

Effects of charged Higgs bosons in the deep inelastic process $\nu_\tau \mathcal{N} \rightarrow \tau^- X$ and the possibility of detecting tau-neutrinos at cosmic neutrino detectors

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Abstract. We study the deep inelastic process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ (with $\mathcal{N} \equiv (n + p)/2$ an isoscalar nucleon), in the context of the two-Higgs doublet model Type II (2HDM(II)). We discuss the contribution to the total cross-section of diagrams, in which a charged Higgs boson is exchanged. We present results which show strong dependence of such contributions on $\tan\beta$ and M_{H^\pm} . We show that for $\tan\beta \approx 150$ and $M_{H^\pm} \approx 300$ GeV, the contribution of the charged Higgs boson exchange diagrams to the cross-section of the charged current inclusive $\nu_\tau \mathcal{N}$ collision can become important. We find that this contribution for an inclusive dispersion generated through the collision of an ultra-high-energy tau-neutrino with $E_\nu \approx 10^{21}$ eV on a target nucleon can be as large as 40% of the value of the contribution of the W^\pm exchange diagrams, provided $M_{H^\pm} \approx 300$ GeV and $\tan\beta \approx 150$. Such enhancement and the induced variation on the mean inelasticity $\langle y \rangle^{\text{CC}}$ could lead to sizeable effects in the acceptance of cosmic tau-neutrino detectors at experiments such as HiRes, PAO, and the CRTNT, which are anchored to the ground, and at experiments such as EUSO and OWL, which are proposed to orbit around the Earth. We also compare the contribution to $\sigma_{H^\pm}^{\text{tot}}$ from the different allowed initial quarks and we show that the contribution from the bottom quark dominates by far. This means that the H^\pm contribution practically always gives a top quark in the final state. Such a large component of the cross-section having a top quark event in the final state could have recognizable features in the EAS experiments.

Keywords. Charged Higgs bosons; ultra-high-energy neutrinos; tau-neutrinos.

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1. Introduction

Although the Standard Model (SM) [1], of the strong and electroweak interactions describes correctly particle physics at energy ranges currently attainable, one of its basic ingredients, the scalar Higgs sector, still remains untested. In the SM, the

Higgs sector consists of a single $SU(2)$ doublet, and after spontaneous symmetry breaking (SSB) it remains a physical state, the Higgs boson (h_{SM}^0), whose mass is not predicted in the theory. On the other hand, the SM is not expected to be the ultimate theoretical structure responsible for electroweak symmetry breaking (EWSB) [2,3]. One of the most simple extensions of the SM is the so-called two-Higgs doublet model (2HDM) [4]. There are four classes of 2HDM which naturally avoid tree-level flavor-changing neutral currents that can be induced by Higgs boson exchange [5]. These models include a Higgs sector with two scalar doublets, which give masses to the up- and down-type fermions as well as the gauge bosons. Model II is particularly interesting, where one of the Higgs scalar doublet couples to the up-components of isodoublets while the other to the down-components. Model II is that utilized in SUSY theories. The 2HDM(II) has a rich structure and predicts interesting phenomenology [2]. The physical spectrum consists of two neutral CP-even states (h^0, H^0) and one CP-odd (A^0), as well as a pair of charged scalar particles (H^\pm). The advantage of such a model is the fact that any Higgs sector built only upon doublets preserves naturally the lowest-order electroweak relation $\rho = 1$, with $\rho = M_{W^\pm}^2 / (M_Z^2 \cos^2 \theta)$, which has been tested with a good accuracy. On the phenomenological side, an important aspect of the 2HDM is that the Higgs sector may provide an additional source of CP violation [6].

Several experimental lower limits on the charged Higgs boson mass in this model have been reported in the literature:

$$M_{H^\pm} > 79.3 \text{ GeV (95\% C.L.) [7],} \tag{1}$$

$$M_{H^\pm} > (0.97, 1.28, 1.89) \times \tan \beta \text{ GeV (95\% C.L.) [8–10],} \tag{2}$$

$$M_{H^\pm} > (0.97, 1.5, 1.9, 2.5, 2.6) \times \tan \beta \text{ GeV (90\% C.L.) [11–15].} \tag{3}$$

There are also recent lower limits on the charged Higgs boson from results of $B^- \rightarrow \tau^- \bar{\nu}$ decay [16]. However, the most restrictive experimental lower limit on the charged Higgs boson mass in this model has been reported from the analysis of $\bar{B} \rightarrow X_s \gamma$ [17]:

$$M_{H^\pm} > 295 \text{ GeV (95\% C.L.).} \tag{4}$$

Further, based on the discussions on $\tan \beta$ given in refs [18,19], we restrict ourselves by taking the following upper limit on $\tan \beta$:

$$\tan \beta \leq 150. \tag{5}$$

Large-scale neutrino telescopes [20] have, as a main goal, the detection of ultra-high-energy (UHE) cosmic neutrinos ($E_\nu \geq 10^{12}$ eV) produced outside the atmosphere (neutrinos produced by galactic cosmic rays interacting with interstellar gas and extragalactic neutrinos) [21,22]. UHE neutrinos can be detected by observing long-range muons and tau-lepton decays produced in charged current neutrino–nucleon interactions. UHE tau-neutrinos are generated through neutrino oscillations [23,24]. The detection of UHE neutrinos will provide us with the possibility to observe $\nu\mathcal{N}$ -collisions with a neutrino energy in the range 10^{12} eV $\leq E_\nu \leq 10^{21}$ eV

and a target nucleon at rest. An enlightening discussion on UHE neutrino interactions is given by Gandhi *et al* [22].

We discuss in this paper the cross-section of the deep inelastic process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ ($\mathcal{N} \equiv (n + p)/2$ an isoscalar nucleon), in the context of the SM and of the 2HDM(II). We perform our numerical calculations using the QCD improved parton model [25,26] with the parton distribution functions reported by Pumplin *et al* [27]. We use the CTEQ PDFs provided in an $n_f = 5$ active flavors scheme. Our aim is to calculate how large is the contribution of diagrams, in which a charged Higgs boson is exchanged, to the total cross-section of the mentioned inclusive process in the frame of the 2HDM(II). In the 2HDM(II) the couplings of the down-type quarks and charged leptons are proportional to $m_f \times \tan\beta$. Hence, for large $\tan\beta$ the contribution of H^\pm -exchange diagrams will be maximal in this model.

This paper is organized as follows. In §2, we review the notation and physical region of the inclusive $\nu_l \mathcal{N}$ production. In §3, we present the formulae for the differential cross-section of the deep inelastic process $\nu_\tau \mathcal{N} \rightarrow \tau^- + X$, in the context of the SM and the 2HDM(II). In §4, we give our results and discuss the total cross-section rates of the charged current deep inelastic process $\nu_\tau \mathcal{N}$ in the frame of the SM and 2HDM(II). In §5, we discuss the contribution of charged Higgs boson to the charged current cross-section of the $\nu_\tau \mathcal{N}$ collision, the induced variation on the mean inelasticity $\langle y \rangle^{\text{CC}}$ and hence the effects on the possibility of detecting cosmic tau-neutrinos. Finally, in §6, we summarize our conclusions.

2. Notation, physical region

In this section we review the notation and the physical region for the parameters of the inclusive process

$$\nu_l + \mathcal{N} \rightarrow l^- + X, \quad (6)$$

where ν_l, \mathcal{N} and l^- stand for the incoming neutrino, the target nucleon and the outgoing lepton l^- , respectively. We will denote the four-momenta of these particles by $p, P_{\mathcal{N}}$ and p' , respectively. In accordance with the kinematics for the collision of a neutrino on a target nucleon, the following construction is chosen:

$$p^\mu = E_\nu(1, 0, 0, 1), \quad P_{\mathcal{N}}^\mu = M_{\mathcal{N}}(1, 0, 0, 0). \quad (7)$$

As usual, we define the invariants [26]

$$\begin{aligned} s &= (p + P_{\mathcal{N}})^2, \\ Q^2 &= -(p - p')^2, \\ \nu &= P_{\mathcal{N}}(p - p')/M_{\mathcal{N}}, \end{aligned} \quad (8)$$

and the dimensionless variables

$$x = \frac{Q^2}{2\nu M_{\mathcal{N}}}, \quad y = \frac{\nu}{E_\nu}. \quad (9)$$

The physical region of these kinematical variables is obtained by requiring that the scalar products of any two particle four-momenta be positive and the determinant

Δ_3 of the three independent four-momenta (whenever possible we will neglect the fermion masses)

$$\Delta_3(p, p', P_N) = \begin{vmatrix} 0 & pp' & pP_N \\ p'p & 0 & p'P_N \\ P_Np & P_Np' & 0 \end{vmatrix} \quad (10)$$

be positive [28].

From the non-negative character of the scalar products we find

$$0 \leq x \leq 1, \quad 0 \leq y \leq 1. \quad (11)$$

Explicit evaluation of Δ_3 using (8), (9) and (10) gives

$$\Delta_3 = (s/2)^3 2xy(1 - y), \quad (12)$$

with $s = 2M_N E_\nu$. The condition $\Delta_3 \geq 0$ does not lead to additional restrictions on the physical region.

The expressions given in (11), define the physical region for the dimensionless variables x and y . We have taken all fermion masses to be zero, which implies that in the calculation of the total cross-section the integration over the momentum transfer square extends up to zero. However, the parton distributions can be used only when Q^2 is not too small. Furthermore, in order to separate deep inelastic from elastic scattering, a cut on the invariant mass W of the unobserved particles in the final state is required. Therefore, besides the kinematical conditions (11), we also have in general the following constraints:

$$\begin{aligned} Q^2 &= sxy \geq Q_c^2, \\ W &= sy(1 - x) \geq W_c. \end{aligned} \quad (13)$$

The cuts for Q^2 and W constrain further the physically allowed region for the process (6). The physical region can now be written in terms of the dimensionless variables as follows:

$$\frac{Q_c^2}{s} \leq x \leq 1 - \frac{W_c}{s}, \quad \max \left\{ \frac{Q_c^2}{sx}, \frac{W_c}{s(1-x)} \right\} \leq y \leq 1. \quad (14)$$

The physically allowed region can also be expressed as

$$\frac{Q_c^2 + W_c}{s} \leq y \leq 1, \quad \frac{Q_c^2}{sy} \leq x \leq 1 - \frac{W_c}{sy}. \quad (15)$$

3. The cross-section for the inclusive process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$

3.1 The differential cross-section for the process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ in the SM

The differential cross-section for the inclusive reaction

$$\nu_\tau(p) + \mathcal{N}(P_N) \rightarrow \tau^-(p') + X, \quad (16)$$

where $\mathcal{N} \equiv (n + p)/2$ is an isoscalar nucleon, at the lowest order in α in the frame of the SM (see figure 1) is given as follows [29]:

$$\begin{aligned} \frac{d^2\sigma_{\text{SM}}}{dx dy} &= \frac{2G_{\text{F}}^2 M_{\mathcal{N}} E_\nu}{\pi} \left(\frac{M_{W^\pm}^2}{Q^2 + M_{W^\pm}^2} \right)^2 \\ &\times [xq_W(x, Q^2) + x\bar{q}_W(x, Q^2)(1 - y)^2], \end{aligned} \quad (17)$$

where Q^2 , x and y are defined in (8) and (9) and $M_{\mathcal{N}}$ stands for the nucleon mass. The quantities $q_W(x, Q^2)$ and $\bar{q}_W(x, Q^2)$ are given as

$$\begin{aligned} q_W(x, Q^2) &= \frac{u_v(x, Q^2) + d_v(x, Q^2)}{2} \\ &\quad + \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} + s_s(x, Q^2) + b_s(x, Q^2), \\ \bar{q}_W(x, Q^2) &= \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} + c_s(x, Q^2), \end{aligned} \quad (18)$$

where the valence and sea parton distribution functions (PDFs), $q_v(x, Q^2)$ and $q_s(x, Q^2)$, can be expressed as

$$\begin{aligned} u_v(x, Q^2) &= u(x, Q^2) - \bar{u}(x, Q^2), \\ d_v(x, Q^2) &= d(x, Q^2) - \bar{d}(x, Q^2), \\ u_s(x, Q^2) &= \bar{u}(x, Q^2), \\ d_s(x, Q^2) &= \bar{d}(x, Q^2), \\ c_s(x, Q^2) &= c(x, Q^2) = \bar{c}(x, Q^2), \\ s_s(x, Q^2) &= s(x, Q^2) = \bar{s}(x, Q^2), \\ b_s(x, Q^2) &= b(x, Q^2) = \bar{b}(x, Q^2), \end{aligned} \quad (19)$$

where the PDFs $q(x, Q^2)$ describe the quark q content of the proton. In other words, the parton distribution functions $q(x, Q^2)$ give the probabilities of finding a quark q inside a proton with the fraction x of the proton momentum: $q^\mu = xP^\mu$, in a scattering process with momentum transfer square Q^2 .

In the case of the Standard Model the couplings of the fermions to the W^\pm boson are given by the Lagrangian

$$\mathcal{L} = -\frac{g}{\sqrt{2}} \sum_{(f_u, f_d)} \left\{ \left(\bar{f}_u \gamma^\mu \frac{1 - \gamma_5}{2} f_d \right) W_\mu^+ + \left(\bar{f}_d \gamma^\mu \frac{1 - \gamma_5}{2} f_u \right) W_\mu^- \right\}, \quad (20)$$

where f_u and f_d stand for the up- and down-components of the fermion doublet.

3.2 The differential cross-section for the process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ in the 2HDM(II)

The differential cross section for the inclusive reaction (16), at the lowest order in α in the frame of the 2HDM(II) (see figure 2), can be written as follows [30]:

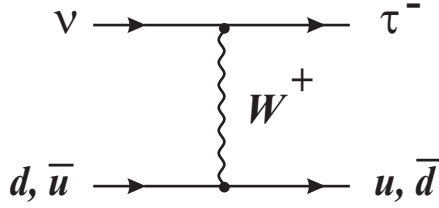


Figure 1. Diagrams which at the quark level contribute to the process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ at the lowest order in α in the SM (d stands for d -, s - and b -quarks; and u stands for u - and c -quarks).

$$\frac{d^2\sigma_{2\text{HDM}}}{dx dy} = \frac{d^2\sigma_{\text{SM}}}{dx dy} + \frac{d^2\sigma_{H^+}}{dx dy}, \tag{21}$$

where for large $\tan\beta$

$$\frac{d^2\sigma_{H^+}}{dx dy} = \frac{G_F^2 M_{\mathcal{N}} E_\nu m_\tau^2 M_{W^\pm}^2 \tan^4\beta}{2\pi (Q^2 + M_{H^\pm}^2)^2} y^2 [xq_H(x, Q^2) + x\bar{q}_H(x, Q^2)], \tag{22}$$

where Q^2 , x and y are defined in (8) and (9) and $M_{\mathcal{N}}$ stands for the nucleon mass. The quantities $q_W(x, Q^2)$ and $\bar{q}_W(x, Q^2)$ are given in (18), whereas $q_H(x, Q^2)$ and $\bar{q}_H(x, Q^2)$ are given as

$$\begin{aligned} q_H(x, Q^2) &= \frac{m_d^2}{M_{W^\pm}^2} \left(\frac{u_v(x, Q^2) + d_v(x, Q^2)}{2} + \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} \right) \\ &\quad + \frac{m_s^2}{M_{W^\pm}^2} s_s(x, Q^2) + \frac{m_b^2}{M_{W^\pm}^2} b_s(x, Q^2), \\ \bar{q}_H(x, Q^2) &= \frac{m_d^2}{M_{W^\pm}^2} \left(\frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} \right) + \frac{m_s^2}{M_{W^\pm}^2} c_s(x, Q^2). \end{aligned} \tag{23}$$

In the case of the 2HDM(II) the couplings of the fermions to the W^\pm boson are given by the Lagrangian in eq. (20), in a way similar to the SM [2]. On the other hand, taking the elements of the CKM-matrix $V_{ij} = \delta_{ij}$, the couplings of the fermions to the H^\pm boson are given by the Lagrangian [2]

$$\begin{aligned} \mathcal{L} &= \frac{g}{M_{W^\pm}} \left\{ m_\tau \tan\beta \left(\bar{\nu} \frac{1 + \gamma_5}{2} \tau \right) + m_u \cot\beta \left(\bar{u} \frac{1 - \gamma_5}{2} d \right) \right. \\ &\quad \left. + m_d \tan\beta \left(\bar{u} \frac{1 + \gamma_5}{2} d \right) \right\} H^+ + \text{h.c.} \end{aligned} \tag{24}$$

We end this section by pointing out that the interference effect of diagrams 2a and 2b can be neglected, because we are neglecting all fermion masses whenever possible (except those fermion masses coming from the Higgs boson–fermion couplings). We can make clear this statement by analyzing just the leptonic part of the interference term, as follows:

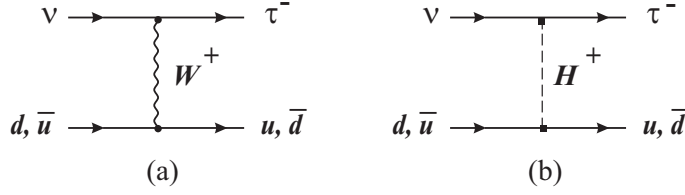


Figure 2. Diagrams which at the quark level contribute to the process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ at the lowest order in α in the 2HDM(II) (d stands for d -, s - and b -quarks, and u stands for u - and c -quarks).

$$\begin{aligned}
 & \sum_{\text{spins}} \text{Re} \{ \mathcal{M}_{2(a)} (\mathcal{M}_{2(b)})^* \}_{\text{leptonic part}} \\
 & \sim \sum_{\text{spins}} \bar{u}_\tau(p', s') \gamma^\alpha (1 - \gamma^5) u_\nu(p, s) \bar{u}_\nu(p, s) (1 + \gamma^5) u_\tau(p', s') \\
 & \sim \text{Tr} \{ \gamma^\alpha (1 - \gamma^5) \not{p} (1 + \gamma^5) \not{p}' \} \\
 & = 0.
 \end{aligned}$$

4. Results for deep inelastic $\nu_\tau \mathcal{N}$ in the SM and the 2HDM(II)

We present results for the case of unpolarized deep inelastic process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$ with a neutrino energy in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$ and the nucleon at rest, i.e. a target nucleon. We take $10^{15} \text{ eV} \leq E_\nu$ which leads to a condition where all fermion masses become negligible with respect to the total energy $s = 2M_{\mathcal{N}}E_\nu$, even the top quark mass. We take cuts of $\sim 2 \text{ GeV}^2$ and 10 GeV^2 for Q^2 and the invariant mass W , respectively. These values for the cuts are suited for the parton distribution functions of Pumplin *et al* [27] which we will use in our calculations. We have checked numerically that the total cross-section rates do not depend on the choice of the cuts on the momentum transfer square Q^2 , when they take on values of a few GeV^2 . This is due to the fact that the propagators involved in the calculation of the cross-section are $1/(M_{W^\pm}^2 + Q^2)$ and $1/(M_{H^\pm}^2 + Q^2)$.

We perform our numerical calculations for the quark masses by taking $m_u = 4 \text{ MeV}$, $m_d = 8 \text{ MeV}$, $m_c = 1.5 \text{ GeV}$, $m_s = 150 \text{ MeV}$, $m_b = 4.9 \text{ GeV}$ and $m_t = 174 \text{ GeV}$. For the evaluation of the $H^\pm \tau^- \nu_\tau$ coupling we take $m_{\nu_\tau} = 0$ and $m_\tau = 1.777 \text{ MeV}$. Also, we take $M_{W^\pm} = 80.4 \text{ GeV}$ for the mass of the charged boson W^\pm [31].

Taking into account the constraints on $M_{H^\pm}/\tan\beta$ given by (2) and (3), we will present results for $M_{H^\pm} \geq 2.0 \tan\beta \text{ GeV}$. The conditions (4) and (5) will also be fulfilled. Hence, we restrict our numerical analysis to the region $50 \leq \tan\beta \leq 150$ for $M_{H^\pm} = 300 \text{ GeV}$.

4.1 Total cross-section $\sigma_{\text{SM}}^{\text{tot}}$ and $\sigma_{\text{2HDM}}^{\text{tot}}$

We display in our numerical results the total cross-section as a function of E_ν in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$, with $E_{\mathcal{N}} = M_{\mathcal{N}}$ in figure 3. We observe in

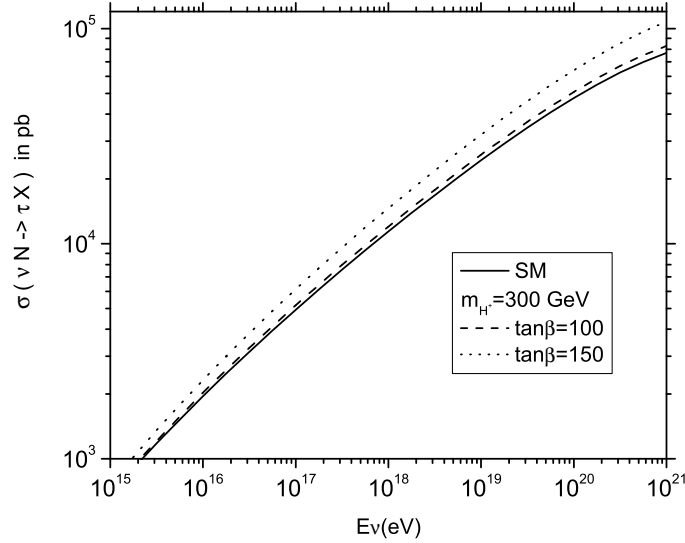


Figure 3. Comparison of $\sigma_{2\text{HDM}}^{\text{tot}}(\nu_\tau \mathcal{N} \rightarrow \tau^- + X)$ with $\sigma_{\text{SM}}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X)$ as a function of E_ν in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$, with $E_{\mathcal{N}} = m_{\mathcal{N}}$. We take for the 2HDM, $M_{H^\pm} = 300 \text{ GeV}$ and $\tan \beta = 100, 150$.

this figure that the total cross-section $\sigma_{2\text{HDM}}^{\text{tot}}$ becomes larger as $\tan \beta$ increases, as expected.

4.2 The ratio $\sigma_{H^\pm}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X) / \sigma_{\text{SM}}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X)$

In figure 4, we present our results for the ratio $\sigma_{H^\pm}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X) / \sigma_{\text{SM}}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X)$ as a function of E_ν in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$, with $E_{\mathcal{N}} = M_{\mathcal{N}}$. We observe in this figure that the ratio $\sigma_{H^\pm}^{\text{tot}} / \sigma_{\text{SM}}^{\text{tot}}$ becomes larger as $\tan \beta$ increases. Furthermore, we notice that $\sigma_{H^\pm}^{\text{tot}} / \sigma_{\text{SM}}^{\text{tot}}$ depends on E_ν . We see that for $50 \leq \tan \beta \leq 150$, the greater the E_ν , the larger the $\sigma_{H^\pm}^{\text{tot}} / \sigma_{\text{SM}}^{\text{tot}}$ ratio becomes.

4.3 Comparison of the contribution to $\sigma_{H^\pm}^{\text{tot}}$ from the different quarks

Finally, in figure 5 we compare the contribution to $\sigma_{H^\pm}^{\text{tot}}$ from the different allowed initial quarks (see figure 2b), taking $M_{H^\pm} = 300 \text{ GeV}$ and $\tan \beta = 150$. We observe in this figure that the contribution from the bottom quark dominates by far (a similar behavior is obtained in the other cases). This fact implies that the contribution of the H^\pm exchange diagrams to the total cross-section of the $\nu_\tau \mathcal{N}$ scattering in the frame of the 2HDM(II) is the same regardless of whether the nucleon is a proton, a neutron or an isoscalar nucleon, because these particles have the same content of b -quark. We would like to emphasize that the H^\pm contribution practically always gives a top quark in the final state, resulting in a much larger value of the mass

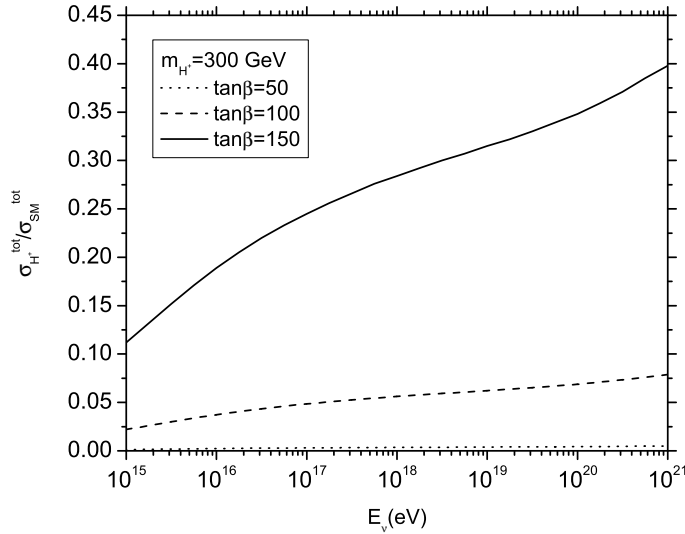


Figure 4. $\sigma_{H^+}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X)/\sigma_{\text{SM}}^{\text{tot}}(\nu_\tau + \mathcal{N} \rightarrow \tau^- + X)$ as a function of E_ν in the range 10^{15} eV $\leq E_\nu \leq 10^{21}$ eV, with $E_{\mathcal{N}} = m_{\mathcal{N}}$, for $M_{H^\pm} = 300$ GeV and $\tan\beta = 50, 100, 150$.

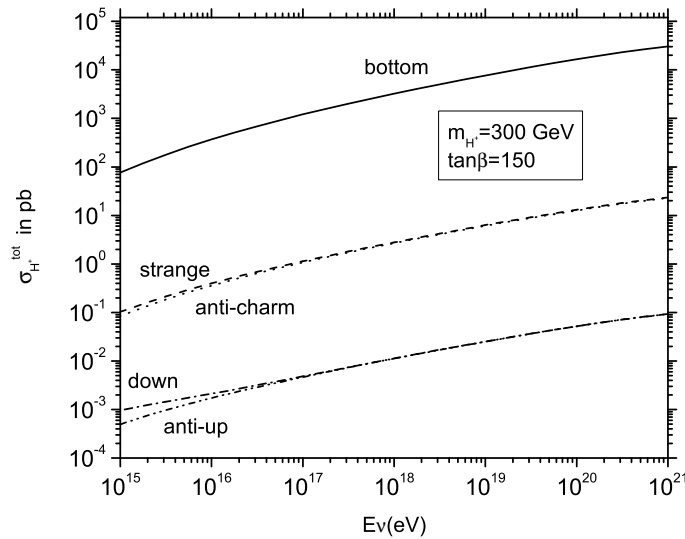


Figure 5. Contribution to $\sigma_{H^+}^{\text{tot}}$ from the different allowed initial quarks.

invariant W and dimensionless parameter y in the 2HDM with respect to the corresponding values in the case of the SM. Such a large component of the cross-section having a top quark event in the final state could have other recognizable features in the EAS experiments, which may be worth looking into.

5. Charged Higgs boson effects in the possibility of detecting ultra-high-energy tau-neutrinos

Cosmic neutrino fluxes can initiate air showers through interaction in the atmosphere, or in the Earth. Neutrino trajectories will be going down in a nearly horizontal way in the former case, whereas in an Earth-skimming manner in the latter case. Thus, it is important to know the acceptance (event rate/flux) of the proposed air shower experiments for detecting both types of neutrino-initiated events (horizontal air-showers (HAS) and upgoing (Earth-skimming) air-showers (UAS)). These acceptances for fluorescence detectors have been calculated in [32] for experiments as High Resolution Fly's Eye Detector (HiRes) [33], Pierre Auger Observatory (PAO) [34] and the Cosmic Ray Tau Neutrino Telescopes (CRTNT) [35], which are anchored to the ground, and at experiments such as Extreme Universe Space Observatory (EUSO) [36] and orbiting wide angle light-collectors (OWL) [37], which are proposed to orbit around the Earth. In ref. [32], it is stated that the very different dependence on the cross-section of the HAS (linear) and UAS (non linear) acceptances will provide a practical method to measure the charged current cross-section of the neutrino nucleon scattering $\sigma_{\nu\mathcal{N}}^{\text{CC}}$. One simply has to exploit the ratio of UAS to HAS event rates. However, one assumption made in ref. [32] was taking $\sigma_{\nu_e\mathcal{N}}^{\text{CC}} = \sigma_{\nu_\mu\mathcal{N}}^{\text{CC}} = \sigma_{\nu_\tau\mathcal{N}}^{\text{CC}}$. This is valid in the SM, but it is no longer valid in the 2HDM(II), because the $\sigma_{\nu_\tau\mathcal{N}}^{\text{CC}}$ could be twice larger than $\sigma_{\nu_e\mathcal{N}}^{\text{CC}} = \sigma_{\nu_\mu\mathcal{N}}^{\text{CC}}$ in this model. Hence, the neutrino charged current interaction mean-free path (introduced in ref. [32]) could be for ν_τ one half of that for ν_e and ν_μ , in the mentioned model. Therefore, the existence of charged Higgs bosons could imply a large deviation from the results for the acceptances for space-based (or ground-based) tau-neutrino detectors presented in figure 7 of ref. [32].

Further, it was already pointed out in ref. [38], that besides the obvious role of the neutrino cross-section in the actual interaction that leads to the possible neutrino detection there are other more subtle effects. The y distribution of the cross-section also has an impact on the detection rate [38–40]. It is clear from eqs (17) and (22) that the y distribution of the charged current interaction could change due to the effects of charged Higgs bosons. In fact, in table 1 we show how the mean inelasticity $\langle y \rangle^{\text{CC}}$ would change in the 2HDM(II) (we present only numerical results for Case I, in which the effects are maximal). Our results obtained for $\langle y \rangle^{\text{CC}}$ in the frame of the SM are in good agreement with those reported in ref. [41]. We observe in table 1 that $\langle y \rangle_{2\text{HDM}}^{\text{CC}} \cong 2.2 \times \langle y \rangle_{\text{SM}}^{\text{CC}}$ for $M_{H^\pm} = 300$ GeV, $\tan\beta = 150$ and $E_\nu = 10^{21}$ eV. Such a drastic variation in $\langle y \rangle^{\text{CC}}$ could induce a large variation in the possibility of detecting ultra-high-energy tau-neutrinos through inclined showers [38]. We end this section by stating that the possible effects of charged Higgs bosons on the detection of cosmic tau-neutrinos deserve future detailed investigations.

6. Conclusions

We have calculated the total cross-section rates for the deep inelastic process $\nu_\tau + \mathcal{N} \rightarrow \tau^- + X$, where $\mathcal{N} \equiv (n + p)/2$ is an isoscalar nucleon, in the frame of the SM

Table 1. Mean inelasticity parameter for charged current interaction $\langle y \rangle^{\text{CC}}$ as a function of E_ν , with $E_{\mathcal{N}} = M_{\mathcal{N}}$. We compare $\langle y \rangle_{\text{SM}}^{\text{CC}}$ with $\langle y \rangle_{\text{2HDM}}^{\text{CC}}$ by taking $M_{H^\pm} = 300$ GeV and $\tan \beta = 50, 100, 150$.

E_ν (eV)	$\langle y \rangle_{\text{SM}}^{\text{CC}}$	$\langle y \rangle_{\text{2HDM}}^{\text{CC}}$		
		$\tan \beta = 50$	$\tan \beta = 100$	$\tan \beta = 150$
10^{15}	0.289	0.290	0.299	0.335
10^{16}	0.255	0.255	0.271	0.329
10^{17}	0.234	0.235	0.256	0.328
10^{18}	0.218	0.221	0.245	0.327
10^{19}	0.204	0.207	0.234	0.324
10^{20}	0.180	0.184	0.215	0.315
10^{21}	0.131	0.135	0.173	0.290

and the 2HDM(II). In the case of the 2HDM(II) we have taken into account the contribution of the diagrams in which a charged Higgs boson ($\sigma_{H^\pm}^{\text{tot}}$) is exchanged.

We obtained our numerical results by considering the nucleon at rest and according to the following conditions: taking $M_{W^\pm} = 80.4$ GeV, $\sin^2 \theta_W = 0.223$, a neutrino energy in the range $10^{15} \text{ eV} \leq E_\nu \leq 10^{21} \text{ eV}$ and setting cuts of $\sim 2 \text{ GeV}^2$ and 10 GeV^2 for the momentum transfer square Q^2 and the invariant mass (W), respectively. We made use of the parton distribution functions of Pumplin *et al*, provided in an $n_f = 5$ active flavors scheme. In order to take into account the experimental data, we restricted our numerical analysis to the region $50 \leq \tan \beta \leq 150$ and $M_{H^\pm} = 300$ GeV, thus the constraint $M_{H^\pm} \geq 2.0 \times \tan \beta$ GeV is also fulfilled.

We have shown that the most important contribution to $\sigma_{H^\pm}^{\text{tot}}$ comes from the H^\pm -exchange diagram with an initial b -quark (and hence an outgoing t -quark). This fact implies that the contribution of the H^\pm exchange diagrams to the total cross-section of the $\nu_\tau \mathcal{N}$ scattering in the frame of the 2HDM(II) is the same regardless of whether the nucleon is a proton, a neutron or an isoscalar nucleon, because these particles have the same content of b -quark. On the other hand, we would like to stress that the H^\pm contribution leads to a top quark in the final state and that such a large component of the cross-section having a top quark event in the final state could have characteristic aspects in the EAS experiments.

We found that the ratio $\sigma_{H^\pm}^{\text{tot}}/\sigma_{\text{SM}}^{\text{tot}}$ becomes larger as $\tan \beta$ increases. We also found that the ratio $\sigma_{H^\pm}^{\text{tot}}/\sigma_{\text{SM}}^{\text{tot}}$ depends on E_ν . In particular, we obtain that for $50 \leq \tan \beta \leq 150$: the greater the E_ν , the larger the $\sigma_{H^\pm}^{\text{tot}}/\sigma_{\text{SM}}^{\text{tot}}$ ratio becomes.

We showed that the contribution of the charged Higgs boson exchange diagrams can lead to a sizeable enhancement with respect to the SM cross-section rates for the charged current $\nu_\tau \mathcal{N}$ deep inelastic scattering. For the case of an ultra-high-energy tau-neutrino with $E_\nu = 10^{21}$ eV colliding on a target nucleon such enhancement reached 40% provided that $\tan \beta = 150$ and $M_{H^\pm} = 300$ GeV. Besides, we obtain $\langle y \rangle_{\text{2HDM}}^{\text{CC}} \approx 2.2 \times \langle y \rangle_{\text{SM}}^{\text{CC}}$ for the same values of E_ν , $\tan \beta$ and M_{H^\pm} . This enhancement and the induced variation in the mean inelasticity $\langle y \rangle^{\text{CC}}$ could lead to sizeable effects in the acceptance of the cosmic tau-neutrino detectors at space-based experiments such as the EUSO and OWL proposals and at ground-based experiments such as PAO, HiRes and CRTNT.

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