NEUTRINOS AS A PROBE OF PHYSICS BEYOND THE STANDARD MODEL*

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Neutrino physics plays a very important role in understanding several aspects of fundamental physics. In this short review, we revise three different aspects connected to them. Neutrino oscillations at low energy (up to few GeV) offer the possibility to study possible leptonic CP violation whereas at ultra-high energy ($E_\nu > 10^3$ TeV) they could enable to check scenarios of Physics Beyond the Standard Model involving extra-dimensions. On the other hand, e^+e^- colliders can help in detecting new heavy Majorana neutrinos suggested by Grand Unified Theories with masses of $\mathcal{O}(10^2\,\mathrm{GeV})$.

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1. Introduction

Neutrinos were initially suggested by Pauli to explain the apparent energy non conservation and wrong spin-statistics relations in nuclear beta decay and since then they have played a fundamental role in various branches of subatomic physics as well as in astrophysics and cosmology. In particular, the confirmation of the massive nature of neutrinos has important implications both in particle physics, where neutrino masses are likely related to a new mass scale not accessible by direct experimental study, and in astrophysics, where they give us important informations regarding the nuclear fusion in the stars and the supernova explosion mechanism. They would also help in explaining the apparent dark matter in the Universe.

In this context, the strong evidence for neutrino oscillations reported by the SuperKamiokande (SK) collaboration in their atmospheric neutrino data assumes a great importance; the deficit in the observed neutrino fluxes

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coming from the atmosphere compared to the expectations was measured in a high statistics experiment, showing a dependence on the neutrino path-length and energy in the way predicted in the case of neutrino oscillations. In more recent years the hypothesis of neutrino oscillations has been strongly confirmed in the atmospheric [1], accelerator [2], solar [3] and reactor [4] sectors. However, we do not have a complete information on the parameters of the neutrino mixing matrix and in particular no information whatsoever has been obtained on the existence of the complex phase of the matrix. With the aim of improving our knowledge of the mixing angles and to answer the fundamental question about the existence of the CP-violation in the leptonic sector, there has been a marked growth of interest in the development of new neutrino facilities; the role played in this context by a low γ β-beam is briefly reviewed in Sec. 2.

Oscillations take place only if neutrinos have non-degenerate masses; the data of neutrino experiments indicate that the neutrino masses are very small compared to the electroweak scale $v \sim \mathcal{O}(100\ \mathrm{GeV})$. This is usually explained invoking the *see-saw* mechanism in which the new physics is assumed to be a very high scale of $M \sim \mathcal{O}(10^{15}\ \mathrm{GeV})$ and its observable effects at energies below the TeV are the very light neutrino masses $m_\nu \sim (v^2/M) \sim \mathrm{eV}$. But independently of the mechanism generating them there may exist new heavy right-handed neutrinos with masses near the electroweak scale and mixings with the light leptons only constrained by experimental data. Large lepton colliders will be able to measure the properties of these heavy neutrinos; this will be the subject of Sec. 3.

The other aspect discussed here is related with scenarios of Physics Beyond the Standard Model (SM) involving large extra dimensions; their existence allow the possibility that the fundamental scale of gravity is at the TeV. If it is the case, gravity could dominate the interactions of ultra high energy cosmic rays making then observable the interaction of very energetic neutrinos of $E_\nu \sim \mathcal{O}(10^{10})$ GeV with nucleons at neutrino telescopes. Gravity-mediated elastic interactions at relatively large impact parameters can be calculated in the eikonal approximation; in Sec. 4 we will show that the calculated number of hadronic showers produced in ν-nucleon elastic interactions can be detected at IceCube to give a clear and model independent signature of the TeV gravity.

2. Low energy neutrinos

If neutrinos oscillate, the flavour and the mass eigenstates are not the same basis but they are related by a unitary 3×3 matrix U_{PMNS} [5]:

$$U_{\mathrm{PMNS}} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_{23} & s_{23} \\ 0 & -s_{23} & c_{23} \end{pmatrix} \begin{pmatrix} c_{13} & 0 & s_{13}e^{i\delta} \\ 0 & 1 & 0 \\ -s_{13}e^{-i\delta} & 0 & c_{13} \end{pmatrix} \begin{pmatrix} c_{12} & s_{12} & 0 \\ -s_{12} & c_{12} & 0 \\ 0 & 0 & 1 \end{pmatrix}, \quad (1)$$

where $c_{ij} = \cos\theta_{ij}$, $s_{ij} = \sin\theta_{ij}$ and δ is a complex phase. The probability to have an oscillation between two flavours α and β depends on the angles and the phase in Eq. (1) and also on the two independent mass differences $\Delta m^2_{12} = m^2_2 - m^2_1$ and $\Delta m^2_{23} = m^2_3 - m^2_2$. Oscillation experiments with low energy neutrinos have measured with quite good precision Δm^2_{12} and θ_{12} whereas we only know the absolute value of Δm^2_{23} and the allowed departure of θ_{23} from maximal mixing, $\sin^2 2\theta_{23} > 0.9$, but not if it is smaller or greater than 45°. Even if different experiments have been designed with the aim to perform precision measurements of the solar and atmospheric angles and mass differences, the ultimate and very hard to reach goal remains the measure of the two still unknown parameters of the PMNS mixing matrix, θ_{13} (for which only an upper limit exists so far [6]) and the leptonic CP violating phase δ, for which we have any information whatsoever. Two classes of problems arise if we take into account the possibility to measure these parameters at future neutrino facilities:

- The flavour changing transition $P^{\pm}_{\alpha\beta}$ obtained for neutrinos (+) and antineutrinos (−) at a fixed energy and baseline with input parameter $(\bar\theta_{13}, \bar\delta)$ has no unique solution. Indeed, the equation $P^{\pm}_{\alpha\beta}(\bar\theta_{13}, \bar\delta) = P^{\pm}_{\alpha\beta}(\theta_{13}, \delta)$ has two intersections: the input pair $(\bar\theta_{13}, \bar\delta)$ and a second, energy dependent, point. This second intersection introduces an ambiguity in the measurement of the physical values of θ_{13} and δ (*the ambiguity problem*).
- Neither in the solar nor in the atmospheric sector, mass differences and mixing angles are known with very high precision so it can be argued that the uncertainties on these quantities can strongly affect the measurements of θ_{13} and δ.

Many papers in the literature have been devoted to the study of the ambiguity problem and we address to them for an exhaustive review [7] of the subject. Here we will only concentrate on the impact of the atmospheric angle on the measurements on θ_{13} and δ in the case only the intrinsic ambiguity is present (a detailed analysis is in progress [8]) being the impact of the solar parameters estimated to be negligible.

All the previous studies devoted to the measurements of θ_{13} and δ at future neutrino facilities have been performed using a two-parameters χ^2 fixing the absolute magnitude of the mass differences and angles in the PMNS mixing matrix to some reasonable value (typically their best fit values). In order to take into account effects coming from the uncertainty on θ_{23}, we consider a three-variables $\chi^2(\theta_{13}, \delta, \theta_{23})$ in which the atmospheric angle is taken in the range corresponding to the usual 90% confidence level coming from the fits to the experimental data (we took as a reference the results claimed in [9]). The analytical structure of the fitting function is then:

$$\chi^2(\theta_{13},\delta,\theta_{23}) = \sum_{p=\pm}\left[\frac{N^p_\beta(\bar\theta_{13},\bar\delta,\bar\theta_{23}) - N^p_\beta(\theta_{13},\delta,\theta_{23})}{\sigma_{N^p}}\right]^2 , \qquad (2)$$

where N^+_β (N^-_β) indicates the number of positively (negatively) charged leptons produced by the neutrino of flavour β after interaction with the nucleons in the detector. The results of the fits are presented in the usual (θ_{13},δ) plane after projection of the three dimensional surface of χ^2 corresponding at the 90% CL. To illustrate the effect of such an analysis, we calculate the number of events in Eq. (2) considering a β-beam setup. The β-beam concept was first introduced in Ref. [10]. It involves producing a beam of β-unstable heavy ions, accelerating them to some reference energy, and allowing them to decay in the straight section of a storage ring, resulting in a very intense neutrino beam. The ions chosen have been ^{6}He, to produce a pure $\bar\nu_e$ beam, and ^{18}Ne (which has three different decay modes), to produce a ν_e beam. The neutrino beam energy depends on the γ of the parent ions in the decay ring. For the scenario considered in this paper the γ ratio for the two ions has been fixed to be $\gamma(^6\mathrm{He})/\gamma(^{18}\mathrm{Ne}) = 3/5$ [11]. More information about the neutrino fluxes (computed for a baseline of $L = 130\,\mathrm{Km}$ and $m_e \neq 0$) and the estimated background and efficiencies for the water detector adopted in this paper can be found in [12].

Let us first consider the sensitivity to the θ_{13} and δ parameters obtained looking for $\nu_e \to \nu_\mu$ ($\bar\nu_e \to \bar\nu_\mu$) oscillation, the signal being the appearance of ν_μ ($\bar\nu_\mu$) charged-currents in the detector. In Fig. 1 the result of the fit has been reported, in which the parameters have been fixed to the following values: $\Delta m^2_{12} = 8.2\times10^{-5}$ eV^2, $\theta_{12} = 32°$, $\Delta m^2_{23} = 2.5\times10^{-3}$ eV^2; the input point chosen to simulate the data is $\bar\theta_{13} = 7°$, $\bar\delta = 45°$ and $\bar\theta_{23} = 40°$. We have also superimposed the result of a usual two-parameters fit, performed at the same input value (but fixed $\theta_{23} = 40°$). As we can see, the error on the atmospheric parameter affects in a different way the measure of θ_{13} and δ: comparing the two and three-parameters fits, no large deviations in the allowed range of δ can be appreciated whereas the measure of θ_{13} appears strongly affected; this is mainly due to the fact that in the transition probability $P_{\nu_e\nu_\mu}$ θ_{23} multiplies the leading term containing θ_{13}. On the other hand the subleading dependence on δ is the reason for such a large vertical spread; in particular, the points on the plane close to 180° are the clone points whose location is analytically calculable as pointed out in [13].

This situation can be strongly improved if we also consider the informations coming from the disappearance channels that consists of the flavour conserving transition $\nu_e \to \nu_e$ ($\bar\nu_e \to \bar\nu_e$), the signal being the appearance of ν_e ($\bar\nu_e$) cherged-currents in the detector. Up to now, a true Monte Carlo

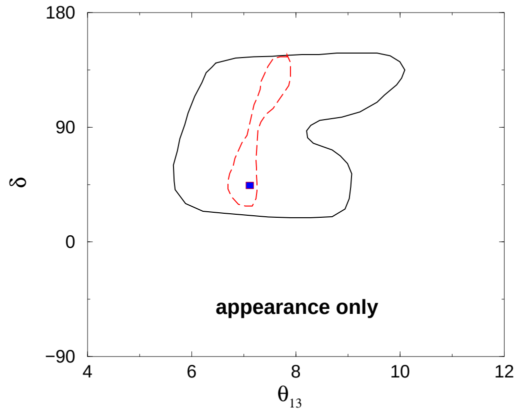

Fig. 1. Projection on the $(\theta_{13},\ \delta)$ plane of three-variables χ^2 fit after a 10 yrs of ν_e and $\bar{\nu}_e$ β-beam run with a 440 Kton water detector. The 90% CL contours are shown for the following input value: $\bar{\theta}_{13} = 7°$ and $\bar{\delta} = 45°$. Continuous line stands for the projection on the (θ_{13}, δ) plane of the three-parameters fit; dashed line stands for the result of a two-parameters fit. The box represents the input point.

simulation for estimating the backgrounds and systematics for the disappearance channels is missing, so that we can adopt the optimistic view of considering an overall systematic of 2% [14], fractional backgrounds at the level of 5×10^{-4} and efficiencies of 0.7 (for both ν_e and $\bar{\nu}_e$); the detector background is estimated to give 69 (11) events for ν_e ($\bar{\nu}_e$); these numbers are based on the discussions made in [12] and [15] . The effect of considering at the same time the appearance and disappearance channels is shown in Fig. 2 in which we plot the projection of the function $\chi^2(\theta_{13}, \delta, \theta_{23}) = \chi^2_{\rm app} + \chi^2_{\rm dis}$.

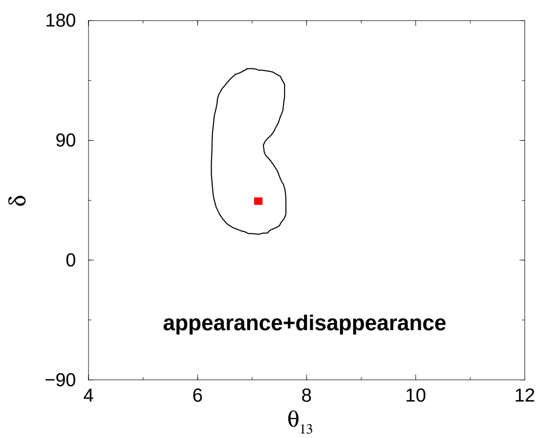

Fig. 2. The same as Fig. 1 but considering the appearance and disappearance channels at the same time.

It is clear that the disappearance channels help in cutting larger values of θ_{13}; as illustrated in [11] and [14], the mild dependence of $P_{\nu_e\nu_e}$ on θ_{13} permits a good sensitivity to this angle only for relatively larger values, depending on the systematic error taken into account. On the other hand, it does not depend on δ and θ_{23} so that the spread on the CP-violating phase remains the same.

3. Massive neutrinos at linear colliders

The phenomenon of neutrino oscillation involves very small neutrino masses of $\mathcal{O}$(eV); a simple way to explain them is to invoke the *see-saw* mechanism [16]. The SM Lagrangian is enlarged to include very heavy right-handed (RH) neutrinos $\nu_{\rm R}$ in such a way that the most general SU(3) $\times$ SU(2) $\times$ U(1) invariant Lagrangian giving rise to Dirac and ν^c Majorana masses is given by:

$$\mathcal{L} = -\,\nu^{cT} y_\nu (Hl) + \frac{1}{2}\nu^{cT} M \nu^c + \text{ h.c.}\,, \tag{3}$$

where H is the SM Higgs and l the lepton doublets. We expect the eigenvalues of M to be of order $M_{\rm GUT}$ or the new physics scale because ν^c Majorana masses are SU(3)$\times$ SU(2)$\times$ U(1) invariant, hence naturally of the order of the cutoff of the low-energy theory. Then, after integrating out the new heavy fields the lowest order terms of the corresponding effective Lagrangian reduce to the SM Lagrangian of dimension 4 plus a unique operator of dimension 5 violating lepton number which after spontaneous symmetry breaking gives masses to the light neutrinos (see for a review [17])

$$m_\nu = m_D^T M^{-1} m_D\,. \tag{4}$$

Automatically the mixing angles are also small since they are suppressed by a factor M^{-1}.

Independently of the mechanism generating the light neutrino masses there may exist new heavy RH neutrinos, singlet under the SM gauge group, with masses near the electroweak scale and relatively large mixings only constrained by present experimental limits. They complicate the explanation of why the observed neutrino masses are so small; however, we can consider some internal symmetry in the neutrino mass matrix by means of which we can recover an explanation of the small neutrino masses but completely disconnected from the values of the mixing angles.

Here we take the phenomenological approach of parametrising the corresponding enlarged Lagrangian, allowing for the largest mass and mixing values permitted by experiments, and investigate what we can learn at future

e^+e^- colliders without wondering about the detailed neutrino mass generation mechanism. If we enlarge the SM Lagrangian introducing n heavy Majorana neutrino singlets N, the corresponding neutrino mass matrix has dimension $3+n$, and will be diagonalised by a unitary matrix

$$U_\nu^T M U_\nu = (M_\nu)_{\rm diag} \equiv {\rm diag}(m_1 m_2 m_3 M_1 ... M_n), \tag{5}$$

where

$$U_\nu = \begin{pmatrix} \mathcal{U} & \mathcal{V} \\ \mathcal{V}' & \mathcal{U}' \end{pmatrix}, \tag{6}$$

with the 3×3 matrix $\mathcal{U}$ ($n\times n$ matrix $\mathcal{U}'$) describing the mixing among the light (heavy) neutrinos and the matrices $\mathcal{V}$ and $\mathcal{V}'$ parametrising the mixing between the light and heavy neutrinos. Thus, the flavour field eigenstates are linear combinations of the mass field eigenstates in such a way that, inserting them in the SM interaction Lagrangian we obtain:

$$\begin{aligned} L_{\rm int} &= -\frac{e}{2\sqrt{2}\sin\theta_W} \sum_{\alpha=e,\mu,\tau} \sum_{i=1}^{n} \bar{l}_\alpha \gamma^\mu (1-\gamma_5) (U_\nu)_{\alpha i} \nu_i W_\mu^- + {\rm h.c.} \\ &\quad - \frac{e}{4\sin\theta_W \cos\theta_W} \sum_{\alpha=e,\mu,\tau} \sum_{i,j=1}^{n} \bar{\nu}_i \gamma^\mu (1-\gamma_5) (U_\nu)^*_{\alpha i} (U_\nu)_{\alpha j} \nu_j \, Z_\mu \,. \end{aligned} \tag{7}$$

Bounds on the new interactions come form universality and by the non-observation of the lepton number violating processes $\mu \to e\gamma, \mu \to ee\bar{e}, Z \to e\bar{\mu}, ...$ [18].

For the sake of simplicity, let us consider the case in which only one Majorana neutrino N is added; then, the limits quoted in [18] can be read as bounds of the type: $\sin^2\theta_{\nu_e N} < 0.0054$, $\sin^2\theta_{\nu_\mu N} < 0.0096$ and $\sin^2\theta_{\nu_\tau N} < 0.016$. These mixings can be studied at $e^+ e^-$ colliders considering the process $e^+e^- \to l\nu W$ [19, 20]. Choosing the flavour of the final charged lepton l, we can select one (or a combination of) mixing angle(s):

$$\begin{aligned} e^+e^- &\to e^+\nu_e W^- \to \theta_{\nu_e N}\,, \\ e^+e^- &\to \mu^+\nu_\mu W^- \to \theta_{\nu_e N},\ \theta_{\nu_\mu N}\,, \\ e^+e^- &\to \tau^+\nu_\tau W^- \to \theta_{\nu_e N},\ \theta_{\nu_\tau N}\,. \end{aligned}$$

On the other hand, the beam polarisation allows to pick up different contributions. If the heavy neutrino is a Majorana particle and it only couples to one lepton flavour, for example to the muon family, then for an initial state in which left-handed positrons and right-handed electrons collide, the contribution to the $e^+e^- \to \mu^+\nu_\mu W^-$ reduces to the diagrams in Fig. 3.

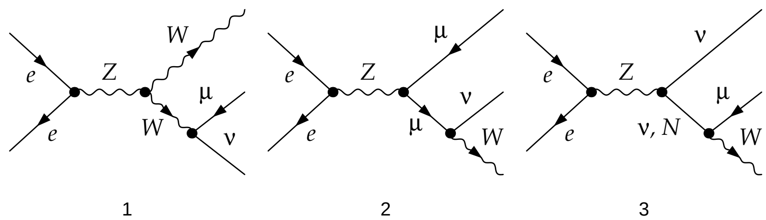

Fig. 3. Feynman diagrams contributing to the $e^+e^- \to \mu^+\nu_\mu W^-$ process for a Majorana neutrino and initial state with left-handed positrons and right-handed electrons.

For this choice of the polarisations of the initial leptons, we can strongly reduce the SM background, which in this case is given by the diagrams 1 and 2. Moreover, the diagram 3 with the heavy neutrino N is the signal of our process and only depends on the mixing angle $\theta_{\nu_\mu N}$ (the diagrams containing $\theta_{\nu_e N}$, which would have to be added if N would also mix with the electron, do not contribute due to the chirality structure of the amplitudes). A detailed analysis of the physics potential of a lepton collider in revealing and then studying the properties of such heavy neutrinos will be presented elsewhere [20].

4. Ultra-high energy neutrinos

When a cosmic ray from outer space hits a nucleon in the upper atmosphere it produces extensive air showers. The observed events [21] have energies of up to 10^{11} GeV, and their profile and distribution are consistent with a primary proton of extragalactic origin. In their way to the Earth these protons interact with the CMB photons and produce pions:

$$p + \gamma_{2.7\mathrm{K}} \to \Delta^+ \to n + \pi^+ \; (p + \pi^0) \, . \tag{8}$$

The flux of cosmogenic neutrinos is created in the decay of the charged pions, and it will appear correlated with observable fluxes of nucleons and photons.

These neutrinos are of great interest as probes of new TeV physics first because they provide large center of mass energies and second because the relative effect of new physics on the weakly interacting neutrinos is larger than on quarks or charged leptons, making it easier to see deviations.

Here we concentrate on scenarios of Physics Beyond the SM involving large extra dimensions. In its simplest picture, the fundamental Planck scale and the number n of compact dimensions where gravity propagates are the only parameters. Making use of the Gauss's and Newton's laws it is possible to relate the D-dimensional Planck scale M_D with the 4-dimensional one M_P and the radius of the compact dimensions: $M_\mathrm{P}^2 \sim R^n \, M_D^{n+2}$. Then, for sufficiently large values of R, the order of magnitude of M_P can be recovered putting $M_D = \mathcal{O}(1 \text{ TeV})$. This opens the very interesting possibility that gravity could dominate the interaction of ultra-high energy cosmic rays. In the energy regime in which the center of mass energy $s > M_D^2$ at the parton level (*transplanckian regime*), a neutrino can experience two different types of interaction with nucleons: *(i)* inelastic processes where it interacts with a parton and forms a black-holes; *(ii)* elastic processes where it transfers to the parton a small fraction $y = (E_\nu - E_\nu^{'})/E_\nu$ of its energy and keeps going. The cross sections for the first class of processes can only be estimated via dimensional arguments ($\sigma_{\rm BH} = \pi R_{\rm S}^2$ where $R_{\rm S}$ is the Schwarzschild radius of the system) and it is affected by large uncertainties coming from factors like the angular momentum, the charge, the geometry of the trapped surface and the radiation before the collapse, that make a quantitative estimate difficult. On the other hand, elastic processes in which neutrinos transfer to the parton a small fraction of energy up to $y \sim \mathcal{O}(0.1)$ can be estimated with the *eikonal approximation*. The eikonal amplitude $\mathcal{A}_{\rm eik}(s,t)$ [22–24] re-sumes the infinite set of ladder and cross-ladder diagrams in which D-dimensional gravitons are exhanged among the neutrinos and the partons of the nucleons. The Born amplitude of the process is computed as an integral over the momentum k_T along the extra dimensions

$$\mathcal{M}_{\rm Born}(q_\perp^2) \;=\; -\frac{s^2}{M_D^{n+2}} \int d^n k_T \frac{1}{t - k_T^2} \,,$$

where $q_\perp^2 \sim -t$ is the transverse momentum of the gravitons. The one-loop amplitude $\mathcal{M}_{1-{\rm loop}}(q_\perp^2)$ is a convolution of two Born amplitudes and one can write:

$$\mathcal{M}_{\rm Born} + \mathcal{M}_{1-{\rm loop}} \;=\; -2is \int d^2 b_\perp e^{i q_\perp \cdot b_\perp} \left(i\chi - \frac{1}{2}\chi^2 \right) .$$

Summing to all orders one gets:

$$\mathcal{M}_{\rm eik} \;=\; = -2is \int d^2 b_\perp e^{i q_\perp \cdot b_\perp} \left(e^{i\chi} - 1 \right) \,,$$

in which $\boldsymbol{b}$ spans the bidimensional impact parameter space and χ is the eikonal phase:

$$\chi(b_\perp) \equiv \frac{1}{2s} \int \frac{d^2 q_\perp}{(2\pi)^2} e^{-i q_\perp \cdot b_\perp} \mathcal{M}_{\rm Born}(q_\perp^2).$$

When a cosmogenic neutrino with energy E_ν enters the atmosphere, the two types of processes described above can start hadronic showers of energy up to $0.8 E_\nu$, from thermal evaporation of the formed black-holes, and of energy

yE_ν, if elastic processes are at work. AGASA and Fly's Eye are able to efficiently detect penetrating air showers of energies above $\approx 10^{10}$ GeV [25]. In these experiments 1 event passes all the cuts, which implies [25] an upper bound of 3.5 neutrino events at 90% CL. Making use of the neutrino flux quoted in [26] and the cross sections of Fig. 4, we obtain 3.5 events (2.1 BH and 1.4 elastic) if $M_D = 1.0$ TeV, whereas for $n = 6$ we have 2.6 BH plus 0.9 eikonal events if $M_D = 1.5$ TeV.

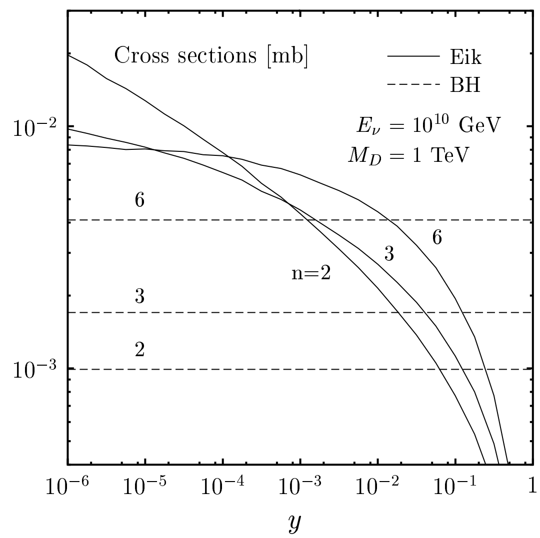

Fig. 4. Eikonal and BH cross sections as a function of the fraction of energy lost by the neutrinos.

Neutrino telescopes can do better in revealing TeV gravity signals from neutrino interactions. The IceCube experiment [27] seems to be the best place to look for them. It is a large scale (km^3) neutrino telescope currently under construction in the Antarctic ice. Its center is at a depth of 1.8 km, which implies that if $\sigma_{\nu N} \leq 0.01$ mb neutrinos can reach it vertically with no previous interactions, whereas if $\sigma_{\nu N} \leq 0.0001$ mb they could also reach it horizontally after crossing 150 km of ice. The detector is sensitive to hadronic showers of energy $E_{\rm sh} > 500$ TeV. Given a cosmogenic neutrino flux Φ_ν, the number $N_{\rm sh}$ of shower events at IceCube can be estimated as

$$N_{\rm sh} = \sum_i 2\pi AT \int d\cos\theta_z \int dE_\nu \frac{d\Phi_\nu}{dE_\nu} P_{\rm surv} P_{\rm int}\ , \tag{9}$$

where the sum goes over the three neutrino and antineutrino species, $A \approx 1\,{\rm km}^2$ is the detector's cross sectional area with respect to the ν flux, and T is the observation time. In the calculation we have to take into account two different aspects connected to the neutrino interaction:

- The probability that the neutrino survives to reach the detector from a zenith angle θ_z, $P_{\rm surv} = \exp[-X(\theta_z)\ \sigma\ N_A]$, where $X(\theta_z) \approx \rho_{\rm ice} L(\theta_z)$ is the column density of material ($L(\theta_z)$ is the length of the column in ice) in its way to the detector, and $\sigma = \sigma_{\rm BH} + \sigma_{\rm SM}$ is the inelastic cross section (we do not include $\sigma_{\rm eik}$ because the elastic processes with small y introduce a negligible distortion in the energy of the neutrinos that reach the detector).

- The probability that, once in the detector, the neutrino experiences an interaction: $P_{\rm int} = 1 - \exp[-L\ \rho_{\rm ice}\ \sigma\ N_A]$, where $L \approx 1\,{\rm km}$ is the linear dimension of the detector and σ is the cross section for the type of interaction process we are interested in.

If we consider values of M_D above the bounds obtained from the absence of penetrating air showers, a prediction on the number of events per year at IceCube can be made [28]. For $n = 2$ (6) we obtain a maximum of 118 (34) elastic events *versus* just 20 (24) short distance events; within the SM we expect 1.4 (0.5) hadronic or electromagnetic events (muons and taus do not shower) per year above 500 TeV. It means that, since the estimate of the eikonal cross section only involves linearised gravity and it is not affected by the uncertainties in the cross section for BH formation, we will face with *a clear and model independent signal of TeV gravity.* On the other hand, the absence of any signal at IceCube would imply a bound on M_D, as we can easily seen looking at Fig. 5. From the behaviour of the number of elastic events per year as a function of M_D we find that IceCube could detect TeV gravity effects above the SM background for M_D up to approximately 5 TeV.

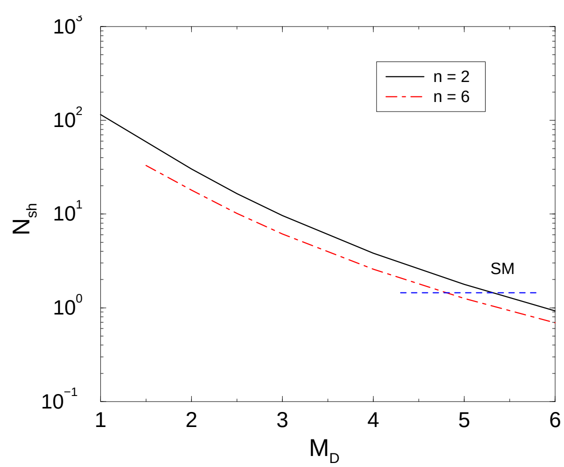

Fig. 5. Number of (eikonal and SM) shower events per year at IceCube for two different numbers of extra dimensions.

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