Antineutrino-neutrino and antineutrino-electron resonant annihilation through rho and other vector mesons

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Abstract

The $\bar{\nu}e^- \rightarrow \pi^-\pi^0$ and $\nu e^+ \rightarrow \pi^+\pi^0$ reactions have a resonant structure whenever the energy of the $s$–channel equals the mass of a $J^P = 1^-$ vector meson. The resonant cross sections are of the order $10^{-38}$ cm$^2$ and correspond to cosmic ray antineutrino energies in the range 0.20 to 2.00 TeV. A similar structure occurs in the annihilation of cosmic ray antineutrinos with a relic (background) neutrino. For $m_\nu \leq 10^{-3}$ eV the resonant energy is above GZK limit and their decay products include multiple $\gamma$–rays. Possible detection schemes are discussed, especially those which rely on Cherenkov radiation.
1 Neutrino–initiated reactions in water.

The first-generation neutrino telescopes (Lake Baikal, AMANDA) are in operation and the next-generation telescopes are under construction. The main purpose of these experiments is to detect cosmic neutrinos (and/or antineutrinos) with energy in TeV-EeV region, which are believed to be coming to the earth from gamma-ray bursts and active galactic nuclei. For their detection the following neutrino interactions in water are usually considered.

1. Charged–current induced neutrino-nucleon reactions

\[ \nu_l N \rightarrow l^- X \quad \sigma(\nu_N) \approx 0.67 \cdot 10^{-38} \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2, \]  

(1)

and antineutrino interactions

\[ \bar{\nu}_l N \rightarrow l^+ X \quad \sigma(\bar{\nu}_N) \approx 0.34 \cdot 10^{-38} \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2. \]  

(2)

Here the cross sections are averaged over proton and neutron; the reactions are sensitive to all neutrino flavours. The above cross-sections are valid at energies in the range \( m_l^2/2m_N \ll E_\nu \ll m_W^2/2m_N \). It is these reactions that the experiments rely upon for detection of atmospheric neutrinos. The high-energy recoil lepton produces Cherenkov light and is detected as the so-called one-ring event. These cross sections are the largest among the reactions available in the experiments. The same reactions on nuclei are calculated as the incoherent sum of protons and neutrons with the nuclear corrections being small.

2. Neutral–current induced neutrino-nucleus interactions have cross sections three times smaller and are not considered, because it is hard to identify the final products.

3. Elastic neutrino-electron and antineutrino-electron scattering

\[ \nu_l e^- \rightarrow \nu_l e^- \quad \bar{\nu}_l e^- \rightarrow \bar{\nu}_l e^- \]

reactions are also sensitive to all three flavours of neutrinos. The high-energy recoil electron is observed through the produced Cherenkov light. The corresponding cross-sections (for \( m_l^2/2m_e \ll E_\nu \ll m_W^2/2m_e = 6.4 \cdot 10^9 \text{ GeV} \)) are as follows:
The purpose of this paper is to call attention to $\bar{\nu}_l l^-$ (and $\nu_l l^+$) resonant annihilation mediated by a negatively– (and positively) charged vector meson $V$ with $J^P = 1^-$. Among the charged leptons only electrons are contained in targets on the earth. Thus only the reaction $\bar{\nu}_e e^-$, which is sensitive to electron antineutrinos in energy range from about $\sim 200 \text{ GeV}$ to $\sim 2000 \text{ GeV}$, is of particular interest for experimental physics.

\[
\sigma(\nu_e e^- \rightarrow \nu_e e^-) = \frac{G_F^2 m_e E_\nu}{2\pi} \left[ (2 \sin^2 \theta_W + 1)^2 + \frac{4}{3} \sin^4 \theta_W \right] \\
\approx 0.9 \cdot 10^{-41} \cdot \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2
\]

\[
\sigma(\bar{\nu}_e e^- \rightarrow \bar{\nu}_e e^-) = \frac{G_F^2 m_e E_\nu}{2\pi} \left[ \frac{1}{3} (2 \sin^2 \theta_W + 1)^2 + \frac{4}{3} \sin^4 \theta_W \right] \\
\approx 0.378 \cdot 10^{-41} \cdot \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2
\]

\[
\sigma(\nu_{\mu(\tau)} e^- \rightarrow \nu_{\mu(\tau)} e^-) = \frac{G_F^2 m_e E_\nu}{2\pi} \left[ (2 \sin^2 \theta_W - 1)^2 + \frac{4}{3} \sin^4 \theta_W \right] \\
\approx 0.15 \cdot 10^{-41} \cdot \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2
\]

\[
\sigma(\bar{\nu}_{\mu(\tau)} e^- \rightarrow \bar{\nu}_{\mu(\tau)} e^-) = \frac{G_F^2 m_e E_\nu}{2\pi} \left[ \frac{1}{3} (2 \sin^2 \theta_W - 1)^2 + \frac{4}{3} \sin^4 \theta_W \right] \\
\approx 0.14 \cdot 10^{-41} \cdot \left( \frac{E_\nu}{1 \text{ GeV}} \right) \text{ cm}^2
\]

Figure 1: Antineutrino-electron resonant annihilation through charged $V$-meson with $J^P = 1^-$. The main contribution to the $\bar{\nu}_e e^-$ annihilation comes from the resonance in the intermediate state, which corresponds to neutrino energies $E_{\nu(\text{res})} = m_V^2 / 2m_e$. The calculation of this process becomes obvious when we keep in mind calculations with the vector meson dominance.
A similar process of high energy neutrino (antineutrino) annihilation takes place with the relic background antineutrino (neutrino) mediated by a neutral vector meson $V^0$. It would be remarkable if this process could serve as a means for the detection of the relic background neutrinos and we shall discuss it in section 3.

2 Antineutrino-electron resonant annihilation

The lightest $J^P = 1^-$-meson is the $\rho^- (770)$-meson and we start our discussion with this particle. Recent observations of the $W^- \rightarrow \rho^-$ transition, which clearly determine the $W - \rho$ coupling, come from the $\tau-$decays $\tau^- \rightarrow \nu_\tau \pi^- \pi^0$ [3].

![Diagram](tau_nu_rho.png)

Figure 2: Decay $\tau^- \rightarrow \nu_\tau \pi^- \pi^0$ through the $\rho^-$-meson.

The process $\tau^- \rightarrow \nu_\tau \rho^- \rightarrow \nu_\tau \pi^- \pi^0$ contributes $(25.4 \pm 0.14)\%$ to the total $\tau-$width, while the process without $\rho^-$ in the intermediate state contributes at most $(0.3 \pm 0.32)\%$. It is clear, that pions in the final state are coming completely from the $\rho-$meson decay, since the $\pi^- \pi^0$ decay channel contributes $\sim 100\%$ [3] to the total width.

We use the central values of the data to extract the phenomenological $\tau \rho \nu_\tau$ coupling. Indeed, following the discussion by Okun [4], we write the matrix element of the $\tau^- \rightarrow \nu_\tau \rho^-$ decay as

$$ M = \frac{G_F}{\sqrt{2}} g_\rho \cos \theta_c \varphi_\rho^{(\rho)} \bar{u}_\nu \gamma^\alpha (1 - \gamma_5) u_\tau, $$

where $\theta_c$ is the Cabbibo angle, $\varphi_\rho^{(\rho)}$ is $\rho-$meson wave function. The result of the width
calculation reads
\[
(0.254 - 0.003) \cdot \Gamma_{\tau(tot)} = \frac{G^2}{16\pi} g^2_{\rho} \cos^2 \theta_c \frac{m^2_\tau}{m^2_\rho} \left(1 - \frac{m^2_\rho}{m^2_\tau}\right)^2 \left(1 + \frac{2m^2_\rho}{m^2_\tau}\right),
\]
(4)

which for \( m_\tau = 1.777 \text{ GeV} \), \( \Gamma_{\tau(tot)} = 2.26 \cdot 10^{-12} \text{ GeV} \) and \( m_\rho = 0.7665 \text{ GeV} \) gives
\[
g^2_{\rho} \cos^2 \theta_c = 0.02435.
\]
(5)

We subtracted in Eq. (4) a 0.3\% contribution for the non–resonant background. Recalling that \( \cos \theta_c = 0.974 \), our result is in agreement with the values obtained from indirect estimates, mentioned in [4]. This value for \( g^2_{\rho} \cos^2 \theta_c \) is obtained for \( \rho \) on–the–mass–shell.

For virtual \( \rho \)–meson, strictly speaking, \( g^2_{\rho} \cos^2 \theta_c \) can vary with center-of-mass energy \( \sqrt{s} \). However, up to now all calculations within the vector meson dominance framework (which are very extensive, for example, for neutrino–nucleus scattering [5]) never take into account this dependence. Since in our case the dominant contribution to annihilation process occurs at \( \sqrt{s} = m_\rho \pm \Gamma_\rho/2 \) (i.e. \( \rho \) is on-mass-shell or almost on-mass-shell), we neglect the variation of the coupling constant with \( \sqrt{s} \).

Another quantity we need to know for the calculation is the \( \rho \pi \pi \) vertex, which is obtained from the \( \rho \)–meson width. The vertex, according to the effective gauge model of hadron interactions, is written as \( if_{\rho\pi\pi}(p^\mu_1 - p^\mu_2) \). The width in the \( \rho \)–rest frame is
\[
\Gamma_{(0)}(\rho \rightarrow \pi^- \pi^0) = \frac{2}{3} \frac{f^2_{\rho\pi\pi}}{4\pi} |p_1|^3 \frac{m^2_\rho}{m^2_\rho}, \quad |p_1| = \sqrt{m^2_\rho - 4 \cdot m^2_\pi}.
\]
(6)

For \( \Gamma_{\rho(0)} = 0.15 \text{ GeV} \), \( m_\pi = 0.137 \text{ GeV} \) one obtains from the width
\[
f^2_{\rho\pi\pi} = 2.896.
\]
(7)

Now it is straightforward to write down the amplitude for the resonant antineutrino–electron annihilation process \( \bar{\nu}_e e^- \rightarrow \rho^- \rightarrow \pi^- \pi^0 \)
\[
M = \frac{G_F}{\sqrt{2}} g_{\rho} \cos \theta_c \bar{\nu}_e \gamma^\alpha (1 - \gamma^5) u_e \frac{-g_{\alpha\beta} + g_{\alpha\beta} \frac{m^2_\pi}{m^2_\rho}}{q^2 - m^2_\rho + i\Gamma_\rho m_\rho} (if_{\rho\pi\pi})(p^\beta_1 - p^\beta_2),
\]
(8)
where \( q = p_e + p_\nu \), and to calculate the cross section

\[
\sigma = \frac{G_F^2 g_{\mu \pi}^2}{4 \pi} f^2 \cos^2 \theta e \frac{2 m_e E_\nu}{(2 m_e E_\nu - m_\rho^2)^2 + 2 m_e E_\nu \cdot \Gamma^2_{\rho(0)}} \left( 1 - \frac{4 m_\pi^2}{2 m_e E_\nu} \right)^{3/2}
\]

At resonance the cross-section is

\[
\sigma_{\text{res}}(\bar{\nu}_e e^- \rightarrow \rho^- \rightarrow \pi^- \pi^0) = 4.4 \cdot 10^{-38} \text{ cm}^2.
\]

With this cross-section one can easily estimate the expected number of events in a water detector of volume \( V \)

\[
N = \int d\Omega N_A \rho_{H_2O} V \int_{E_\nu(\text{min})}^{E_\nu(\text{max})} j(E_\nu, \theta) \sigma(E_\nu) dE_\nu
\]

The resonant neutrino energy for this process is \( E_{\nu(\text{res})} = m_\rho^2/2m_e = 580 \text{ GeV} \). The main contribution to the expected number of events comes from the energy interval \( E_{\text{min}} = (m_\rho - \Gamma_\rho)^2/2m_e = 370 \text{ GeV} \) to \( E_{\text{max}} = (m_\rho + \Gamma_\rho)^2/2m_e = 830 \text{ GeV} \). As the lowest estimate of the neutrino flux at these energies we can take the atmospheric neutrino flux \( d j(E_\nu)/dE_\nu(\theta, E_\nu) \), calculated by Volkova [6]. In what follows we neglect the dependence on the incident angle \( \theta \) and make calculations for the vertical flux.

In the energy region \( 10^2 \text{ GeV} - 10^6 \text{ GeV} \), two analytic formulas are given in Ref. [6]. The more accurate one (\( E_\nu \) must be given in GeV)

\[
j_{\bar{\nu}_e}(E_\nu) = \frac{1}{2} \cdot 2.4 \cdot 10^{-3} \cdot E_\nu^{-2.69} \left( \frac{0.05}{1+1.5 E_\nu/8760} + \frac{0.185}{1+1.5 E_\nu/1890} ight) + \frac{11.4 E_\nu^{0.083 - 0.215 \log E_\nu}}{1+1.21 E_\nu/1190}
\]

underestimates the flux in comparison to the values presented in the table of the same reference; the less accurate one

\[
j_{\nu_e}(E_\nu) = \frac{1}{2} \cdot 1.26 \cdot E_\nu^{-3.69}
\]

overestimates the flux. Another point of ambiguity is related to the values of \( E_{\text{min}} \) and \( E_{\text{max}} \), a convenient definition for narrow resonances is to take one width around the
resonance. Since the width of $\rho$–meson is about 20% of its mass and since we know the behaviour of the cross section and the flux in the whole energy region, the integration should be done rather over two- or even three-width intervals, instead of one–width only. The factor $[1 \div 3]$ in Eq. (14) reflects these ambiguities. Finally, the number of events obtained is

$$N = 20 \cdot [1 \div 3] \cdot \frac{V}{10^3 \text{ m}^3} \cdot \frac{\int d\Omega}{2\pi} \left( \frac{\text{events}}{\text{year}} \right)$$

(14)

The process under consideration will provide a distinct signature in the detector. For the resonant netrino energy $E_\nu = 580$ GeV the characteristic energy of outgoing pions is $E_{\pi^-} \sim E_{\pi^0} \sim E_\nu/2 \sim 290$ GeV. The characteristic longitudinal momentum of each pion has the same value, but the characteristic transverse momentum is about $\sqrt{2m_\pi E_\nu}/2 \sim 0.38$ GeV. Thus, the two pions move nearly collinearly with the angle between them being $\delta = \sqrt{2m_\pi/E_\nu} \approx 1.8 \cdot 10^{-3}$ rad.

Charge pions with such energy do not decay within the detector (the mean free path is about 15 kilometers), but can produce a nuclear shower. If not, the pion is detected by the Cherenkov light produced in a Cherenkov ring of almost maximal radius (the pion velocity is $\beta = (1 - 10^{-7})$ to be compared to the accuracy $\delta\beta/\beta \sim 10^{-4} - 10^{-5}$ that can be resolved nowadays at best). The half-angle $\theta_{\text{cher}}$ of the Cherenkov cone is

$$\theta_{\text{cher}}(\pi^-) = \arccos \frac{1}{n\beta},$$

(15)

where $n = 1.35$ is index of refraction for water.

The neutral pion decays on two photons, each photon having approximate energy of 145 GeV and producing an electromagnetic shower in the water. Since the energies of photons are high (> 0.4 GeV), they are not distinguishable by the detector. Since the initial photon energy is shared between many produced electrons, the energy of every electron is of a few MeV, the velocity is less than $\approx 0.95$, so the half-angle $\theta_{\text{cher}}$ of Cherenkov cone is smaller than that for the $\pi^-$. Thus, the final event from the $\rho^-$ decay is seen as two nearly collinear Cherenkov rings and can be easily separated from the elastic scattering initiated one–ring events.
The next vector meson to serve as intermediate state for the process in Fig. 1 is $K^{*-}(872)$ with $m_{K^{*-}} = 0.892$ GeV. Its total width $\Gamma_{K^{*-}} = 0.0508$ GeV originates completely from the channel $K^* \rightarrow K\pi$, so the resonant annihilation produces a pion and a kaon in the final state: $\bar{\nu}_e e^- \rightarrow K^{*-} \rightarrow (K\pi)^-$. 

The coupling for the $W - K^*$ transition also appears in $\tau-$decays. Indeed, the partial width of $\tau^- \rightarrow \nu_\tau K^{*-} \rightarrow \nu_\tau \pi K$ equals $(0.9 \pm 0.4)\%$ of the total $\tau-$width, while without $K^{*-}$ in the intermediate state the partial width is less than $< 0.17\%$. This gives

$$g_{K^*}^2 \sin^2 \theta_c = 1.124 \cdot 10^{-3},$$

which is one order smaller than for the $\rho-$meson.

To derive the $K^* K\pi$ vertex, we introduce for the $K^*$ decay the coupling $i f_{K^* K\pi} (p_1^\mu - p_2^\mu)$ and obtain

$$\frac{f_{K^* K\pi}^2}{4\pi} = 2.554.$$ 

The neutrino-antineutrino annihilation $\bar{\nu}_e e^- \rightarrow K^{*-} \rightarrow (K\pi)^-$ cross section in this case is

$$\sigma = \frac{G_F^2}{3} \frac{f_{K^* K\pi}^2}{4\pi} f_{K^*}^2 \sin^2 \theta_c \left( \frac{2m_e E_\nu}{(2m_e E_\nu - m_{K^{*-}})^2 + 2m_e E_\nu \cdot \Gamma_{K^{*-}(0)}^2} \times \left( 1 - \frac{m_\pi^2 + m_K^2}{m_e E_\nu} + \frac{(m_K^2 - m_\pi^2)^2}{m_e E_\nu^2} \right) \right)^{3/2}$$

The corresponding resonant neutrino energy is $E_{\nu(\text{res})} = 780$ GeV. At resonance the cross section equals $1.0 \cdot 10^{-38}$ cm$^2$ (which is 4 times smaller than for the $\rho^-$meson). The final states $\pi^0 K^-$ and $\pi^- K^0$ occur with equal probabilities. In the former case the
charged kaon and neutral pion can be observed as two nearly collinear rings (the same signature as if the $\rho^-$ was in the intermediate state). This increases the number of events in the detector by about 5%. In the latter case the neutral kaon can be either $K_L$ or $K_S$. The long–lived kaon $K_L$ does not decay within the detector (the mean free path is about 12 kilometers), so only the charged pion can be observed as one–ring event, which is impossible to distinguish from the muon resulting from the deep inelastic scattering.

The short–lived kaon $K_S$ with probability 31.4% decays (the mean free path is 20 meters) on 2 neutral pions and the event as a whole is again seen as two nearly collinear rings. This again increases the number of events by about 1.5%. With the probability 68.6% $K_S$ decays on two charged pions; in this case the final state is 3 charged pions which are seen as one-ring event.

Besides $\rho(770)$ and $K^*(892)$, several other charged $J^P = 1^-$ mesons are known: $\rho(1700)$, $K^*(1680)$, $\rho(1450)$, $K^*(1410)$. They can also be intermediate states of antineutrino-electron annihilation, but up to now there are no experimental data, which we can use in order to calculate their transition probability to $W^-$meson. However, it is not unnatural to suppose that their contribution to the antineutrino–electron annihilation cross section is of order $10^{-38} \text{ cm}^2$. Thus cross sections of this magnitude also appear in the interactions of cosmic neutrinos with energies up to $3 - 4 \text{ TeV}$.

Thus, the resonant neutrino annihilation events, when detected with neutrino telescopes, provide an estimate for the number of cosmic electron–type antineutrinos. Combined together with the total cross section events, it provides an estimate of the neutrino fluxes for electron–type and muon–type neutrinos separately. The importance of estimating the cosmic neutrino fluxes of various species and of antineutrinos with high enough accuracy is obvious because they are relevant for understanding their origin, i.e. how and where they are produced and the ratios of the fluxes when they arrive on the earth.

Finally we mention the Glashow resonance [7], which is mediated by the $W$–boson $\bar{\nu}_e e^- \rightarrow W \rightarrow \bar{\nu}_e e^-, \bar{\nu}_\mu \mu^-, \ldots$. The resonant reaction occurs for $E_\nu \approx 6 \cdot 4 \times 10^{15} \text{ eV}$
and has a large cross section $\approx 10^{-32} \text{ cm}^2$.

3 Antineutrino-neutrino annihilation.

As mentioned earlier, for experiments on the earth only the electrons are possible leptonic targets. On a larger scale, like regions of the universe, there are also relic neutrinos and antineutrinos which are part of the cosmic background radiation and have not been detected yet. Their density is estimated to be $54/\text{cm}^3$ for each species having a temperature of $1.9^\circ \text{ K}$.

A high energy neutrino of each species can annihilate with the corresponding relic antineutrino and vice versa with neutral vector mesons forming the intermediate state. At present we have good evidence from oscillation experiments that

$$\Delta m_{\text{atm}}^2 \sim 10^{-3} \text{ eV}^2 \quad \text{and}$$

$$\Delta m_{\text{solar}}^2 \sim 10^{-5} \text{ eV}^2 \quad \text{or less.} \quad (19)$$

They allow neutrino masses of the order of $0.01 - 1 \text{ eV}$. These values are not unique because there is still an arbitrary overall scale. Such small masses were justifiably neglected when we calculated the neutrino energy in the previous sections. The masses must be kept, however, when the relic neutrinos are considered as targets.

The annihilation of cosmic neutrinos with relic neutrinos with the $Z$–boson being an intermediate state was considered twenty years ago [8, 9] and its importance for cosmology and cosmic ray physics was described recently in [10] and [11], where possible implications for neutrino masses are discussed. More precisely, the resonant cosmic neutrino energy, $E_c$, for annihilation through the $Z$–boson is

$$E_c = \frac{M_Z^2}{2m_\nu} = 4 \times 10^{21} \text{ eV} \left(\frac{1 \text{ eV}}{m_\nu}\right) \quad (21)$$

with $m_\nu$ the mass of the lightest neutrino and the energy–averaged cross section is

$$\sigma(Z) = 4 \times 10^{-32} \text{cm}^2. \quad (22)$$
The Z–boson created in the reaction decays into several dozens of secondary particles (baryons, mesons, leptons and neutrinos) known as “Z–bursts”. A typical “Z–burst” contains 2.7 hadrons, 30 photons and 28 neutrinos. These secondaries form a highly collimated flux, the half-angle of the the cone being about $2 \cdot 10^{-11}$ rad. If the neutrino–antineutrino annihilation takes place and the following $Z$–decay occurs within a 20 pc zone around the earth and the flux is directed towards the earth, then the secondaries arrive on the earth simultaneously. Some of them are able to initiate the multiple air showers which could be searched for on the earth with air shower arrays. Therefore, showers of particles which arrive on earth in coincidence are unique signatures for the neutrino–antineutrino annihilation, with either the neutrino or antineutrino belonging to the cosmic background.

For the annihilation process through vector mesons, the resonant neutrino energy is smaller $E_\nu = M_V^2/2m_\nu$. For the $\rho^0$–meson the relevant reaction

$$\bar{\nu}_\ell + \nu_\ell \rightarrow \rho^0 \rightarrow \pi^+\pi^-$$

occurs at $E_{\text{res}} = 3 \times 10^{17} \text{ eV} (1 \text{ eV}/m_\nu)$. Comparing this reaction with the annihilation in figure 1, we note that the kinematics of the two reactions and $f_{\rho\pi\pi}$ are the same. The couplings $g_{\rho^-}$ (above it was denoted simply as $g_\rho$) and $g_{\rho^0}$ to the $W^–$ and the $Z^0$–bosons, respectively, are different. Both vector bosons couple to the isovector part of the weak vector current and are related by CVC.

The effective Lagrangian for strangeness conserving processes, relevant to our problem, is given by

$$\mathcal{L} = \frac{G}{\sqrt{2}} \left[ \bar{\nu}_\alpha(1 - \gamma_5)\nu (J_1^\alpha + iJ_2^\alpha) + \text{h.c.} + \bar{\nu}_\alpha(1 - \gamma_5)\nu (xV_3^\alpha - A_3^\alpha + yJ_s^\alpha) \right]$$

where $J_i = V_i - A_i$ is one of the isospin components of the usual $V - A$ currents and $J_s$ the isoscalar current. In the electroweak theory

$$x = 1 - 2 \sin^2 \theta_W \quad \text{and} \quad y = -2 \sin^2 \theta_W$$

with $J_s = \frac{1}{\sqrt{3}}V_8$, $\theta_W$ the weak mixing angle and we neglect strange quark contributions.
The $\rho$-mesons couple to the $V_1^\mu \pm iV_2^\mu$ and $V_3^\mu$ parts of vector current and their couplings are defined as

$$g_{\rho^-} = \langle \rho^- | V_1^\mu + iV_2^\mu | 0 \rangle$$

and

$$g_{\rho^0} = \langle \rho^0 | V_3^\mu | 0 \rangle.$$  

They are related by Clebsch–Gordon coefficients

$$\frac{g_-}{g_0} = \frac{\sqrt{2}}{\sqrt{3}} = -\sqrt{2}.$$  

The complete coupling of $\rho^0$ to neutrinos through the $Z$-boson is

$$x g_0 = -\frac{1}{\sqrt{2}} (1 - 2 \sin^2 \theta_W) g_{\rho^-} = -0.010 \quad \text{for} \quad \sin^2 \theta_W \approx 0.223.$$  

The resonant $\bar{\nu}\nu$ annihilation cross section, mediated by hadronic vector mesons is eight times smaller than the corresponding reaction with electrons as targets, i.e. $\sigma_{\bar{\nu}\nu} \sim 0.5 \times 10^{-38}\text{cm}^3$. This process contributes beyond the GZK cutoff only when the mass for one of the relic neutrinos is very small, $m_\nu < 10^{-3}\text{eV}$. In comparison to the Weiler process, the new cross sections are six order of magnitude smaller.

Multiparticle decays with baryons in the final state are impossible for $\rho$-mesons. The main decay channel $\rho^0 \rightarrow \pi^+\pi^-$ has no observable consequences, since the produced pions decay within $10^{-6}$ pc to muons and neutrinos of energy about $10^{16}$ eV, which are impossible to detect as specific events. The rare decays $\rho^0 \rightarrow \pi^0\gamma$ and $\rho^0 \rightarrow \eta\gamma$ (different channels are summarized in Table 1) produce 3 to 7 high-energy photons, with an opening angle between them being $\sim 10^{-9}$ rad.

The probabilities are summarized below in Table 2. Strictly speaking, the resonant cross section is calculated for two pions in the final state; however, there is no reason to believe that it differs significantly for other $\rho-$decays. If the annihilation events occur within a 0.2 pc distance from the earth, then these photons can be observed on the earth as simultaneously arriving $\gamma-$initiated showers. The number of showers is three or seven if the decaying $\rho$-meson moves toward the earth and fewer if it moves in another direction.
Table 1: Chains of $\rho^0$ meson decays and the corresponding probabilities.

| Decay                        | Probability |
|------------------------------|-------------|
| $\rho \to \pi^0 \gamma$     | 6.8 $\times$ 10$^{-4}$ |
| $\rho \to \eta \gamma$      | 2.4 $\times$ 10$^{-4}$ |
| $\pi^0 \to 2 \gamma$        | 1           |
| $\eta \to 3 \pi^0$          | 0.322       |
| $3 \pi^0 \to 6 \gamma$      | 1           |
| $\eta \to \pi^+ \pi^- \pi^0$| 0.283       |
| $\pi^0 \to 2 \gamma$        | 1           |

To estimate the relative probability to observe a multiparticle shower for different “bursts” it is enough to compare the two counting rates

$$P = \sigma \cdot \Delta E_{\nu} \cdot j \cdot w_{\text{earth}} \cdot l_{\text{earth}}^3 \cdot w_{\text{decay}},$$

with the variables and their values in Eq. (30) being specified in Table 2. We find

$$\frac{P(\rho)}{P(Z)} \sim 0.1 \cdot 10^{-2}$$

Thus the probability to observe a signal of $\rho$ or $K^*$ bursts is even smaller than that for $Z$–bursts.

Other light neutral vector mesons with $J^{PC} = 1^{--}$ series are the $\omega$, whose main decay channel is $\omega \to \pi^+ \pi^- \pi^0$, and the $\phi$, whose main decay channel is $\phi \to K^0_L K^0_S$ with the subsequent decays of $K^0_L$ and $K^0_S$ containing several photons in the final state. The resonant neutrino annihilation through these mesons also contributes to the events under discussion; but the $Z^0$–boson coupling to these mesons are unknown.

Considering the $K^{*0}$ meson as the intermediate state, its contribution to the resonant processes is negligible, because the $Z^0$–boson coupling to this meson is suppressed, since it involves a flavour–changing–neutral coupling. Heavy neutral mesons $\rho(1450)$, $\omega(1420)$, $\phi(1680)$, $\rho(1700)$, $\omega(1650)$ can also contribute to the resonant annihilation, but at the moment the values of their couplings to $Z^0$ are hard to estimate.
Table 2: Characteristic values for cosmological $Z$ and $\rho$ bursts.

| Characteristic                        | $Z$-burst | $\rho$-burst |
|---------------------------------------|-----------|--------------|
| resonant neutrino energy for $m_e = 1$ eV, $E_\nu$ | $4 \cdot 10^{21}$ eV | $3 \cdot 10^{17}$ eV |
| “one-width” energy interval, $\Delta E_\nu$ | $4.5 \cdot 10^{20}$ eV | $2.3 \cdot 10^{17}$ eV |
| energy–averaged annihilation cross section $\sigma_{av}$ | $4.2 \cdot 10^{-32}$ cm$^2$ | $0.3 \cdot 10^{-38}$ cm$^2$ |
| differential atmospheric neutrino flux at resonant neutrino energy if it satisfies the $E_{\nu^{-3}}$ behaviour estimated in [3], $j_0$ | $j_0$ | $2.5 \cdot 10^{12} \cdot j_0$ |
| angle of lab-frame cone | $\sim 10^{-11}$ | $\sim 10^{-9}$ |
| the probability to be directed towards the earth, $w_{\text{earth}}$ | $(\sim 10^{-11})^2$ | $(\sim 10^{-9})^2$ |
| distance to observe multishowers, $l_{\text{earth}}$ | $\sim 20$ pc | $\sim 0.2$ pc |
| probability of multishower per one annihilation act, $w_{\text{decay}}$ | hadronic | 0.7 |
| $3\gamma$ | | $8.4 \cdot 10^{-4}$ |
| $7\gamma$ | | $7.8 \cdot 10^{-5}$ |

4 Conclusions.

The cross section of $\bar{\nu}_l l^-$ and $\nu_l l^+$ reactions have a resonant behaviour in $s$–channel since the intermediate vector bosons couple, through vector meson dominance, to mesons with quantum numbers $J^{PC} = 1^{--}$, in particular with the $\rho(770)$ and $K^*(892)$. The resonant cross sections of reactions $\bar{\nu}_e e^- \rightarrow \rho^- \rightarrow \pi^- \pi^0$ and $\bar{\nu}_e e^- \rightarrow K^{*-} \rightarrow (K\pi)^-$ are calculated to be $4.4 \cdot 10^{-38}$ cm$^2$ and $1.0 \cdot 10^{-38}$ cm$^2$ respectively, and the resonant energies are $580$ GeV and $780$ GeV.

The $\bar{\nu}_e e^-$ reaction can be used for the experimental detection of the electron antineutrinos. In particular, the flux of atmospheric neutrinos with energies in the region of $\sim 200$ GeV to $\sim 2000$ GeV, may be visible at neutrino telescopes using water Cherenkov
detectors.

A similar process of resonant annihilation can occur in the universe. A high energy neutrino of any flavour can annihilate with the corresponding relic antineutrino, and vice versa, with neutral vector mesons in the intermediate state. If at least one of the neutrino species has a mass \( \leq 10^{-3} \) eV, then the resonant neutrino energy is \( \geq 3 \cdot 10^{20} \) eV. These events can lead to multiple the \( \gamma \)-primaries in cosmic rays beyond the GZK–cutoff. The neutrino–antineutrino resonant annihilation could also serve as an indication for the existence of relic neutrinos in the universe. The probability to observe some signatures on the earth is lower than that for the Weiler process [10].

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