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PION AND STRANGE PARTICLE PRODUCTION BY 1 GEV NEUTRINOS

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A B S T R A C T

Estimates are made of cross-sections for the production of pions or strange particles by high-energy neutrinos incident on nuclei. At 1 GeV the production of pions via the (33) resonance is found to augment the total cross-section by about 50 %.

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1. Introduction

There have been several discussions ¹⁾ of "elastic" neutrino induced processes of the type

$$\nu + n \rightarrow p + e \text{ (or } \mu \text{)}. \quad (1.1)$$

Here we present some crude estimates for inelastic processes in which pions or strange particles are produced. The replacement of the final nucleon by a hyperon is considered in Section 2. Single pion production is discussed in Section 3. In Section 4 some observations are made on the case of a nucleus rather than a nucleon as target.

In general we ignore the lepton mass, so strictly speaking we deal with electrons rather than muons. Calculations are made particularly for a neutrino energy of about 1 GeV, in the range relevant for current experiments ¹⁰⁾. The elastic process (1) is then expected ¹⁾ to have a cross-section of about $0.85 \times 10^{-38} \text{ cm}^2$. We use units such that $\hbar = c = M$ (nucleon mass) = 1, and the metric

$$q^2 = q_1^2 + q_2^2 + q_3^2 + q_4^2 = q_1^2 + q_2^2 + q_3^2 - q_0^2.$$

2. Strange particle production

Perhaps the simplest inelastic processes involving neutrinos or anti-neutrinos are those where the target nucleon is converted to a hyperon, i.e.,

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$$\nu + n \rightarrow \Sigma^+ + e^-(\mu^-) \quad (2.1a)$$

$$\bar{\nu} + p \rightarrow \Lambda(\Sigma^0) + e^+(\mu^+) \quad (2.1b)$$

$$\bar{\nu} + n \rightarrow \Sigma^- + e^+(\mu^+) \quad (2.1c)$$

If the $\Delta S = \Delta Q$ rule is valid then (2.1a) is forbidden.

The baryon current for a particular reaction (2) will have the general structure

$$J_\mu = G \left[(F_1 + iF_2) \gamma_\mu + iF_2 K_\mu + F_3 q_\mu - G_1 \gamma_\mu \gamma_5 - iG_2 K_\mu \gamma_5 - iG_3 q_\mu \gamma_5 \right] \quad (2.2)$$

where G is the weak coupling constant, K_μ the nucleon 4-momentum, q_μ the hyperon momentum, and the form factors F and G are functions of q^2 .

The contributions from F_3 and G_3 contain the lepton mass as a factor, and will be ignored. With the interaction

$$J_\mu \bar{\Psi}_e \gamma_\mu (1 + \gamma_5) / \sqrt{2} \Psi_\nu \quad (2.3)$$

$$J_\mu \bar{\Psi}_e \gamma_\mu (1 - \gamma_5) / \sqrt{2} \Psi_\nu \quad (2.4)$$

the differential cross-section in the laboratory frame is ¹¹⁾:

$$\begin{aligned}
\frac{d\sigma}{d\Omega} = & \frac{G^2}{2\pi^2} \frac{e'^2}{1+2es^2} \left\{ \left[|F_1|^2 + |\mu F_2|^2 \gamma^2 + \text{Re} \mu F_1 F_2^* (1-\gamma) \right] [1-s^2] \right. \\
& + 2s^2 \gamma |F_1 + \mu F_2|^2 \\
& + \left[|G_2|^2 \gamma^2 - \text{Re} G_1 G_2^* \gamma \right] [1-s^2] \\
& + \frac{1}{4} q^2 \left[2 |F_1 + \mu F_2|^2 s^2 + (|\mu F_2|^2 + |G_2|^2)(1-s^2) \right] \\
& + G_1^2 \left[1-s^2 + (1-\gamma)^2 s^2 + \frac{1}{2} q^2 s^2 \right] \\
& \left. \pm 2 \text{Re} \left[G_1^* (F_1 + \mu F_2) s^2 \right] \left[2e - \frac{1}{2} q^2 + 2\gamma(1-\gamma) \right] \right\} \quad (2.5)
\end{aligned}$$

where

$$\begin{aligned}
s &= \sin \frac{1}{2} \theta \\
\theta &= \text{lab. scattering angle} \\
e &= \text{initial neutrino energy} \\
e' &= \text{final electron energy} \\
&= [e - 2\gamma(1-\gamma)] / [1 - 2es^2] \\
2\gamma &= \text{hyperon-nucleon mass difference in nucleon mass units} \\
q^2 &= (4\text{-momentum transfer})^2 \\
&= 4e^2 s^2 [1 - 2\gamma(1-\gamma)/e] [1 + 2es^2]^{-1}.
\end{aligned}$$

Little is known about the form factors F and G . In the universal Feynman - Gell-Mann theory, ignoring renormalization, $F_1 = G_1 = 1$ and the others are zero. Leptonic decay rates of the hyperons predicted on this basis are much too large compared with experiment²⁾. The disagreement is by a factor of about ten for Λ , about fifty for Σ^- , and by a still larger factor for

4.

Σ^+ (where no leptonic decay has been seen, in agreement with $\Delta S = \Delta Q$). It has been shown by Berman ⁷⁾, and by Fujimara, Katani, and Saito ⁹⁾, that the slowness of decay cannot to any great extent be attributed to accidental cancellation among the various contributions. The individual form factors (or the coupling constant) must be much smaller than for the nucleon. Unless the q^2 dependence of form factors is greatly different from in the nucleon case, we can then expect that the reactions (2.1) are very much less probable than the elastic process (1.1).

3. Pion production

a) Via the (33) resonance

We consider the process

$$\nu + n \rightarrow (p \text{ or } n) + e + \pi^-.$$

At the energies of interest a substantial contribution to this must be made by the formation and decay of the (33) resonance. This can be treated very simply in the static theory ⁴⁾, following the Chew - Low account of photoproduction ³⁾.

First an appropriate form of the weak-interaction current J_μ is required. Between real nucleon states it has the matrix element ¹⁾

$$iG \bar{u} \gamma^+ \left\{ F_1(q^2) \gamma_\mu - (1/2M) F_2(q^2) \sigma_{\mu\nu} q_\nu + \lambda F_A(q^2) \gamma_\mu \gamma_5 \right\} u. \quad (3.1)$$

The induced pseudoscalar does not contribute appreciably for electrons and is omitted. The corresponding non-relativistic form involving 2-spinors is

$$G u \gamma^+ \left\{ \begin{array}{l} \left[\frac{1}{2}(\underline{K} + \underline{K}') - \frac{i}{2}(1+\mu) \underline{\sigma} \times \underline{q} - \lambda \underline{\sigma} \right] \\ i \left[1 - \frac{1}{2} \lambda \underline{\sigma} \cdot (\underline{K} + \underline{K}') \right] \end{array} \right\} u. \quad (3.2)$$

This is correct as far as terms linear in the initial neutron momentum $\underline{K}$ and final proton momentum $\underline{K}'$. But the terms involving $(\underline{K} + \underline{K}')$ cannot be retained in a no-recoil theory. However, $(\underline{K} - \underline{K}') = \underline{q}$ can be referred to the leptons, by momentum conservation, so that the $\underline{q}$ dependence can be retained. The term $\underline{\sigma} \times \underline{q}$ must certainly be retained, because $(1+\mu)$ is quite large, but for the present we drop the q^2 dependence of form factors. Then the remaining terms in (3.2) can be obtained in the static theory as the matrix element between real nucleon states of

$$J_{\mu}^{\prime}(q) = \begin{pmatrix} J_{\mu}^{\prime} \\ J_4 \end{pmatrix} = Z^{-1} \bar{N} \tau^+ \left\{ \begin{array}{l} \left[-\lambda \underline{\sigma} - \frac{i}{2}(1+\mu) \underline{\sigma} \times \underline{q} \right] \\ i \end{array} \right\} N. \quad (3.3)$$

Here $\bar{N}$ and N create and absorb bare nucleons, and $Z^{+1/2}$ is the projection of the real on the bare nucleon. The Fourier transform of the complete current operator in the static theory can therefore be written

$$J_{\mu}(q) = J_{\mu}^{\prime}(q) + J_{\mu}^{\prime\prime}(q) \quad (3.4)$$

where $J^{\prime\prime}$ has zero matrix elements between real nucleon states. In "elastic" scattering $J^{\prime}$ alone is operative. Moreover, in photopionproduction Chew and Low found that the analogue of $J^{\prime\prime}$ was little enhanced by the (33) resonance and could be treated in Born approximation. This should be again true here, and in a first orientation we omit $J^{\prime\prime}$. We omit also the fourth component of

6.

J' in (3.3), because it cannot excite the (33) resonance. There remains

$$J_\mu = \begin{pmatrix} J \\ \bar{w} \\ 0 \end{pmatrix} \quad J = Z^{-1} \bar{N} \gamma^+ (-\lambda \bar{\sigma} - \frac{i}{2} \xi \bar{\sigma} \times q) N \quad (3.5)$$

with $(1+\mu) = \xi$.

Let $k = (\underline{k}, e)$ and $k' = (\underline{k}', e')$ be initial and final lepton 4-momenta. Let $|I\rangle$ be the initial neutron and $|F\rangle$ a final nucleon plus pion state. With the interaction

$$- \frac{1}{\sqrt{2}} J_\mu \bar{\Psi}_e^i \gamma_\mu (1 + \gamma_5) \Psi_\nu$$

the cross-section for producing electrons in dk' is

$$d\sigma = \left(dk' / (2\pi)^3 \right) 2\pi |\mathcal{M}|^2 \quad (3.6)$$

where

$$\mathcal{M} = M_{\mu\nu} M_{\mu\nu}^* \quad (3.7)$$

The lepton factor is

$$\begin{aligned} M_{\mu\nu} &= \frac{1}{2} \text{Trace} \quad i \gamma_\mu (1 + \gamma_5) \frac{i \gamma_{k'}}{2e'} \quad i \gamma_\nu (1 + \gamma_5) \frac{i \gamma_k}{2e} \\ &= \frac{1}{e'e} \left\{ k'_\mu k_\nu + k_\mu k'_\nu - \delta_{\mu\nu} (k' \cdot k) + \epsilon_{\mu\nu\rho\sigma} k'_\rho k_\sigma \right\}. \end{aligned} \quad (3.8)$$

The heavy particle factor is

$$M_{\mu\nu} = \sum_I \sum_F (I | J_{\mu}^+(q) | F) (F | J_{\nu}(q) | I) \delta(E_F - E_I - \omega) \quad (3.9)$$

where $\omega = e - e'$. Summation over all final states is indicated, and averaging over initial spin.

With (3.5) only the space part of $M_{\mu\nu}$ survives

$$M_{ij} = M(\omega) \left\{ \left[\lambda^2 + \frac{1}{4} \right] \xi_{\underline{q} \cdot \underline{q}}^2 \right\} \delta_{ij}^{-\frac{1}{4}} \xi_{\underline{q}_i \underline{q}_j + i}^2 \xi_{\lambda} \epsilon_{ijk} q_k \quad (3.10)$$

The function $M(\omega)$ of energy transfer ω is defined by

$$M(\omega) \delta_{ij} = \sum_I \sum_F (I | \sigma_i \tau^- | F) (F | \sigma_j \tau^+ | I) \delta(E_F - E_I - \omega). \quad (3.11)$$

Now, this same quantity occurs in the total cross-section for the scattering of positive pions by neutrons. The pion interacts with a nucleon current

$$\sqrt{2} \frac{f_0}{m} \frac{1}{\sqrt{2\omega}} \bar{N} \tau^+ \underline{\sigma} \cdot \underline{p}_N$$

where $\underline{p}$ is the pion momentum, $\omega = (\underline{p}^2 + m^2)^{\frac{1}{2}}$ m is the pion mass and f_0 the unrenormalized coupling constant. The total cross-section for pions of energy ω is

$$\begin{aligned} \sigma_{\pi n}(\omega) &= \frac{\omega}{p} 2\pi \sum_I \sum_F (I | \frac{f_0}{m} \frac{1}{\sqrt{\omega}} \bar{N} \underline{\sigma} \cdot \underline{p} \tau^- | F) (F | \frac{f_0}{m} \frac{1}{\sqrt{\omega}} \bar{N} \underline{\sigma} \cdot \underline{p} \tau^- | I) \\ &= 2\pi (\omega^2 - m^2)^{\frac{1}{2}} \left(\frac{f_0}{m} \right)^2 M(\omega). \end{aligned} \quad (3.12)$$

Combining (3.6) - (3.12) we have the following cross-section for leptons emerging in solid angle $d\Omega$, integrated over direction and energy ω of the pion

$$d\sigma = d\Omega \int d\omega \frac{\sigma_{\pi n}(\omega)}{\sqrt{\omega^2 - m^2}} \left(\frac{m}{f}\right)^2 \frac{g^2 e'^2}{4\pi^3} \left\{ \lambda^2(1+s^2) + \xi^2 \left(\frac{1}{2}(e^2 + e'^2)s^2 + e'e s^4 \right) + 2\xi\lambda(e' + e)s^2 \right\} \quad (3.13)$$

where $s = \sin \frac{1}{2} \theta$, θ is the angle of the electron relative to the neutrino; $f = f_0 Z$ is the renormalized coupling constant, experimentally nearly equal to unity. For comparison, the static theory yields for the elastic process

$$d\sigma = d\Omega \frac{g^2 e^2}{2\pi^2} \left\{ (1-s^2) + \lambda^2(1+s^2) + \xi^2 e^2(s^2 + s^4) + 4\xi\lambda e s^2 \right\}. \quad (3.14)$$

The first term arises from the fourth component of J_μ , which did not contribute to (3.13). Equ. (3.14) can be obtained from the complete relativistic formula¹⁾ by suppressing terms of relative order (q/M) but retaining those in $(\xi q/M)$.

We make a crude numerical comparison of (3.13) and (3.14) for neutrino energy ≈ 940 MeV. The pion nucleon cross-section $\sigma_{\pi n}(\omega)$ is peaked at the resonant energy $\omega_0 = 330$ MeV. So elsewhere in (3.13) we set $\omega = \omega_0$, $e' = e - \omega_0$. In the forward direction the ratio of inelastic to elastic scattering is

$$\frac{1}{2\pi} \frac{\lambda^2}{1+\lambda^2} \left(\frac{e - \omega_0}{e} \right)^2 \left(\frac{m}{f} \right)^2 \frac{1}{\sqrt{\omega_0^2 - m^2}} \int d\omega \sigma_{\pi n}(\omega). \quad (3.15)$$

The resonance has a width at half height $\Gamma = 140$ MeV = 0.15 proton masses. It has a height h , for positive pions on neutrons, of about $66 \text{ mb} = 150 (\text{proton Compton wavelength})^2$. Thus

$$\int \delta(\omega) d\omega = \frac{\pi}{2} \hbar \Gamma \approx 35.4. \quad (3.16)$$

Then one finds with $f = 1$, $\lambda = 1.25$, $e = 1$, $\omega_0 = 0.32$.

$$\frac{\text{inelastic}}{\text{elastic}}(\text{forward}) \approx 0.10. \quad (3.17)$$

For angles of more than a few degrees both (3.13) and (3.14) should be substantially reduced by recoil effects ⁴⁾, and more importantly by the q^2 dependence of form factors. With certain assumptions Yamaguchi finds that $d\sigma/d\Omega$ in the elastic case falls to half its forward value at about 30° . These effects should cancel to a large extent in the inelastic to elastic ratio. However, one can guess that the inelastic falls off somewhat less rapidly than the elastic for two reasons:

- 1) the argument of the form factors

$$q^2 = 4ee' \sin^2 \frac{1}{2}\theta$$

is smaller at a given angle.

- 2) the recoiling object, essentially the (33) isobar, is heavier.

Moreover, at angles of about 30° , the expressions (3.13) and (3.14) are already dominated by the weak magnetic (ξ) and interference ($\xi\lambda$) terms, rather than the Gamow - Teller (λ^2) terms. The ratio (inelastic/elastic) increases somewhat for this reason also. One expects therefore that the ratio of total cross-sections is larger than (3.17). For the elastic process Yamaguchi quotes a total cross-section $0.86 \times 10^{-38} \text{ cm}^2$, whence

$$\sigma(\text{inelastic}) \gtrsim 0.9 \times 10^{-39} \text{ cm}^2. \quad (3.18)$$

We derived (3.13) for neutrinos on neutrons. With a proton target the pion cross-section $\sigma_{\pi n}$ should be for positive pions incident on protons, rather than neutrons. For incident antineutrinos one takes for σ the cross-section for incident negative pions, and changes the sign of the $(\xi\lambda)$ term.

b) Comparison with electron-pion production

Processes induced by neutrinos or electrons are closely related if the CVC (conserved vector current) is assumed and the axial part of the weak interaction set aside. Experimental data for electrons could then be used to predict the cross-section for neutrinos. This would give a lower limit, if we believe the sign of the vector-axial interference term from the static model. Consider the diagrams shown in Figs. 1a and 1b. To obtain the (b) cross-section from the (a) the following considerations are involved:

- 1) the square of the photon propagator $(1/q^2)^2$ is removed.
- 2) the electromagnetic ($\alpha = 1/137$) coupling is replaced by the weak coupling: $(4\pi\alpha)^2 \rightarrow G^2$.
- 3) the lepton trace involves $\gamma_\mu(1+\gamma_5)/\sqrt{2}$ instead of γ_μ . This causes no change. But because the neutrinos are polarized there is an extra 2.
- 4) the isoscalar part of the electric current does not excite the (33) resonance, so there the important isospin matrix element is

$$(I = 3/2, I_3 = \frac{1}{2} \mid \frac{1}{2}\gamma_3 \mid I = \frac{1}{2}, I_3 = \frac{1}{2}).$$

For the neutrino we have

$$(I = 3/2, I_3 = \frac{1}{2} \mid \gamma_+ \mid I = \frac{1}{2}, I_3 = -\frac{1}{2}).$$

These are equal.

The result is that the differential cross-sections for final leptons, in solid angle $d\Omega$ and energy range de' , are related by ⁸⁾

$$\frac{d\sigma_\nu}{d\Omega de'} = 2 \left(\frac{Gq^2}{4\pi\alpha} \right)^2 \frac{d\sigma_e}{d\Omega de'}. \quad (3.19)$$

Unfortunately, electron data exist only for a few fixed large angles ⁵⁾. This has not much value for predicting the outcome of such neutrino experiments as are likely to be done soon. But we can make some comparison with the static model formula (3.13). For large angles (3.13) can be good only at low energy, because of form factor and recoil effects. For small e' , omitting the axial part, it becomes

$$\frac{d\sigma_\nu}{d\Omega de'} = \frac{\sigma_{nn}(e)}{\sqrt{(e^2 - m^2)}} \left(\frac{m}{f} \right)^2 \frac{G^2 e'^2 e^2 q_s^2}{8\pi^3}. \quad (3.20)$$

From (3.19) we then infer for electrons

$$\frac{d\sigma_e}{d\Omega de'} = \frac{\sigma_{nn}(e)}{\sqrt{(e^2 - m^2)}} \left(\frac{m}{f} \right)^2 \frac{\alpha^2}{s^2} \frac{e^2}{16\pi}. \quad (3.21)$$

In one experiment ⁵⁾ the angle of scattering was fixed at 75° , and the initial energy chosen to give maximum excitation of the resonance. The variable is then q^2 , or e'^2 , and for small values we can compare the result with (3.21). With $e \approx 330$ MeV, $\sigma_{nn} \approx 70$ mb, this formula gives

$$\frac{d\sigma_e}{d\Omega de'} \approx 3 \cdot 10^{-34} \text{ cm}^2/\text{ster/MeV}.$$

This agrees sufficiently with the value obtained by Panofsky and Allton, both experimentally and by inference from photoproduction:

$$4 \cdot 10^{-34} \text{ cm/ster/MeV.}$$

Of course, formulae of the type (3.21) are already well known in the electron context.

c) One-pion exchange

The simplest one-pion exchange diagram is shown in Fig. 2. It describes the "elastic" scattering of the lepton on a virtual pion, which is ejected. For the moment we consider the exchange of a neutral pion. The vector part of the lepton-pion interaction is known on the CVC hypothesis, and the axial part we again ignore. The weak interaction vertex then contributes to the amplitude a factor

$$G \bar{\Psi}_e \gamma_\mu (1 + \gamma_5) \Psi_\nu (\pi - q)_\mu F_\pi ((\pi + q)^2) \quad (3.22)$$

where F_π is the pion-electromagnetic form factor. In this calculation we use the simple Clementel - Villi form ⁶⁾

$$F(x) = x_r / (x + x_r) \quad (3.23)$$

$$x_r = 22m_\pi^2.$$

The differential cross-section is

$$\frac{d\sigma}{dk' dK' d\pi} = \left(\frac{1}{2\pi}\right)^5 \frac{4\pi g^2 G^2}{E' E \omega e' e} \frac{q^2 (2\pi \cdot k \pi \cdot k' + m^2 k \cdot k')}{(m_\pi^2 + q^2)^2} \quad (3.24)$$

$$\delta^4(K' + k' + \pi - K - k) F_{\pi N}^2 ((k' - k)^2) F_{\pi N}^2 (q^2).$$

where $F_{\pi N}$ is the form factor for the pion nucleon vertex, $F_{\pi N}(-m_\pi^2) = 1$, and $g^2 (\approx 15)$ is the strong coupling constant. Integration over the electron and pion yields

$$\frac{d\sigma}{dK'} = \left(\frac{1}{2\pi}\right)^3 \frac{g^2 G^2}{E' E e} \frac{q \cdot k (q^2 + 2q \cdot k + m_\pi^2)^2 q^2}{(q^2 + 2q \cdot k)(q^2 + m_\pi^2)^2} \quad (3.25)$$

$$\frac{x_r}{2q \cdot k + x_r} F_{N\pi}^2 (q^2).$$

The remaining integration, to obtain the total cross-section, was done numerically. The pion mass was here made zero. This is permissible because the pole has been removed (by the γ_5 interaction) and the total comes mainly from $q^2 \gg m_\pi^2$. It also means, unfortunately, that the result is sensitive to the assumed form for $F_{N\pi}$, which has been taken here as a sharp cut-off at the nucleon mass.

If the target is a proton only a neutral pion can be exchanged, emerging as a positive pion. With a neutron a negative pion can also be exchanged, emerging as a neutral. Allowing for this, by a factor 3, the following cross-sections are obtained

$$\begin{aligned} \sigma &= 0.06 \times 10^{-38} \text{ cm}^2 & \text{for neutrino energy } e &= M \\ &= 0.26 \times 10^{-38} \text{ cm}^2 & e &= 4M \\ &= 1.0 \times 10^{-38} \text{ cm}^2 & e &= 10M. \end{aligned}$$

14.

4. Nuclei as targets

When the target is a nucleus, rather than a neutron, a number of factors enhance the inelastic relative to the elastic process.

In the first approximation we simply add the cross-sections of the individual nucleons. Now while the elastic process goes only on neutrons

the inelastic process goes on both neutrons and protons

Moreover, in so far as the (33) resonance is dominant, the cross-section on protons is three times the inelastic cross-section on neutrons. So in going from a neutron to a nucleus the magnitude of the inelastic relative to the elastic process is increased by a factor of about four.

The exclusion principle works in the same direction. It requires that the recoiling nucleon have enough momentum to lie outside the Fermi sphere. This appreciably reduces the elastic cross-section²⁾. But it has less effect on the inelastic, where recoil momenta are typically larger. The nucleon recoil in the decay of the (33) isobar is about 0.32M, which is about equal to the Fermi momentum, $\approx 0.28 M$. The exclusion principle is relevant for intermediate as well as final states, but this effect should be small. In intermediate states (at least in the static model) nucleon momenta are typically of order M.

The untypical case of small momentum transfer to the nucleon should be considered separately. For then there is the possibility of different nucleons contributing coherently. This requires that it is a π^+ which is produced, so that the nucleon charge is unchanged. Rough estimates, forgetting the reabsorption and scattering of the pion, indicate that the extra coherent contribution might be comparable with the incoherent. In fact, the coherent contribution must be greatly conditioned by the final state interaction of the pion, which is strongest precisely in the energy range where most pions are produced.

Conclusion

For neutrinos incident on neutrons, we expect the inelastic cross-section at 1 GeV to be about 10 % of the elastic. This is based on the production of pions via the (33) isobar. This process goes more readily on protons. Mainly for that reason we expect the ratio of inelastic to elastic to increase to perhaps about 50 % when the target is a nucleus. This would reduce to about 4 % for an incident neutrino energy of 500 MeV.

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Fig. 1a

Electron on proton

Fig. 1b

Neutrino on neutron

Fig. 2

Simplest one-pion exchange diagram

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only because the static approximation is made in the pion-nucleon
rest system, rather than in the laboratory. The improved kinematics
reduces the forward inelastic cross-section by some 20 % at 1 GeV.
A new factor (M/W) , where W is essentially the isobar mass, causes
a further reduction. The retention of such a factor after a no-recoil
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