Vetoing atmospheric neutrinos in a high energy neutrino telescope

Stefan Schönert,\textsuperscript{1} Thomas K. Gaisser,\textsuperscript{2} Elisa Resconi,\textsuperscript{1} Olaf Schulz\textsuperscript{1}

\textsuperscript{1}Max-Planck-Institut für Kernphysik
Saupfercheckweg 1, 69117 Heidelberg, Germany

\textsuperscript{2}Bartol Research Institute and Department of Physics \& Astronomy
University of Delaware, Newark, DE 19716 USA

Abstract

We discuss the possibility to suppress downward atmospheric neutrinos in a high energy neutrino telescope. This can be achieved by vetoing the muon which is produced by the same parent meson decaying in the atmosphere. In principle, atmospheric neutrinos with energies $E_\nu > 10$ TeV and zenith angle up to $60^\circ$ can be vetoed with an efficiency of $> 99\%$. Practical realization will depend on the depth of the neutrino telescope, on the muon veto efficiency and on the ability to identify downward moving neutrinos with a good energy estimation.
Neutrino telescopes such as IceCube at the South Pole, ANTARES in the Mediterranean Sea and the Lake Baikal detector search for neutrinos of extraterrestrial origin by using the Earth as a filter to suppress the background of atmospheric muons. The charged current $\nu_\mu$ channel is favored at TeV to PeV energies because of the large muon range, up to several kilometers of matter. To prevent downward-going atmospheric muons ($\mu^-$) from being mis-identified as muons induced by neutrino interactions, neutrino telescopes typically search for neutrino-induced upward or horizontal muons ($\nu_\mu^+$-$\mu^+$) (Markov and Zheleznykh).

The principle of detecting neutrinos by looking for $\nu_\mu^+$ implies that the field of view of a neutrino telescope is limited to half of the sky i.e. the opposite hemisphere with respect to the geographical position of the detector. Moreover, with this approach, atmospheric neutrinos that penetrate through the entire Earth become an irreducible background for the search of extraterrestrial neutrinos. There are two standard approaches to separating extraterrestrial from atmospheric neutrinos in the $\nu_\mu^+$ sample. One is based on seeing an excess of events in a particular direction/time interval (point source search), while the other is based on the assumption that astrophysical neutrinos have a harder spectrum than the atmospheric $\nu_\mu^+$ background (diffuse search).

The new generation of neutrino telescopes like IceCube, to be completed in 2011, and the R&D project KM3NeT, will operate instrumented volumes of about one km$^3$. Among the new opportunities offered by such large detectors, one feature seems especially interesting to us: the use of part of the instrumented volume as an active veto for $\mu^-$. This opens the field of view of neutrino telescopes to the hemisphere above the detector.

At the energies involved in a neutrino telescope (TeV-PeV), the opening angle between the $\nu_\mu^-$ and the $\mu^-$ produced in an atmospheric meson decay is very small. This implies that an atmospheric $\nu_\mu^-$ has a certain probability to arrive in the detector accompanied by its partner $\mu^-$. In this letter, we discuss the conditions under which a $\mu^-$-veto will also veto atmospheric $\nu_\mu^-$'s. This opens the opportunity to suppress what it is considered so far the irreducible background in neutrino telescopes.

At high energy, neutrinos come predominantly from decay of charged kaons and pions. Charged pions decay with a probability of 99.99% through the reaction $\pi^\pm \rightarrow \mu^\pm + \nu_\mu(\bar{\nu}_\mu)$ and charged kaons with a probability of 63.4% via $K^\pm \rightarrow \mu^\pm + \nu_\mu(\bar{\nu}_\mu)$. There are small con-
tributions from other channels, but we start by analyzing the simplest and most important case of pion and kaon decay and discuss the other cases at the end.

![Diagram of two-body decay of the parent meson into muon and neutrino.]

FIG. 1: Two body decay of the parent meson into muon and neutrino. The left figure displays the back-to-back kinematics in the meson center of mass (cm) frame. The right figure shows the momenta after Lorentz transformation into the laboratory frame.

Momentum conservation requires that the neutrino and companion muon are emitted back-to-back in the center of mass (cm) frame. The energies of the muon and neutrino in this frame are

\[ E_{\mu}^{cm} \approx \frac{m_i}{2} (1 + r_i) \]
\[ E_{\nu}^{cm} \approx \frac{m_i}{2} (1 - r_i) = |p^{cm}| \]

(c = 1). Here \( r_i = m_{\mu}^2 / m_i^2 \) are the quadratic mass ratios, where the parent meson masses are \( m_i \) (\( i = \pi^{\pm}, K^{\pm} \)), the muon mass is \( m_\mu \) and we neglect the neutrino mass. Numerically, \( r_\pi = 0.573 \) and \( r_K = 0.046 \). Given the higher rest mass of the kaon with respect to that of the muon, the energy is quasi-equally shared between neutrino and muon for the kaon two-body decay, while for pion decay the energy balance is shifted more towards the muon. It should be noted that the muon energy is always larger than the neutrino energy in the cm-frame.

Figure 1 displays the kinematical relations in the cm-frame and after Lorentz-transformation into the laboratory (lab) frame. After Lorentz-transformation along the positive x-axis from the cm- to the lab-frame, the muon and neutrino energies are given by:

\[ E_{\nu} = \gamma E_{\nu}^{cm} + \beta \gamma p_{2\nu}^{cm} \]
\[ E_{\mu} = \gamma E_{\mu}^{cm} + \beta \gamma p_{2\mu}^{cm} \]

where \( \gamma \) and \( \beta \) are the Lorentz factor and speed of the parent pion or kaon.
Taking into account the back-to-back emission of neutrino and muon in the cm-frame,

\[ p_{x\nu}^{cm} = |p^{cm}| \cos \theta_\nu \quad \text{and} \]
\[ p_{x\mu}^{cm} = |p^{cm}| \cos(\theta_\nu - \pi) = -|p^{cm}| \cos \theta_\nu, \]

where \( \theta_\nu \) is the angle of the neutrino in the cm-frame of the parent meson relative to its direction in the lab-frame. In the approximation \( \beta \to 1 \), which is valid for meson energies above several GeV, we can then rewrite Eq. (2) as

\[ E_\nu = \gamma |p^{cm}| (1 + \cos \theta_\nu) \quad \text{and} \]
\[ E_\mu = \gamma |p^{cm}| \left( \frac{1 + r_i}{1 - r_i} - \cos \theta_\nu \right). \]

The extreme condition for a given parent energy \( E_i = E_\nu + E_\mu \) is obtained for \( \cos \theta_\nu = 1 \) for which \( E_{\mu,\text{min}} = r_i E_{\nu,\text{max}} / (1 - r_i) \). Thus, for any \( (E_i, E_\nu) \) the minimum energy of the companion muon is

\[ E_\mu \geq \left( \frac{r_i}{1 - r_i} \right) E_\nu \quad \text{and} \]
\[ E_i = E_\mu + E_\nu \geq E_\nu \left( \frac{1}{1 - r_i} \right) \]

with \( i = \pi^\pm, K^\pm \). The corresponding numerical values for pion and kaon decays are

\[ \pi^\pm: \quad E_\pi \geq 2.342 \cdot E_\nu \]
\[ K^\pm: \quad E_K \geq 1.048 \cdot E_\nu. \]

The corresponding numeric values for the muon–neutrino relations are \( E_\mu \geq 1.342 \cdot E_\nu \) for \( \pi^\pm \) decays, and \( E_\mu \geq 0.048 \cdot E_\nu \) for \( K^\pm \). It should be noted that - unlike the situation in the cm-frame - the companion muon energy can be as low as 5% of the neutrino energy, if the parent meson is a kaon. The difference between pion and kaon decays which is related to the quadratic mass ratios \( r_i \), has an important impact on the energy dependence of the veto efficiency as discussed below. In particular, in addition to satisfying Eq. (5) the muon must have enough energy to penetrate to the detector.

To veto \( \downarrow \nu_\mu \), the \( \downarrow \mu \) track must also be relatively near to the neutrino trajectory. From the values of the transformed angles \( \Theta_\mu, \Theta_\nu \) we obtain typical distances between the neutrino and companion muon tracks after 10 km path length of less than 1 m (0.1 m) for neutrino energies above 1 TeV (10 TeV) if the parent meson was a pion, and less than 10 m (1 m)
for kaons. High energy atmospheric neutrinos and their companion muons can therefore be treated as quasi-aligned given the typical granularity of optical modules in a neutrino telescope.

The production spectrum of atmospheric neutrinos from $\pi^\pm (K) \rightarrow \mu^\pm + \nu_\mu$ is obtained from the convolution of the parent meson decay probability with the spectrum of mesons and the phase space distribution of the neutrinos:

$$P_\nu(E_\nu, X) = \int_{E_{\pi,\text{min}}}^\infty \left[ \frac{B_{\pi \rightarrow \mu \nu}}{E (1 - r_\pi)} \right] \left\{ \frac{\epsilon_\pi}{E X \cos(\theta)} \right\} \Pi(E, X) dE + \int_{E_{K,\text{min}}}^\infty \left[ \frac{B_{K \rightarrow \mu \nu}}{E (1 - r_K)} \right] \left\{ \frac{\epsilon_K}{E X \cos(\theta)} \right\} K(E, X) dE.$$  

Here $P_\nu dE_\nu$ is the number of neutrinos ($\nu_\mu + \bar{\nu}_\mu$) with energy between $E_\nu$ and $E_\nu + dE_\nu$ produced per g/cm$^2$ along the direction defined by zenith angle $\theta$. $\Pi(E, X)$ and $K(E, X)$ are the differential energy spectrum of charged pions and kaons at slant depth $X$ (in g/cm$^2$), and $E$ is the energy of the parent pion or kaon. The factors in curly brackets are the decay probabilities of pions (kaons) at vertical depth $X \cos(\theta)$. The factors in square brackets give the branching ratios times the normalized distributions of neutrino energies. The decay distributions are isotropic in the parent rest frame (uniform in $\cos(\theta)$), so from Eq. 4 the distributions of neutrino energy are flat over the kinematically allowed regions of phase space: $0 \leq E_\nu \leq E_\pi (1 - r_\pi)$ and $0 \leq E_\nu \leq E_K (1 - r_K)$. The total differential intensity of neutrinos is obtained by integrating Eq. 8 over the whole atmosphere with $E_{\pi,\text{min}}$ and $E_{K,\text{min}}$ given by Eqs. 6 and 7.

The contributions from pions and kaons in Eq. 8 are identical in form, but the kinematics are significantly different because of the difference in mass ratios. The critical energies below which decay is favored over hadronic re-interactions are also different, $\epsilon_\pi \approx 115$ GeV and $\epsilon_K \approx 850$ GeV. As a consequence of these differences the contribution of pions decreases and kaon decay becomes the main source of atmospheric neutrinos at high energy. (See $\pi$ fraction in Fig. 2)

If we require a minimum muon energy at production so that the muon can penetrate to the detector at slant-depth $X$, then, in addition to Eq. 5

$$E \geq E_\nu + E_{\mu,\text{min}}(X)$$  

(9)
FIG. 2: Unaccompanied vertical atmospheric $\downarrow \nu_\mu$ flux (upper); probability of accompaniment (lower) for seven depths. The grey curve is the fraction of $\downarrow \nu_\mu$ from $\pi$ decay.

must also be satisfied. The latter condition governs until $E_\nu > E_{\nu,\text{min}}(X) \times (1 - r)/r$, which is $E_\nu > 0.75 \times E_{\nu,\text{min}}(X)$ for pion decay but $E_\nu > 20 \times E_{\nu,\text{min}}(X)$ for kaon decay. For sufficiently high neutrino energy, the accompanying muon is guaranteed at depth, but the asymptotic energy occurs later for the dominant kaon component.

Fig. 2 illustrates the behavior of the proposed atmospheric $\downarrow \nu_\mu$ veto as a function of depth at vertical incidence. The lower plot shows the probability that the partner muon reaches various depths as a function of $E_\nu$. In each case the early onset of veto for the $\pi \rightarrow \mu \nu$ component is indicated by the shoulder of the curve. The upper panel shows the remaining flux of atmospheric $\downarrow \nu_\mu$ after the $\downarrow \mu$ veto is applied. Fig. 3 shows how the veto probability depends on zenith angle at two different depths, 1.8 and 3.5 km.w.e. The first approximates the center of IceCube at 2 km in ice, while the second represents the center of a detector at the NEMO site, which is the deepest candidate location for Km3NeT.

For this illustration we use a simple energy-independent, average relation from MMC for the muon energy at production needed to reach depth $X$:

$$E_{\mu,\text{min}}(X) = 0.73 \text{ TeV} \times \{\exp[X/2.8 \text{ km.w.e.}] - 1\}.$$ 

The $\downarrow \mu$ veto probability decreases with depth for a given neutrino direction and with in-
increasing zenith angle at a fixed vertical detector depth. These effects are both the result of the increased muon energy needed to penetrate to the deep detector. The dominant kaon contribution reaches its asymptotic value about a factor 30 higher in energy than the pion component.

The values in Figs. 2 and 3 are obtained from an analytical approximation applicable for a power law primary cosmic-ray spectrum \([10]\). For pions

\[
\Pi(E, X) = e^{-\left(\frac{X}{\Lambda \pi}\right)} \frac{Z_{\pi N}}{\Lambda_N} N_0(E)
\]

\[
\times \int_0^X \exp \left[ \frac{X'}{\Lambda \pi} - \frac{X'}{\Lambda_N} \right] \left( \frac{X'}{X} \right)^{\epsilon_\pi/E \cos \theta} \mathrm{d}X',
\]

with a corresponding expression for kaons. \(N_0(E)\) is the power-law primary spectrum of nucleons evaluated at the pion energy. Cross coupling between kaon and pion channels has been neglected as well as production of anti-nucleons. \(Z_{\pi N}\) is the spectrum-weighted moment for pion production, \(\Lambda_N\) is the nucleon interaction length and \(\Lambda_N, \Lambda_\pi, \Lambda_K\) are attenuation lengths for nucleons, charged pions and charged kaons, respectively.

Eq. 10 gives the production spectrum of neutrinos as a function of slant depth \(X\) in the atmosphere. The total differential neutrino spectrum requires evaluating \(\int_0^{\text{ground}} \mathcal{P}_\nu(E, X) \mathrm{d}X\). High-energy muons originate high in the atmosphere, so it is a good approximation to ex-
tend the upper limit of this integral to infinity. The result for the contribution from charged pions is

$$\phi_{\nu,\pi}(E_\nu) = N_0(E_\nu) Z_{N\pi} \frac{\Delta z}{\lambda N} \xi(E_\nu) \times \int_{z_{\min}}^{\infty} \frac{dz}{z + \xi(E_\nu)} \right)$$

with a similar expression for kaons. Here $z = E/E_\nu$ and $z_{\min}$ is the greater of $1 + E_{\mu,\min}/E_\nu$ from Eq. 9 or $1/(1 - r_i)$ (Eq. 5). Also, $\xi_i(E) = \epsilon_i/E\cos(\theta)$.

The integral in Eq. 11 can be evaluated analytically in the limits of low $(E_\nu \cos(\theta) \ll \epsilon_{\pi})$ and high $(E_\nu \cos(\theta) \gg \epsilon_{\pi})$ energy. One can then combine the low and high energy limits into a single approximation,

$$\phi(E_\nu) \approx \frac{N_0(E_\nu)}{1 - Z_{NN}} \sum_{i=\pi,K} \left[ A_i \frac{1}{1 + B_i/\xi_i(E)} \right],$$

where

$$A_i = \frac{Z_{N,i}}{1 - r_i} \frac{1}{\gamma + 1} \left( z_{i,\min} \right)^{\gamma + 1}$$

and

$$B_i = z_{i,\min} \gamma + 2 \frac{\Lambda_i - \Lambda_N}{\gamma + 1} \ln(\Lambda_i/\Lambda_N).$$

The expression 12 with numerical values of the parameters from Ref. [10] is used to make Figs. 2 and 3.

The new opportunity to veto high energy atmospheric muons using part of a neutrino telescope [8], opens the possibility to suppress downward going atmospheric neutrinos. Extraterrestrial neutrinos, which are in general not accompanied by a muon, will not be discarded by the veto system. If a downward event starts inside a fiducial volume of the detector and has no activity in the outer veto region, and has sufficiently high energy that an atmospheric $\downarrow\nu_\mu$ would be accompanied by its partner muon, then the neutrino can be classed (with a certain probability) as being of extraterrestrial origin.

So far we have assumed that all neutrinos come from decay of charged kaons and pions. We now estimate the contribution of atmospheric $\downarrow\nu_\mu$ from minor channels and assess their effects on the veto probability. The $K_{\mu3}$ semi-leptonic decay has a branching ratio of $0.27$ ($K^0_L$) and $0.033$ ($K^\pm$). The final state is $\pi, \mu, \nu_\mu$, so the minimal kinematic configuration occurs as before when the muon is backward in the cm system of the decaying kaon, balanced by the forward moving neutrino and pion. In this case, however, the forward momentum is shared by two particles so a lower $E_\nu$ guarantees the presence of the partner muon for a $\downarrow\nu_\mu$. This channel therefore does not weaken the veto probability. At the other extreme are
the $\nu_{\mu}$ from muon decay, in which case there is no muon. The contribution is small at high energy and can be estimated as in Ref. [11]. For $E_{\nu} \sim 10$ TeV and zenith angles $< 60^\circ$, this contribution is less than one per mil.

Contributions of prompt neutrinos from charm are more complex and more difficult to assess. A model for charm production that nearly saturates existing limits [12] is the Re-combination Quark Parton Model (RQPM) [13]. The charm contribution to atmospheric $\nu_{\mu}$ is less than 10% for $E_{\nu} = 10$ TeV in this model and crosses over the conventional contribution just above 100 TeV. Relevant decay channels include $D^+ \rightarrow \bar{K}_0 \mu^+ \nu_{\mu}$ (7%), $D^+ \rightarrow K^- \pi^+ \mu^+ \nu_{\mu}$ (4%), and $\Lambda_c^+ \rightarrow \Lambda \mu^+ \nu_{\mu}$ (2%). Because the charmed mass is higher, the energy fraction carried by the muon can be even lower than in the case of the kaon. On the other hand, the final states involve several particles so the neutrino energy required to see the partner muon is relatively low as well. Given the branching ratios and the upper limits on charm contribution, unaccompanied prompt $\nu_{\mu}$ should weaken the veto probability by less than one percent.

We conclude therefore that it should be possible in principle to veto atmospheric neutrinos with energies in the multi-TeV range and zenith angles less than $60^\circ$ with an efficiency of 99% or somewhat better. This paves the way for a sensitive and new type of search of neutrinos from astronomical sources. The main limitation will be the extent to which real detectors can put a lower limit on the energy of a neutrino interaction that starts in the detector. Because of the steep spectrum of atmospheric neutrinos, most of the interactions will be near the nominal threshold, thus enhancing the smearing effect of fluctuations. Detailed simulations of individual detectors, including surface air shower arrays, or alternative veto systems at shallow depths, are needed to assess the veto capability in practice. Such simulations will also be able to include the case where the $\nu_{\mu}$ is vetoed by a muon on a different branch in the accompanying air shower, which enhances the veto probability to some extent at high energy.

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