Cosmic Rays and New Physics at the TeV: the Neutrino-Nucleon Cross Section

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Abstract

Ultrahigh energy neutrinos can be used to explore the physics at the TeV scale. We study the neutrino-nucleon cross section in models with extra dimensions and the fundamental scale at the TeV. In particular, we discuss the production of string resonances and the gravitational interactions (multigraviton exchange and production of microscopic black holes) in these models. We show that the new TeV physics could give observable signals in horizontal air showers and neutrino telescopes.

1 Introduction

We observe extensive air showers produced when a cosmic ray from outer space hits a nucleon in the upper atmosphere. The energy of the particle starting the shower can be very large, with observed events of up to 10^{12} GeV. In particular, there are events above the so called GZK cutoff energy $E_{\rm GZK} \approx 5 \times 10^{19}$ eV. These events present a problem because, although the profile of the shower is consistent with a primary proton, the process

$$p + \gamma_{2.7{\rm K}} \to \Delta^+ \to p + \pi^0 \ (n + \pi^+) \tag{1}$$

is very effective on reducing the energy of a proton propagating in the cosmic background. Since in principle there are no near sources of such energetic particles, these cosmic rays should not be protons. One could then especulate if they are neutrinos. Neutrinos can come from a very far source with no loss of energy. The problem with them is that their standard cross section (c.s.) with the atmosphere is five orders of magnitude too small. A first motivation to study the ν-N c.s. is then: Can new TeV physics increase this c.s. up to hadronic size? In addition, there are experiments designed to measure the standard model (SM) ν-N c.s. at ultrahigh energies. A second motivation would be: Can new TeV physics give observable signals in horizontal air showers or neutrino telescopes? Note that since they are weakly interacting the relative effect of new physics on neutrino interactions is going to be much larger than in proton or charged lepton interactions.

Let us make a naive estimate of the impact of new physics on the cross section. The ν-q amplitude is mediated in the SM by a Z boson in the t channel: $\mathcal{A}_{\rm SM}(s,t) \propto g^2 s/(t - M_Z^2)$, where s and t are the usual Mandelstam variables ($t = -q^2 = -\frac{1}{2}(1-\cos\theta)$). In the limit $s \gg M_Z^2$ the c.s.

$$\sigma_{\rm SM}(s) = \frac{1}{16\pi}\frac{1}{s^2}\int_{-s}^{0} {\rm d}t\ |\mathcal{A}_{\rm SM}|^2 \propto \frac{g^4}{M_Z^2} - \frac{g^4}{s + M_Z^2} \tag{2}$$

is dominated by t of order M_Z^2 and becomes just $\sigma_{\rm SM}(s) \propto g^4/M_Z^2$. This simple result tells us that even at large center of mass energies the c.s. is very sensitive to the mass of the exchanged particle. It explains why the c.s. would be much larger if the exchanged particle were a gluon, and also that a Z' boson would introduce only a small correction to the SM result.

The presence of extra dimensions would have a more promising impact. A graviton-mediated Born amplitude will present two main features. First, the spin 2 of the intermediate field gives amplitudes growing like s^2, versus just s for the spin 1 Z boson. Second, one has to sum the contributions of the infinite tower of KK gravitons. Actually, this gives a divergence for more than one extra dimension. Although both effects push the cross section in the right direction, they both imply the presence of an ultraviolet cutoff Λ above which the model is not consistent. It is then easy to conclude that at the cutoff the c.s. $\sigma_{4+n} \approx 1/\Lambda^2$ does not seem large enough.

String theory provides another scenario with an infinite tower of higher spin fields (the string excitations) which, in addition, does not require an ultraviolet cutoff. In the brane world picture matter and gauge fields correspond to the zero modes of open strings, whereas the graviton is the massless mode of a closed string. A four fermion amplitude will include diagrams exchanging open and closed strings, with the later subleading in the string coupling g (the exchange of a closed string can be also seen as the one-loop exchange of two open strings). We would expect:
(i) At $s \lesssim M_S^2$ the diagram with exchange of an open string dominates the amplitude, giving at $s \approx 0$ a Z boson in the t channel. The diagram with an intermediate closed string gives at low energy the gravitational interactions.
(ii) At $s \approx M_S^2$ both types of diagrams give string resonances (Regge excitations). Closed string excitations, however, couple weaker ($\sim g^2$) to fermions.
(iii) At $s \gg M_S^2$ both diagrams give the usual soft behaviour of the string in the ultraviolet: the amplitudes go to zero exponentially at fixed angle (t/s fixed) and like a power law in the Regge limit (t fixed). Essentially, the string amplitude goes to zero everywhere except forward ($-t \lesssim M_S^2$), where only survives the contribution of the massles mode of the intermediate string. In

this regime the exchange of open and closed string gives, respectively,

$$\mathcal{A}(s,t) \approx g^2 \frac{s}{t} \quad \text{and} \quad \mathcal{A}(s,t) \approx g^4 \frac{s^2}{M_S^{n+2}} \int \frac{\mathrm{d}^n\, q_T}{t - q_T^2}\ . \tag{3}$$

Although the amplitude mediated by a closed string is subleading in g, it grows faster with s. As the center of mass energy increases the string amplitude is dominated by the long distance (small t) contributions of the higher dimensional graviton.

In the next sections we will evaluate the contributions to the ν-N c.s. from the production of open string excitations and from graviton exchange. We will go beyond the Born level and use the eikonal approximation to resumate the dominant long distance graviton contributions.

2 String excitations

We will consider a simple brane model where matter and gauge fields are the zero modes of open strings with both ends attached to a set of N 4-dimensional branes sitting at a fixed point of a higher dimensional bulk. We will asume that an orbifold projection eliminates the extra symmetry of the massless modes leaving just the SM.

The four fermion amplitude corresponding to the exchange of an open string is then very simple:

$$\begin{aligned} \mathcal{A}(1,2,3,4) &= g^2 \mathcal{S}(s,t) F^{1243}(s,t,u) \mathrm{Tr}[t^1 t^2 t^4 t^3 + t^3 t^4 t^2 t^1] \\ &+ g^2 \mathcal{S}(s,u) F^{1234}(s,u,t) \mathrm{Tr}[t^1 t^2 t^3 t^4 + t^4 t^3 t^2 t^1] \\ &+ g^2 \mathcal{S}(t,u) F^{1324}(t,u,s) \mathrm{Tr}[t^1 t^3 t^2 t^4 + t^4 t^2 t^3 t^1] \end{aligned} \tag{4}$$

In this expression g is the gauge coupling, $\mathcal{S}(s,t) = \Gamma(1-\alpha' s)\Gamma(1-\alpha' t)/\Gamma(1-\alpha' s - \alpha' t)$ with $\alpha' = M_S^{-2}$ is the Veneziano factor (M_S is the string scale), the factors F^{abcd} depend on the helicity of the external fermions, and the Chan-Paton traces describe the gauge numbers of the fermions (t^a are representation matrices of $U(N)$).

Let us consider the process $\nu_L u_L \to \nu_L u_L$. We find that if the Chan-Paton traces satisfy $T_{1243} - T_{1324} = -1/10$ and $T_{1234} = T_{1324} \equiv -a/10$ then at low energies the string amplitude reproduces the SM result: a massless Z boson in the t channel with $s_W^2 = 3/8$. In particular, for $a = 0$ our amplitude is

$$\mathcal{A}(s,t) = \frac{2}{5} g^2 \frac{s}{t} \mathcal{S}(s,t) \ . \tag{5}$$

This amplitude has poles at $s = nM_S^2$; near the n pole it is

$$\mathcal{A}_n \approx \frac{2}{5} \frac{g^2}{s - nM_S^2} \frac{nM_S^4}{t} \frac{(t/M_S^2)(t/M_S^2+1)\cdots(t/M_S^2+n-1)}{(n-1)!} \tag{6}$$

The residue is a polinomial of order $n-1$ in t, indicating that the spin of the resonances in this mass level goes up to $J = n-1$. To separate the contribution of each resonance we express the amplitude in terms of the scattering angle and parametrice it in terms of rotation matrix elements:

$$\mathcal{A}_n = \frac{2}{5} g^2 \ \frac{nM_S^2}{s - nM_S^2} \ \sum_{J=0}^{n-1} \alpha_n^J \ d_{00}^J(\theta) \ . \tag{7}$$

The coefficient α_n^J gives the contribution to $\mathcal{A}(\nu_L, u_L \to \nu_L, u_L)$ of the resonance with mass $\sqrt{n} M_S$ and spin J. For example, at the first mass level we obtain a scalar resonance with $\alpha_1^0 = 1$, at $s = 2M_S^2$ there is a single vector resonance with $\alpha_2^1 = 1$, whereas at $s = 3M_S^2$ there are modes of spin $J = 2$ ($\alpha_3^2 = 3/4$) and $J = 0$ ($\alpha_3^0 = 1/4$). We find an interesting sum rule for the α_n^J coefficients:

Using the narrow-width approximation we find

$$\sigma(\nu_L u_L \to X_n^J \to \text{all}) \equiv \sigma_n^J(\nu_L u_L) = \frac{4\pi^2\Gamma_n^J}{\sqrt{n}M_S}(2J+1)\delta(\hat{s}-nM_S^2)\ , \tag{8}$$

where $\Gamma_n^J = (g^2|\alpha_n^J|\sqrt{n}M_S)/(40\pi(2J+1))$ is the partial width of the spin J resonance at the n mass level. The c.s. to create in the collision any resonance at the n mass level is independent of n:

$$\sigma_n(\nu_L u_L) = \sum_{J=0}^{n-1}\sigma_n^J = \frac{\pi g^2}{4}\ \frac{2}{5}\ \delta(\hat{s}-nM_S^2)\ . \tag{9}$$

We repeat the procedure for the rest of partons in the nucleon and plot in Fig. 1 the total ν-N c.s. for a particular choice of Chan-Paton traces.

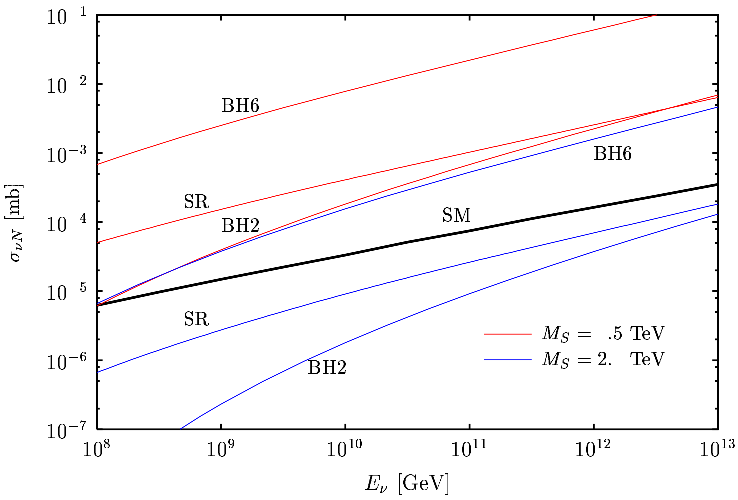

Figure 1: Cross section to produce string excitations (SR) and black holes in 2 or 6 extra dimensions (BH2, BH6). We include the SM cross section.

3 Graviton exchange: eikonal approximation

Although matter and gauge fields are attached to a 4-dimensional brane, gravity propagates in the 10 dimensions. We will asume that n of the 6 extra compact dimensions are large and gravity is D-dimensional ($D = 4+n$) at the distances of interest. The higher dimensional Newton's constant is usually defined as

$$G_D = V_n G_N \equiv \frac{(2\pi)^{n-1}}{M_D^{n+2}} , \tag{10}$$

where V_n is the volume of the compact space, G_N is the 4-dimensional constant, and M_D the higher dimensional Planck scale. Obviously, M_D is related to the string scale M_S. From the low-energy limit of the closed string amplitude we obtain

$$G_N = \frac{g^4}{64\pi^2} \frac{1}{M_S^{n+2} R^n} \quad \Rightarrow \quad M_D^{n+2} = \frac{1}{2\pi\alpha^2} M_S^{n+2} . \tag{11}$$

For $n = 2$ $M_D = 3.5 M_S$, whereas for $n = 6$ $M_D = 1.9 M_S$.

As explained in the introduction, at transplanckian energies the ν-q amplitude will be dominated by forward (long distance) graviton-mediated contributions. This is precisely the regime where we can use the eikonal approximation, that resumates all the ladder and cross-ladder contributions. Essentially, it is the exponentiation of the Born amplitude in impact parameter space:

$$\begin{aligned} \mathcal{A}_{eik}(s,t) &= \frac{2s}{i} \int d^2b \, e^{i\mathbf{q}\cdot\mathbf{b}} \left(e^{i\chi(s,b)} - 1 \right) \\ &= \frac{4\pi s}{i} \int db \, b J_0(bq) \left(e^{i\chi(s,b)} - 1 \right) , \end{aligned} \tag{12}$$

where $\chi(s,b)$ is the eikonal phase fixing the amplitude and $\mathbf{b}$ spans the (2-dimensional) impact parameter space. The Born amplitude corresponds to the limit of small $\chi(s,b)$:

$$\mathcal{A}_{Born}(s,t) = \frac{4\pi s}{i} \int db \, b J_0(bq) \, i\chi(s,b) . \tag{13}$$

We can then deduce the eikonal phase from the Fourier transform to impact parameter space of the Born amplitude:

$$\chi(s,b) = \frac{i}{2s} \int \frac{d^2q}{(2\pi)^2} \, e^{i\mathbf{q}\cdot\mathbf{b}} \, i\mathcal{A}_{Born} . \tag{14}$$

In our case the $\mathcal{A}_{Born}$ comes from the exchange in the t channel of a higher dimensional graviton:

$$i\mathcal{A}_{Born} = i\frac{s^2}{M_D^{n+2}}\int \frac{\mathrm{d}^n\, q_T}{t - q_T^2} = i\pi^{n/2}\Gamma\left(1 - \frac{n}{2}\right)\frac{s^2}{M_D^{2+n}}(-t)^{\frac{n}{2}-1}\;, \qquad (15)$$

where the integral over momentum along the extra dimensional q_T (equivalent to the sum over KK modes) gives an ultraviolet contribution that we have regularized using dimensional regularization. The *magic* of the eikonal approximation is that it will be well defined although we obtain it from an ultraviolet dependent Born amplitude: the contributions from large q_T introduce corrections to the phase $\chi(s,b)$ only at small b ($\approx 1/q_T$), but this small b region (see Eq. (12)) gives no sizeable contribution to the eikonal amplitude.

We obtain that the eikonal phase is

$$\chi(s,b) = (b_c/b)^n\;, \quad b_c^n = \frac{(4\pi)^{\frac{n}{2}-1}}{2}\Gamma\left(\frac{n}{2}\right)\frac{s}{M_D^{2+n}}\;. \qquad (16)$$

We note that $\chi(s,b)$ introduces a new scale, $q_c = 1/b_c$, that sets the size of the total cross section (which is proportional to b_c^2).

$\mathcal{A}_{eik}(s,t)$ can be evaluated numerically from Eq. (12) for any values of s and t. We can also obtain approximate expressions in the limits of small and large $q \equiv \sqrt{-t}$. For $q \ll q_c$ we get an expression analogous to the Born amplitude but with an effective cutoff of $\mathcal{O}(q_c)$. For $q \gg q_c$ the integral over impact parameter space is dominated by a saddle poing at $b_s = b_c(qb_c/n)^{-1/(n+1)} \ll b_c$, resulting

$$\mathcal{A}_{eik}(s,q) = Z_n\left(\frac{s}{qM_D}\right)^{\frac{n+2}{n+1}}\;. \qquad (17)$$

We give in Fig. 2 the ν-N eikonal cross section in terms of the fraction of energy $y = (E_\nu - E'_\nu)/E_\nu$ lost by the incident neutrino. For low values of n the neutrino would interact with the atmosphere, but losing only a small fraction of its energy (the cross section could be very large but *soft*).

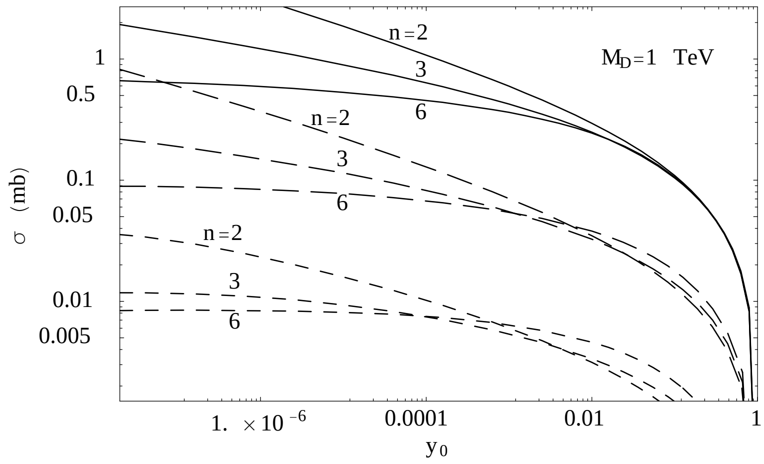

Figure 2: Elastic cross section vs. minimum fraction of energy lost by the neutrino for $E_\nu = (10^{14}, 10^{12}, 10^{10})$ GeV (solid, long dashes, short dashes).

4 Black hole production

The eikonal approximation provides an acceptable description of the scattering as far the dominant impact parameters (given by the position of the saddle poing b_s) are larger than the Schwarzschild radius R_S of the system:

$$R_S = \left(\frac{2^n \pi^{\frac{n-3}{2}} \Gamma\left(\frac{n+3}{2}\right)}{n+2} \right)^{\frac{1}{n+1}} \left(\frac{s}{M_D^{2n+4}} \right)^{\frac{1}{2(n+1)}} . \tag{18}$$

Diagramatically, at $b \lesssim R_S$ one finds that H diagrams (non-linear effects) become as important as the eikonal ladder diagrams. In this regime the system would collapse into a microscopic black hole with an approximate c.s.

$$\sigma(s) = \int_{M^2/s}^{1} dx \, \left(\sum_i f_i(x,\mu) \right) \, \pi R_S^2 \, . \tag{19}$$

This estimate would be reduced by graviton emision during the collapse but enhanced by the fact that a black hole acts as a somehow larger scatterer, so it should not be off by any large factors. Note also that the scale in the parton distribution functions must be $\mu = R_S^{-1}$ and not the black hole mass. The process becomes softer for larger center of mass energy and black hole mass. Actually, for cross sections of order $1/m_p^2$ the neutrino would not see partons at all, it would interact coherently with the whole proton. We give in Fig. 1 the c.s. for black hole production.

Comparing this process with the eikonal scattering we find that the energy transfer from an ultrahigh energy neutrino to the atmosphere would be dominated by hard processes:
(i) For a given flux of neutrinos of fixed energy, the total energy deposited in the atmosphere via small y scatterings will be smaller than through the (less frequent) events with large y or black hole formation.
(ii) In addition, for neutrino fluxes $J(E) \sim E^{\alpha}$ with $\alpha \lesssim -2$ a shower event of given energy would more likely come from a neutrino of similar energy than from a more energetic neutrino that lost only a fraction of its energy.

To conclude, the elastic exchange of higher dimensional gravitons or the production of string excitations or microscopic black holes can not explain the cosmic ray events above the GZK limit. Contrary to some claims, the last two processes are similar in size (see Fig. 1). Black hole production does not dominate (especially for low values of n) because the center of mass energy at the parton level is never too large: the increasing number of partons at small x favors the production of light black holes, with masses around the fundamental scale. All these processes, nevertheless, could give deviations to the signals expected in horizontal air showers and neutrino telescopes.

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