Can one distinguish $\tau$-neutrinos from antineutrinos in neutral-current pion production processes?

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Received 10 August 2006; received in revised form 19 January 2007; accepted 7 February 2007
Available online 28 February 2007

Abstract

A potential way to distinguish $\tau$-neutrinos from antineutrinos, below the $\tau$-production threshold, but above the pion production one, is presented. It is based on the different behavior of the neutral-current pion production off the nucleon, depending on whether it is induced by neutrinos or antineutrinos. This procedure for distinguishing $\tau$-neutrinos from antineutrinos neither relies on any nuclear model, nor it is affected by any nuclear effect (distortion of the outgoing nucleon waves, etc.). We show that neutrino–antineutrino asymmetries occur both in the totally integrated cross sections and in the pion azimuthal differential distributions. To define the asymmetries for the latter distributions we just rely on Lorentz-invariance. All these asymmetries are independent of the lepton family and can be experimentally measured by using electron or muon neutrinos, due to the lepton family universality of the neutral-current neutrino interaction. Nevertheless and to estimate their size, we have also used the chiral model of [E. Hernández, J. Nieves, M. Valverde, hep-ph/0701149] at intermediate energies. Results are really significant since the differences between neutrino and antineutrino induced reactions are always large in all physical channels.

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PACS: 12.15.Mm; 25.30.Pt; 13.15.+g

1. Introduction

Distinguishing $\tau$-neutrinos from antineutrinos, can be easily done above the $\tau$-production threshold energy, since then the charged-current (CC) channel is open. Conservation of lepton number implies that neutrinos generate a tau lepton, while an antineutrino produces an antilepton. The charge of the tau lepton unambiguously reveals the nature of the incident neutrino beam. Below the $\tau$-threshold, $\tau$-neutrinos or antineutrinos only interact with matter via neutral-current (NC) driven processes, and the outgoing lepton is also a neutrino or antineutrino, and the difference does not manifest itself in this clear way.

Nonetheless, it is out of any doubt the great interest of distinguishing between neutrinos and antineutrinos in different scenarios: CP-violation in $\nu_\mu \rightarrow \nu_\tau$ oscillations [1], supernova-neutrino physics [2].

Very recently, Jachowicz and collaborators suggested that in neutrino–nucleus processes, the helicity-related differences between neutrino and antineutrinos induce some asymmetries in the polarization of the ejected nucleons [3]. From this fact, the authors of this reference conclude that these asymmetries represent a potential way to distinguish neutrinos from antineutrinos in NC neutrino scattering on nuclei.

We present in this Letter an alternative manner to distinguish $\tau$-neutrinos from antineutrinos, below the $\tau$-production threshold ($\approx 3.5$ GeV in the Laboratory (LAB) frame for production off the nucleon), but above the pion production one ($\approx 0.15$ GeV in...
the LAB frame). The method is based on the different behavior of the NC pion production reaction in the nucleon, depending on whether it is induced by neutrinos or antineutrinos. This procedure for distinguishing \( \tau \)-neutrinos from antineutrinos relies on any nuclear model, nor it is affected by any nuclear effect (distortion of the outgoing nucleon waves, etc.). We show that neutrino–antineutrino asymmetries occur both in the totally integrated cross sections and in the pion azimuthal differential distributions. To define the asymmetries for the latter distributions we just rely on Lorentz-invariance, and experimentally it requires determining the neutrino scattering plane, detecting either a neutral or a positively charged pion and measuring its momentum. Since the outgoing neutrino will not likely be detected, to fix the neutrino scattering plane would also require measuring the outgoing nucleon momentum and some knowledge of the incoming neutrino momentum.

All these asymmetries are independent of the lepton family and can be experimentally measured by using electron or muon neutrinos, due to the lepton family universality of the NC neutrino interaction. Nevertheless and to estimate their size, we have used the chiral model of Ref. [4] to compute them up to neutrino/antineutrino energies of 2 GeV.

2. Kinematics and cross sections

We will focus on the neutrino–pion production reaction off the nucleon driven by neutral-currents,

\[
v_l(k) + N(p) \rightarrow v_l(k') + N(p') + \pi(k_\pi)
\]

though the generalization of the obtained expressions to antineutrino induced reactions is straightforward.

The unpolarized differential cross section, with respect to the outgoing neutrino and pion kinematical variables is given in the LAB frame (kinematics is sketched in Fig. 1) by\(^2\)

\[
d\sigma_{\nu} = \frac{2G^2}{16\pi^2} \int_0^{+\infty} \frac{dk_\pi k_\pi^2}{E_{\pi}} L_{\mu\sigma}^{(v)} W_{NC}^{\mu\sigma}_{\pi}(v)
\]

with \( \bar{k} \) and \( \bar{k}' \) the LAB neutrino momenta, \( E' = |\bar{k}'| \), the energy of the outgoing neutrino, \( \bar{k}_\pi \) and \( E_\pi \), the momentum and the energy of the pion in the LAB system, \( G = 1.1664 \times 10^{-11} \mathrm{MeV}^{-2} \), the Fermi constant and \( L \) and \( W \) the leptonic and hadronic tensors, respectively. The leptonic tensor is given by (in our convention, we take \( \epsilon_{0123} = +1 \) and the metric \( g^{\mu\nu} = (+, -, -, -) \):

\[
L_{\mu\sigma}^{(v)} = (L_k^{(v)})_{\mu\sigma} + i(L_a^{(v)})_{\mu\sigma} = k'_\mu k_\sigma + k'_\sigma k_\mu - g_{\mu\sigma} k' \cdot k' + i\epsilon_{\mu\sigma\alpha\beta} k'^\alpha k'^\beta
\]

and it is orthogonal to \( q^\mu = (k - k')^\mu \) for massless neutrinos, i.e., \( L_{\mu\sigma}^{(v)} q^\mu = L_{\mu\sigma}^{(v)} q^\sigma = 0 \).

The hadronic tensor includes all sort of non-leptonic vertices and it reads

\[
W_{NC}^{\mu\sigma}(v) = \frac{1}{4M} \sum_{\text{spins}} \int \frac{d^3p'}{(2\pi)^3} \frac{1}{2E_N'} \delta^4(p' + k_\pi - q - p) \langle N'\pi | j_{\mu}^{nc}(0) | N \rangle \langle N'\pi | j_{\nu}^{nc}(0) | N \rangle^*\]

with \( M \) the nucleon mass and \( E_N' \) the energy of the outgoing nucleon. In the sum over initial and final nucleon spins,\(^3\) the bar over the sum denotes the average over the initial ones. As for the one particle states they are normalized so that \( \langle \bar{p}|\bar{p}' \rangle = (2\pi)^3 2p_0 \delta^3 (\bar{p} - \bar{p}') \), and finally for the neutral-current we take

\[
j_{\nu}^{nc} = \bar{\Psi}_\nu \gamma^\mu \left(1 - \frac{8}{3} \sin^2 \theta_W - \gamma_5\right) \Psi_u - \bar{\Psi}_d \gamma^\mu \left(1 - \frac{4}{3} \sin^2 \theta_W - \gamma_5\right) \Psi_d - \bar{\Psi}_s \gamma^\mu \left(1 - \frac{4}{3} \sin^2 \theta_W - \gamma_5\right) \Psi_s
\]

\( ^2 \) To obtain Eq. (2) we have neglected the four-momentum carried out by the intermediate Z-boson with respect to its mass \( (M_Z) \), and have used the existing relation between the gauge weak coupling constant, \( g = e/\sin \theta_W \), and the Fermi constant: \( G/\sqrt{2} = g^2/8M_W^2 \), with \( e \) the electron charge, \( \theta_W \) the Weinberg angle, \( \cos \theta_W = M_W/M_Z \) and \( M_W \) the W-boson mass.

\( ^3 \) Since right-handed neutrinos are sterile, only left-handed neutrinos are considered.

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with $\Psi_u$, $\Psi_d$ and $\Psi_s$ quark fields, and $\theta_W$ the Weinberg angle ($\sin^2\theta_W = 0.231$). Note that with all these definitions, the matrix element $\langle N'\pi|J_{\mu\nu}(0)|N\rangle$ is dimensionless. Both, lepton and hadron tensors are independent of the neutrino lepton family, and therefore the cross section for the reaction of Eq. (1) is the same for electron, muon or tau incident neutrinos. As the quantity $L^{(v)}_{\mu\sigma}(W^{NC}_\sigma) (v)$ is a Lorentz scalar, to evaluate it we take for convenience $\vec{q}$ in the $Z$-direction (see Fig. 1). Referring now the pion variables to the outgoing $\pi N$ pair center of mass frame (as it is usual in pion electroproduction) would be readily done by means of a boost in the $Z$ direction. Note that the azimuthal angle $\phi_\pi$ is left unchanged by such a boost.

For antineutrino induced reactions we have

$$L^{(\bar{v})}_{\mu\sigma} = L^{(v)}_{\sigma\mu}, \quad (W^{NC}_\sigma) (\bar{v}) = (W^{NC}_\sigma) (v).$$

For the sake of simplicity, from now on, we will omit in the hadronic tensor the explicit reference to $\nu$ or $\bar{\nu}$ and the NC$\pi$ label. By construction, the hadronic tensor accomplishes

$$W^{\mu\nu} = W^{\mu\nu}_s + i W^{\mu\nu}_a,$$

where $W^{\mu\nu}_s$ ($W^{\mu\nu}_a$) real symmetric (antisymmetric) tensors. The hadronic tensor is completely determined by up to a total of 19 Lorentz scalar and real, structure functions $W_i(q^2, p \cdot q, p \cdot k_\pi, k_\pi \cdot q)$

$$W^{\mu\nu}_s = (W^{\mu\nu}_s, s, a)^{PC} + (W^{\mu\nu}_s, s, a)^{PV},$$

$$W^{\mu\nu}_a = (W^{\mu\nu}_a, s, a)^{PC} + (W^{\mu\nu}_a, s, a)^{PV},$$

$$W^{(\nu)}_{\mu\sigma} = W^{\mu\sigma}_s \Leftrightarrow (W^{\mu\sigma}_s - W^{\mu\sigma}_a) \Leftrightarrow (W^{\mu\sigma}_a).$$

With $W^{\mu\sigma}_s$ ($W^{\mu\sigma}_a$) real symmetric (antisymmetric) tensors. The hadronic tensor is completely determined by up to a total of 19 Lorentz scalar and real, structure functions $W_i(q^2, p \cdot q, p \cdot k_\pi, k_\pi \cdot q)$, respectively. The triple differential cross section $\sigma_{NC}$ is not orthogonal to $q^\mu$, the W3, W5, W6, W8, W11, W12 and W17 terms do not contribute to the differential cross section since the leptonic tensor is orthogonal to the four vector $q^\mu$ for massless neutrinos.

The tensor $(W^{\mu\nu})^{PV} = (W^{\mu\nu})^{PC} + i(W^{\mu\nu})^{PV}$ when contracted with the leptonic one, $L^{(v)}_{\mu\nu}$, provides a pseudo-scalar quantity, i.e., such contraction is not invariant under a particle transformation. Indeed, under a parity transformation we have

$$L^{(\nu)}_{\mu\nu} \rightarrow (L^{(\nu)}_{\mu\nu})^{PV}, \quad (W^{\mu\nu})^{PV} \rightarrow -(W^{\mu\nu})^{PV}$$

whereas the tensor $(W^{\mu\nu})^{PC} = (W^{\mu\nu}_s)^{PC} + i(W^{\mu\nu}_a)^{PC}$ transforms as $(L^{\mu\nu})^{(v)}$. This explains the origin of the adopted labels PC and PV, which stand for parity violating and conserving contributions to the fifth differential cross section $d^5\sigma/d\Omega(k')dE'd\Omega(k')$, respectively.

The triple differential cross section $d^5\sigma/d\Omega(k')dE'd\Omega(k')$ is a scalar, up to the factor $|\vec{k}'|^2$. Thus all parity-violating contributions must disappear when performing the pion solid angle integration. Note that the coordinate system used to define $d\Omega(k')$ involves the pseudo-vector $\vec{k} \times \vec{k}'$ to set up the $Y$-axis, which induces the non-parity-invariant nature of $d^5\sigma/d\Omega(k')dE'd\Omega(k')$. In electroproduction processes, the leptonic tensor is purely symmetric, and the symmetric part of the hadronic one can not contain terms involving the Levi-Civita tensor, since the electromagnetic interaction preserves parity. Hence in case $d^5\sigma/d\Omega(k')dE'd\Omega(k')$ turns out to be a scalar under parity.

A final remark concerns the time-reversal ($T$) violation effects apparently encoded in the decomposition of the hadronic tensor in Eqs. (8)–(12). Under a time reversal transformation, and taking into account the antineutrino character of the $T$-operator, we have

$$L^{(\nu)}_{\mu\nu} \rightarrow (L^{(\nu)}_{\mu\nu})^{PV}, \quad (W^{\mu\nu})^{PV} \rightarrow -(W^{\mu\nu})^{PV}$$

and therefore $L^{(\nu)}_{\mu\nu} W^{PV}$ is not $T$-invariant either, because of the presence of the PV terms in the hadronic tensor. This does not necessarily means that there exists a violation of $T$-invariance in the process [5]. Time reversal is the transformation that changes the direction of time: $t \rightarrow -t$. Thus invariance under time reversal means that the description of physical processes does not show
any asymmetry if we look backward in time. The invariance under time reversal is equivalent to

\[ |M_{i \rightarrow f}|^2 = |M_{T_f \rightarrow T_i}|^2 \]  

where \( i \rightarrow f \) denotes the transition from the initial state \( i \) to the final state \( f \), being \( M_{i \rightarrow f} \) its corresponding transition amplitude, and \( T_f \) and \( T_i \) denote the states obtained from \( f \) and \( i \), respectively, by reversing the momenta, spins, etc. When one is dealing with electromagnetic or weak interactions, which may be treated to first order in the interaction Hamiltonian \((H_I)\), the transition matrix operator can be approximated by \( H_I \), being it then Hermitian. In these circumstances \( |M_{T_f \rightarrow T_i}| \approx |M_{T_i \rightarrow T_f}| \) and the invariance under time reversal may be written in the following form

\[ |M_{i \rightarrow f}|^2 = |M_{T_f \rightarrow T_i}|^2 \approx |M_{T_i \rightarrow T_f}|^2. \]  

The above equation implies that the transition \( i \rightarrow f \) cannot show up correlations that change sign under time reversal \((T\text{-odd correlations})\). Thus at first order in the interaction Hamiltonian, the tensor \((W^{\text{PV}})\) should vanish, since it would lead to violations of time reversal invariance. However, here we have strong interacting final states (pion and nucleon) and in fact the transition matrix operator is not Hermitian, since it cannot be approximated by \( H_I \). Therefore Eq. (16) does not hold and time reversal invariance does not forbid a non-vanishing \((W^{\text{PV}})\) tensor. In summary, besides genuine time reversal violations \([6], \) the existence of strong final state interaction effects \([7] \) allows for the existence of \( T\text{-odd correlations in } (L^{(o)}_{\mu \nu} W^{\mu \nu}) \), induced by the \((W^{\text{PV}})\) term of the hadronic tensor.

After this discussion, we are in conditions of studying the pion azimuthal angle dependence of the differential cross section. With our election of kinematics \((\hat{k}, \hat{k}')\) in the \(XZ\text{-plane})\), we find \((L^{(o)}_{\mu \nu})_{0y} = (L^{(o)}_{\mu \nu})_{xy} = (L^{(o)}_{\mu \nu})_{yz} = (L^{(o)}_{\mu \nu})_{0x} = (L^{(o)}_{\mu \nu})_{0z} = (L^{(o)}_{\mu \nu})_{xz} = 0 \), and then

\[
\int_0^{+\infty} \frac{d\kappa_x}{E_\pi} k_x^2 \left( L^{(o)}_{\mu \nu} \right)_{\mu \nu} W_{\mu \nu} = \int_0^{+\infty} \frac{d\kappa_x}{E_\pi} k_x^2 \left\{ \left( L^{(o)}_{\mu \nu} \right)_{00} W_0^0 + 2 \left( L^{(o)}_{\mu \nu} \right)_{0x} W_0^x + 2 \left( L^{(o)}_{\mu \nu} \right)_{0z} W_0^z + \left( L^{(o)}_{\mu \nu} \right)_{xx} W_x^x + \left( L^{(o)}_{\mu \nu} \right)_{yy} W_y^y + \left( L^{(o)}_{\mu \nu} \right)_{zz} W_z^z \right\} \\
= A_s + B_s \cos \phi_\pi + C_s \cos 2\phi_\pi + D_s \sin \phi_\pi + E_s \sin 2\phi_\pi,
\]

\[
\int_0^{+\infty} \frac{d\kappa_x}{E_\pi} k_x^2 \left( L^{(o)}_{\mu \nu} \right)_{\mu \nu} W_{\mu \nu} = 2 \int_0^{+\infty} \frac{d\kappa_x}{E_\pi} k_x^2 \left\{ \left( L^{(o)}_{\mu \nu} \right)_{0y} W_0^y + \left( L^{(o)}_{\mu \nu} \right)_{xy} W_x^y + \left( L^{(o)}_{\mu \nu} \right)_{yz} W_z^x \right\} = -A_a - B_a \cos \phi_\pi - D_a \sin \phi_\pi
\]

where we explicitly show the \( \phi_\pi \) dependence. \(^8\) The PV term of the hadronic tensor has led to the parity violating \( \sin \phi_\pi \) and \( \sin 2\phi_\pi \) contributions (all of them proportional to \( k_\pi \)). They disappear when the pion solid angle integration is performed, as anticipated.

Thanks to Eq. (6), we obtain for neutrino and antineutrino reactions

\[
(v, v) \rightarrow d^3\sigma_v \rightarrow d\Omega(k') dE' d\Omega(k_\pi) = \frac{E'}{k' \left| k' \right|} \frac{G^2}{16\pi^2} \left\{ A_s + A_a + (B_s + B_a) \cos \phi_\pi + C_s \cos 2\phi_\pi + (D_s + D_a) \sin \phi_\pi + E_s \sin 2\phi_\pi \right\},
\]

\[
(\bar{v}, \bar{v}) \rightarrow d^3\sigma_v \rightarrow d\Omega(k') dE' d\Omega(k_\pi) = \frac{E'}{k' \left| k' \right|} \frac{G^2}{16\pi^2} \left\{ A_s - A_a + (B_s - B_a) \cos \phi_\pi + C_s \cos 2\phi_\pi + (D_s - D_a) \sin \phi_\pi + E_s \sin 2\phi_\pi \right\}.
\]

3. Neutrino–antineutrino asymmetries

From Eqs. (19) and (20), we see that neutrino–antineutrino asymmetries may occur both in the totally integrated cross sections and in the pion azimuthal differential distributions. For the neutrino and antineutrino induced total cross sections we have

\[
(v, v) \rightarrow \sigma_v = \frac{G^2}{16\pi^2 \left| k' \right|} \int d\Omega(k') dE' d\Omega(k_\pi) \left| k' \right| (A_s + A_a),
\]

\[\text{Footnotes:}\]

\(^6\) One is faced with the problem of eliminating final-state-interaction effects, or at least of having them under control, in order to have a significant test of \( T \) violation \([5,6]\).

\(^7\) For instance, at intermediate energies and for CC driven processes, the \( \Delta(1232) \) resonance plays a central role \([7]\). The inclusion of the resonance width accounts partially for the strong final state interaction effects.

\(^8\) All structure functions, \( W_{i=1,...,19} \), depend on the Lorentz scalar \( p \cdot k_\pi \) and \( k_\pi \cdot q \) factors, which are functions of the angle formed between the \( \hat{q} \) and \( \hat{k}_\pi \) vectors, and thus they are independent of \( \phi_\pi \), when \( \hat{q} \) is taken along the \( Z\)-axis.
\( (\bar{\nu}, \bar{\nu}) \to \sigma_{\bar{\nu}} = \frac{G^2}{16\pi^2 |k|} \int d\Omega(\hat{k}) dE' \frac{d\Omega(\hat{k})}{dE} |\hat{k}|(A_s - A_d), \) \hspace{1cm} (22)

while for the azimuthal distributions, relations of the type

\[
\frac{d\hat{\sigma}}{d\Omega(\hat{k}) dE' d\Omega(\hat{k}_{\pi})} (\phi_{\pi}) - \frac{d\hat{\sigma}}{d\Omega(\hat{k}) dE' d\Omega(\hat{k}_{\pi})} (\phi_{\pi} + \pi) \bigg|_\nu = \left\{ (B_s + B_a) \cos \phi_{\pi} + (D_s + D_a) \sin \phi_{\pi} \right\} |\hat{k}| \frac{G^2}{8\pi^2}, \hspace{1cm} (23)
\]

\[
\frac{d\hat{\sigma}}{d\Omega(\hat{k}) dE' d\Omega(\hat{k}_{\pi})} (\phi_{\pi}) - \frac{d\hat{\sigma}}{d\Omega(\hat{k}) dE' d\Omega(\hat{k}_{\pi})} (\phi_{\pi} + \pi) \bigg|_{\bar{\nu}} = \left\{ (B_s - B_a) \cos \phi_{\pi} + (D_s - D_a) \sin \phi_{\pi} \right\} |\hat{k}| \frac{G^2}{8\pi^2}, \hspace{1cm} (24)
\]

enhance neutrino–antineutrino asymmetries. To increase the statistics in these azimuthal asymmetries, it is interesting to study the relations of Eqs. (23) and (24) for integrated cross sections. From the discussion in Appendix A it follows

\[
\frac{1}{2} \left( \frac{d\sigma(\phi_{\pi})}{d\phi_{\pi}} - \frac{d\sigma(\phi_{\pi} + \pi)}{d\phi_{\pi}} \right) \bigg|_\nu = (B_s + B_a) \cos \phi_{\pi} + (D_s + D_a) \sin \phi_{\pi}, \hspace{1cm} (25)
\]

\[
\frac{1}{2} \left( \frac{d\sigma(\phi_{\pi})}{d\phi_{\pi}} - \frac{d\sigma(\phi_{\pi} + \pi)}{d\phi_{\pi}} \right) \bigg|_{\bar{\nu}} = (B_s - B_a) \cos \phi_{\pi} + (D_s - D_a) \sin \phi_{\pi}, \hspace{1cm} (26)
\]

where for each event, \( \phi_{\pi} \) is the angle formed by the neutrino and pion scattering planes. Since the outgoing neutrino will not likely be detected, if the incoming neutrino momentum is known, the neutrino scattering plane will be set up by detecting also the outgoing nucleon. Indeed, the neutrino scattering plane is also determined by the incoming neutrino and the transferred momenta, and in the LAB system this latter one is given by the sum of the outgoing nucleon and pion momenta \( \vec{q} = \vec{p'} + \vec{k}_{\pi} \). Thus the strategy will be detecting in coincidences the outgoing pion and nucleon momenta, use those to determine the neutrino scattering plane and then measure the corresponding \( \phi_{\pi} \) angle.

All these relations can be used to distinguish \( \tau \)-neutrinos from antineutrinos, below the \( \tau \)-production threshold, but above the pion production one. Thus and using a proton target, the total cross section or the dependence on the angle formed by the neutrino and pion planes in either of the channels,

\[
\nu p \to \nu n \pi^+, \quad \bar{\nu} p \to \bar{\nu} n \pi^+, \quad \nu p \to \nu p \pi^0, \quad \bar{\nu} p \to \bar{\nu} p \pi^0, \hspace{1cm} (27)
\]

may be used to determine the nature of the incident \( \tau \)-neutrino beam below the \( \tau \)-production threshold.

Moreover, we would like to point out that the \( (A_s \pm A_d), (B_s \pm B_a) \) and \( (D_s \pm D_a) \) response functions, the integrated cross sections \( \sigma_{\nu}, \sigma_{\bar{\nu}} \) and the partially integrated quantities \( (B_s \pm B_a) \) and \( (D_s \pm D_a) \) are independent of the lepton family and can be experimentally measured by using electron or muon neutrinos, due to the lepton family universality of the NC neutrino interaction. Thus one would not have to rely on any theoretical model to determine the asymmetry relations proposed in this section.

3.1. Neutrino–antineutrino asymmetries from the chiral model of Ref. [4]

The above suggestion for distinguishing neutrino from antineutrino in NC pion production reactions will be of no value if the terms \( A_s, B_s \) and \( D_s \) are much smaller than \( A_d, B_a \) and \( D_a \). To estimate these response functions and address this issue, we have used the model developed in Ref. [4] to study the weak pion production off the nucleon at intermediate energies. In this model, besides the Delta pole mechanism \( \Delta P \) (weak excitation of the \( \Delta(1232) \) resonance and its subsequent decay into \( NN \)), some background terms, required by chiral symmetry, are also considered. These chiral background terms are calculated within an SU(2) non-linear \( \sigma \) model involving pions and nucleons, which implements the pattern of spontaneous chiral symmetry breaking of QCD. The model gives a fair description of the neutrino and antineutrino CC and NC available data. Details can be found in Ref. [4].

In Fig. 2, we show neutrino and antineutrino NC pion production total cross sections predicted by the model of Ref. [4], as a function of the incoming neutrino or antineutrino energy. There, 68% CL bands (external lines), inferred from uncertainties of the model of Ref. [4] are also displayed. Those uncertainties come from re-adjusting the \( C_4^A(q^2) \) form-factor that controls the largest term of the \( \Delta \)-axial contribution. On the other hand, since the main dynamical ingredients of the model are the excitation of the \( \Delta \) resonance and the non-resonant contributions deduced from the leading SU(2) non-linear \( \sigma \) Lagrangian involving pions and nucleons, we will concentrate in the pion–nucleon invariant mass \( W \leq 1.4 \) GeV region. For larger invariant masses, the chiral expansion will not work, or at least the lowest order used in Ref. [4] will not be sufficient and the effect of heavier resonances will become much more important [8]. For this reason in the figure, we always plot cross sections calculated with an invariant pion–nucleon mass cut of 1.4 GeV, it is to say we compute \( \int_{m_{\pi} + M}^{1.4 \text{ GeV}} dW d\sigma/dW \). Such cut is commonly used in the CC pion production experimental analyses at intermediate energies. See for instance [9] where it can be also been seen that up to incoming neutrino LAB energies of the order of 1 GeV the implementation of this cut hardly changes the measured cross section. In the figure, we also display low energy data measured in the Argonne National Laboratory (ANL) [10] for the \( \nu n \to \nu p \pi^- \) reaction, which do not
Fig. 2. Neutrino and antineutrino NC pion production total cross sections predicted by the model of Ref. [4], as a function of the incoming neutrino or antineutrino energy. An invariant pion–nucleon mass cut of $W \leq 1.4$ GeV has been implemented. Central lines stand for the result of the model. 68% CL bands (external lines), inferred from uncertainties of the model, are also displayed. Experimental cross sections for neutrino $\pi^-$ production are taken from Ref. [10].

include any cut in $W$. As shown in Ref. [4], including or not in the theoretical calculation the cut in $W$ only influences significantly the result for the highest energy data-point, leading to variations of the order of 20%, which are around half the experimental error.

Results of Fig. 2 are really significant since the differences between neutrino and antineutrino induced reactions are always large, about a factor of two, in all physical channels. This confirms that this observable can be used to distinguish neutrino from antineutrino above the pion production threshold.

Next we show results, within this model, for $B_s^\pm B_a$ and $D_s^\pm D_a$ defined in Eqs. (25) and (26). In Fig. 3 we plot these response functions for a $p\pi^-$ final state. The rest of channels lead to similar results. Clear neutrino–antineutrino asymmetries can be appreciated in both panels of the figure. In particular $D_s + D_a$ and $D_s - D_a$ have opposite signs. The difference between the number of events for which the pion comes out above the neutrino plane ($N_{abv}$) and those where the pion comes out below this plane ($N_{blw}$) is 4 ($D_s + D_a$) [4 ($D_s - D_a$)] for a neutrino [antineutrino] induced process. Thus, the sign of this difference would determine whether one has a neutrino or an antineutrino beam. Nevertheless, we should mention that the neutrino–antineutrino asymmetry based in the total cross section discussed above (Fig. 2) would provide a signal with much better statistical significance. This is because the ratio $N_{abv} - N_{blw} / N_{abv} + N_{blw}$ would be of the order of $10^{-2}$ or smaller. Asymmetries based on the $B$ response function will not be as much unfavorable, from the statistical point of view.

These results point out that the neutrino–antineutrino asymmetries based on the total cross sections are certainly much more useful than those based on the pion azimuthal response functions, for intermediate neutrino energies. The situation could be different for energies larger than those explored with the model of Ref. [4].

Acknowledgements

We thank M.J. Vicente-Vacas for useful discussions. J.N. acknowledges the hospitality of the School of Physics & Astronomy at the University of Southampton. This work was supported by DGI and FEDER funds, under contracts FIS2005-00810, BFM2003-00856 and FPA2004-05616, by Junta de Andalucía and Junta de Castilla y León under contracts FQM0225 and SA104/04, and it is a part of the EU Integrated Infrastructure Initiative Hadron Physics Project contract RII3-CT-2004-506078.
Appendix A. Partially integrated azimuthal distribution

To define an azimuthal pion distribution, by integrating the outgoing neutrino variables and pion polar angle in Eqs. (19) and (20), requires for each event a rotation to guarantee that \( \hat{q} \) is in the \( Z \)-direction and that \( \vec{k} \) and \( \vec{k}' \) are contained in the \( XYZ \)-plane. The totally integrated unpolarized neutrino cross section reads

\[
\sigma_v = \frac{G^2}{16\pi^2|k|} \int \frac{d^3k'}{|k'|} \frac{d^3k_\pi}{E_\pi} L_{\mu\sigma}^{(v)}(k,k') W^{\mu\sigma}(p,q,k_\pi).
\]  

(A.1)

To perform the above integrals, we take a fixed system of coordinates \( XYZ \), and since we have \( k_\pi \) momenta fix the latter one. Thus, \( \phi_\pi \) is the angle formed by the neutrino and pion planes. The first of these planes is determined by the incoming (\( k \)) and pion planes. The antineutrino cross section can be obtained in a similar way, and thus it follows the neutrino–antineutrino asymmetry relations given in Eqs. (25) and (26) similar to those in Eqs. (23) and (24), but for integrated cross sections.
References