frac{2M_W^4}{M_Z^2} \left( 24 + 28y - 18y^2 - y^3 + \frac{48}{\beta y}(1 + 2y)L \right) + \frac{1}{y^2} \left[ y^2 + 20y - 48 - \frac{48}{\beta y}(2 - y)L \right] \right],$$

(2.6)

with $\beta \equiv \sqrt{1 - 4M_W^2/s}$, $L \equiv \ln \left( \frac{14y}{12y} \right)$ and $y \equiv s/M_W^2$. All (anti)neutrino flavours may also scatter elastically off DM (anti)neutrinos by $t$-channel $Z$ exchange (the scattering off massless flavours will be negligible with respect to that off DM $\nu$‘s). The corresponding cross section (short dashes) is

$$\sigma(\nu_i \nu_j \rightarrow \nu_i \nu_j) = \frac{G_F^2 M_Z^2}{2\pi} \frac{s}{s + M_Z^2}, \quad i \neq j,$$

(2.7)
while for \(i = j\) there is also the contribution of the \(u\)-channel, giving \((\text{dot-short dashes})\)

\[
\sigma(\nu_i\nu_i \rightarrow \nu_i\nu_i) = \frac{G_F^2 M_Z^2}{\pi} \left( \frac{s}{s + M_Z^2} + \frac{2M_Z^2}{(2M_Z^2 + s)} \ln \left( 1 + \frac{s}{M_Z^2} \right) \right).
\]  

(2.8)

We again neglect the secondary \(\nu\) flux that this process produces.

The sum of all the contributions to the \(\nu_i\) cross section, shown by the full line in fig. 1, determines the absorption probability of the neutrino flux with the same flavour, \(i\), as the DM neutrinos. On the other hand, if the CR neutrino flavour differs from that of the DM, the absorption will be due to the processes in eqs. (2.4), (2.2) and (2.7).

### III Absorption by the cosmic DM \(\nu\) background

As was shown in the previous section, the \(\nu\bar{\nu}\) annihilation cross section reaches a maximum of \(5 \times 10^{-31} \text{ cm}^2\) at the \(Z\) resonance peak. However, even UHE neutrinos of this energy are not absorbed by the diffuse background of DM neutrinos if they are of non-cosmological origin. In fact, the corresponding mean free path \(l_o = (n_o \sigma_{vZ})^{-1} \approx 4 \times 10^{28}\) cm is just above the present size of the horizon, \(H_0^{-1} \approx 10^{28}\) cm.

If instead the UHE neutrinos are produced at large redshifts, the effects of absorption are much more important [11]. In this case one must take into account that the DM \(\nu\) density increases as \((1+z)^3\) and also that the UHE \(\nu\) energy is affected by redshift, \(E(z) = (1+z)E_0\). As long as redshifts \(z < 10^5\) are considered, the DM neutrinos will still be non-relativistic and the cross sections obtained in the previous section can still be used with \(s = 2m_\nu E_\nu(z)\). Taking this into account, the absorption time can be computed as a function of redshift for a given present neutrino energy \(E_0\),

\[
\tau_v(z, E_0) \equiv [c n_o(z) \sigma_{vZ}(s = 2m_\nu(1+z)E_0)]^{-1}.
\]  

(3.1)

The lowest redshift at which absorption of the UHE neutrinos becomes important, \(z_{\text{abs}}\), can be obtained by equating the absorption time to the Hubble time of the Universe \(H^{-1}\), which is, in the matter dominated era \(z < z_{eq} \approx 2 \times 10^4 h^2\): \(H^{-1}(z) = \frac{3H_0}{2(1+z)^3/2}\).

\[
H^{-1}(z) = \frac{3H_0}{2(1+z)^{3/2}}.
\]  

(3.2)
with \( t_0 = 2.1 \times 10^{17} h^{-1} \) s the present age of the Universe. The resulting absorption redshift is plotted as a function of the present \( \nu \) energy \( E_0 \) in fig. 2 (taking \( h = 1 \)). If a source produced UHE neutrinos at redshifts \( z_* \) with energy \( E_* = (1 + z_*)E_0 \), the neutrino absorption by the DM \( \nu \) background will be sizeable as long as \( z_* > z_{\text{abs}}(E_0) \). The resonant effect in the annihilation cross section is apparent in the figure, and may lead to absorption even for redshifts as low as \( z_{\text{min}} = 2 \) and for present energies \( E_0 \lesssim E_{\text{res}}/(z_{\text{min}} + 1) \simeq 3 \times 10^{19} \) eV/\( m_\nu \). In particular, if a high-redshift quasar or a bright phase of galaxy formation produced UHE neutrinos at a redshift of, say, \( z_* = 5 \), the neutrinos with present energy \( 1-3 \times 10^{19} \) eV/\( m_\nu \) would have been almost completely absorbed. The effect becomes even stronger for sources at larger redshifts, and the Universe becomes completely opaque to UHE \( \nu \)'s of the same flavour as the DM neutrinos for \( z \gtrsim 300 \).

Figure 2 also shows, with dotted lines, the absorption redshift for neutrinos with a flavour different from the DM neutrinos, they are attenuated by the processes in eqs. (2.4), (2.2) and (2.7). It follows that the ‘horizon’ of UHE neutrinos is constrained in this case to \( z \lesssim 3 \times 10^2-10^3 \).

IV Absorption by neutrino haloes

If massive neutrinos constitute the DM, it is also possible that the large enhancement in the \( \nu \) density in galactic haloes may lead to a sizeable absorption of UHE neutrinos. The relevant quantity to consider here is the \( \nu \) column density that an UHE neutrino has to cross while traversing the halo. For extragalactic neutrinos, absorption can occur both in the halo of our galaxy and in the halo around the source. In particular, if the \( \nu \) source is in the centre of a galaxy (such as on AGN), with a halo density described by eq. (1.1), the UHE neutrinos will cross a column density

\[
N_\nu = N_\tau = \frac{\pi \rho_0 v_\nu}{4 m_\nu} \simeq 1.8 \times 10^{30} \left( \frac{\rho_0}{0.1 M_\odot/pc^2} \right) \frac{r_{10}}{m_\nu} \text{ cm}^{-2}
\]

on their way out of the source, with \( r_{10} \equiv r_\nu/10 \) kpc. The probability for the UHE neutrinos to annihilate with the DM on their way through the halo is then

\[
P_{\text{abs}}(E_\nu) = 1 - e^{-N_\nu \sigma_{\nu\nu}(E_\nu)}.
\]
The ‘survival’ probability, $1 - P_{\text{abs}}$, is shown as a function of $E_0$ in fig. 3, assuming that the source is at $z = 0$ and for $m_\nu = 5.0$ eV (for $z \neq 0$, the dip in the spectrum due to the absorption at the source would be correspondingly redshifted to $E = E_{\text{res}}/(1 + z)$). The three curves correspond to $X \equiv \rho_0 r_c[(M_\odot / \text{pc}^3) \text{kp}^{-1}]^{-1} = 0.1, 1$ and $10$, spanning the range expected for DM haloes of non-dwarf galaxies ($X \simeq 0.1$ for our galaxy and it is expected to increase in early-type galaxies. In particular, ellipticals are generally more massive and compact than spirals). From fig. 3 we see that the halo of our own galaxy can only marginally absorb neutrinos coming from the other side of the galaxy with energies close to the resonance value. One should keep in mind however that fits to halo densities without assuming a ‘minimal halo’ lead to significantly larger column densities than the usually quoted ones, that rely on that assumption. We also note that in rich clusters, where AGN are likely to be found, the DM density might be smeared to Mpc scales (tidal forces between nearby galaxies can strip individual galactic haloes), and although the average DM density of the cluster ($\lesssim 10^{-3} M_\odot / \text{pc}^3$) is lower than in galactic halo models, the DM column density of the cluster could also correspond to $X \gtrsim 0.1$.

One can obtain an upper bound for the absorption that a halo made of neutrinos can produce by considering the ‘maximum neutrino halo’ allowed by the Tremaine–Gunn phase-space constraint [14]. If the halo is modelled as an isothermal sphere [15] with velocity dispersion $\sigma = \sqrt{\langle v_x^2 \rangle} = \sqrt{\langle v_y^2 \rangle} = \sqrt{\langle v_z^2 \rangle}$ (related to the circular velocity of rotation curves by $v_c = \sqrt{2}\sigma$), the the Tremaine–Gunn constraint is

$$\frac{\rho_0}{m_\nu^2 (2\pi\sigma^2)^{3/2}} < \frac{g_\nu}{h^2},$$

leading to

$$\rho_0 r_c < 4.0 g_\nu m_\nu^4 c^2 \sigma_{200}^2 r_{10}^3 \frac{M_\odot}{\text{pc}^3} \text{kp},$$

where $g_\nu$ is the number of occupied spin states contributing to the DM density, i.e. $g_\nu = 1$ if just one flavour of left-handed neutrinos contributes and $\sigma_{200} \equiv \sigma / 200 \text{ km/s}$. It is then clear that to have significant absorption, i.e. $X \gtrsim 1$, large values of $m_\nu$ would be preferable and the effects can be expected to be more important in haloes where the neutrinos have large dispersion velocities. (We should note, however, that if phase-space anisotropies were allowed, the bound in eq. (4.4) would be looser [16].)
In the isothermal sphere model, the following relation holds

\[ \rho_\text{cc} = \frac{9\sigma^2}{4\pi G_N r_c} \simeq 0.61 \frac{\sigma \, M_\odot}{r_\text{10} \, \text{pc}^2} \text{ kpc}, \tag{4.5} \]

where \( G_N \) is Newton’s constant, from which we see that values of \( X \) of order unity seem quite reasonable for the expected ranges of \( \sigma (\lesssim 300 \text{ km/s}) \) and \( r_c (1-10 \text{ kpc}) \). (Actually, eq. (4.5) defines the King radius \( r_c \), sometimes called core radius since \( \rho(r_c) \simeq \rho_0/2 \), but eq. (1.1) is not a good fit to \( \rho(r) \) for the isothermal sphere at \( r \gg r_c \) [15].)

V Discussion

It was shown that DM neutrinos can give rise to a sizeable absorption of UHE neutrinos with energies close to the \( Z \) resonance value. If the absorption happened at high redshifts \( (z > 2) \) because of annihilations with the smoothly distributed cosmic neutrino background, there would be a dip in the UHE \( \nu \) spectrum that should appear at the redshifted energies \( E_{\text{res}}/(1+z) \lesssim E \lesssim E_{\text{res}}/3 \). If the source produces neutrinos at a very early epoch, \( z \gtrsim 300 \), the absorption of the UHE neutrinos will be total. It was also shown that UHE \( \nu \)'s may be absorbed while crossing haloes of very massive galaxies, leading to a dip in the spectrum around \( E_{\text{res}} \simeq 8 \times 10^{13} \text{ eV}/m_{\nu/0} \). Although it is unlikely that a high-redshift astrophysical source be powerful enough to individually produce large UHE \( \nu \) fluxes, it could be that the sum over large \( z \) sources does lead to a sizeable \( \nu \) flux [17], but in this case the dip in the spectrum will be somewhat smeared due to the difference in the energies of the dips for each individual redshifted source. On the other hand, if the absorption is produced in haloes around sources at small \( z \) (that need not be so powerful since they are closer), their combined fluxes will show the dip always at \( E_{\text{res}} \) but, however, the absorption may be reduced by the contribution of sources surrounded by haloes with DM column density not sufficiently large to produce a sizeable absorption \( (X \ll 1) \).

On the observational side, we know that no UHE neutrinos have yet been detected, although neutrino astronomy is still in its infancy. At \( E_\nu \gtrsim 10^{15} \text{ eV} \), one should take into account the fact that the Earth is no longer transparent to neutrinos [18], but on the other hand the atmospheric \( \nu \) flux induced by hadronic cosmic rays drops to very low values at ultra-high energies. Neutrinos of UHE should produce giant horizontal air
shower by interacting deep in the atmosphere or in the crust of the Earth, and the non-observation of these showers has already been used to set bounds \([19]\) on energetic neutrinos from AGN \([20]\). Other methods proposed for UHE neutrino detection are acoustic detection and the observation of radio waves generated by the showers of charged particles produced in dielectric media (the Antarctic ice or even the surface of the Moon) \([21]\). These techniques could allow for much larger detector sizes (as required for UHE \(\nu\) detection) than those of the present Cherenkov detectors. In the near future, large optical Cherenkov detectors deep in the sea (DUMAND) or in the ice (AMANDA) remain a promising tool for UHE \(\nu\) searches. It is not clear whether detection methods sensitive to the features of the UHE neutrino spectrum discussed here will become available (provided that \(\nu\) fluxes of \(10^{21}\) eV exist), but certainly a positive determination of these effects would have far-reaching consequences, giving an indication in favour of neutrino dark matter, even pointing to a value for the neutrino mass, besides providing information regarding the UHE neutrino sources themselves.

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References

Figure captions

Fig. 1: Total $\nu_1\bar{\nu}_1$ annihilation cross section (continuous line) and individual contributions to the different neutrino processes, as a function of the cosmic ray neutrino energy $E_\nu$ (times $m_{\nu_1} \equiv m_\nu / 50 \text{ eV}$).

Fig. 2: Absorption redshift for neutrinos of the same flavour as the DM ones (solid line) and for those of different flavour (dotted line), as a function of the present neutrino energy $E_0$.

Fig. 3: Survival probability for a neutrino leaving from the centre of a halo with density given by eq. (1.1) and for $X \equiv \rho r_c / (M_\odot / pc^3 \text{ kpc}) = 0.1, 1$ and $10$, as a function of the neutrino energy.