νN and $\bar{\nu} N$ interactions at high energies in a theory without suppression of charm-changing currents

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It is shown that a change in the slope of the energy dependence of the νN interaction cross section should be expected at energies above the threshold for the production of charmed particles in a theory without suppression of weak hadron currents in transitions with change of charm ($\Delta c = 1, \Delta s = 0$).

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In[1,2] we formulated hypotheses concerning the inclusion of weak hadronic currents in a regular representation of the $SU(4)$ group, $SU(2)$ current algebra for charges, and universality of weak interactions for leptons and hadrons. On the basis of these hypotheses we obtained two types of solutions for the form of weak hadron currents with change of strangeness and charm. The first of the solutions corresponds to the usual GIM scheme[3] formulated in the quark model. The second solution differs appreciably from the GIM scheme for $n \to p'$ transitions. The corresponding current is given by

$$J_\mu^{\pm} = \cos\theta_c (V_\mu^{1\pm i2} - A_\mu^{1\pm i2}) + \sin\theta_c (V_\mu^{4\pm i5} - A_\mu^{4\pm i5}) + (1 - \sin\theta_c) + V_\mu^{11\mp i12} + (1 + \sin\theta_c) A_\mu^{11\mp i12} + \cos\theta_c (V_\mu^{13\mp i14} - A_\mu^{13\mp i14}). \qquad (1)$$

It is seen from (1) that the current of the $n \to p'$ transition is not suppressed (comparable with the GIM scheme) and has a $V+A$ character. It is quite obvious that this property of the $n \to p'$ current leads to definite physical consequences.

We consider in this article certain conclusions that follow from the form of the current (1) as applied to the physics of high-energy neutrinos. It is well known that within the framework of Bjorken scaling the differential distributions of the secondary muons with respect to the dimensionless variables $x = Q^3/2M\nu$ and $y = \nu/E_\nu$ in the reactions

$$\nu_\mu(\tilde{\nu}_\mu) + N \to \mu^-(\mu^+) + \text{hadrons} \qquad (2)$$

take the form

$$\frac{d^2\sigma^{\nu,\bar{\nu}}}{dx\,dy} = \frac{G^2 M E_\nu}{\pi}\left[\left(1 - y + \frac{y^2}{2}\right)F_2(x) + \frac{y^2}{2}(2xF_1 - F_2) \pm y\left(1 - \frac{y}{2}\right)xF_3(x)\right] \qquad (3)$$

where F_i are structure functions. Within the framework of the quark parton model the structure functions F_1 and F_2 satisfy the relation[4] $2xF_1 = F_2$, and the functions F_2 and F_3 are expressed in this case in terms of the distribution functions of the various types of partons $p(x)$, $n(x)$, $\lambda(x)$, and $p'(x)$ and antipartons $\bar{p}(x)$, $\bar{n}(x)$, $\bar{\lambda}(x)$, and $\bar{p}'(x)$ in the nucleon.

In the scheme considered by us in[1,2], the cross sections $\sigma(\nu N) = [\sigma(\nu p) + \sigma(\nu n)]/2$ and $\sigma(\bar{\nu} N) = [\sigma(\bar{\nu} p) + \sigma(\bar{\nu} n)]/2$ are expressed in terms of the distribution functions of the partons in the proton in the following manner:

m_c	2	2.5	3	3.5	4
E_ν boson	59	83	107	143	178
E_ν fermion	24	41	57	86	115

$$\frac{d^2\sigma^{\nu N}}{dxdy} = \frac{G^2ME_\nu}{\pi} 2x\left[\left(1 - y + \frac{y^2}{2}\right)\left(p(x) + n(x) + \lambda(x) + \frac{\bar{p}(x)+\bar{n}(x)}{2} + 2\bar{p}'(x)\right) + \left(y - \frac{y^2}{2}\right)\left(\lambda(x) - \frac{\bar{p}(x)+\bar{n}(x)}{2}\right)\right], \quad (4)$$

$$\frac{d^2\sigma^{\bar{\nu} N}}{dxdy} = \frac{G^2ME_\nu}{\pi} 2x\left[\left(1 - y + \frac{y^2}{2}\right)\left(\frac{p(x)+n(x)}{2} + \bar{p}(x) + \bar{n}(x) + \bar{\lambda}(x) + 2p'(x)\right) - \left(y - \frac{y^2}{2}\right)\left(\frac{p(x)+n(x)}{2} - \bar{\lambda}(x)\right)\right]. \quad (5)$$

Formulas (4) and (5) are valid at energies exceeding the threshold for the production of charmed particles

$$E_{\text{thresh}} \approx m_c + \frac{m_c^2}{2M} \quad \text{(boson)},$$
$$E_{\text{thresh}} \approx \frac{m_c^2}{2M} - \frac{M}{2} \quad \text{(fermion)}, \quad (6)$$

where m_c is the mass of the charmed particle and M is the mass of the nucleon, neglecting the small contributions from the sea of the parton-antiparton pairs, we get

$$\frac{d^2\sigma^{\nu N}}{dxdy} \simeq \frac{G^2ME_\nu}{\pi} 2x(p(x) + n(x))\left(1 - y + \frac{y^2}{2}\right), \quad (7)$$

$$\frac{d^2\sigma^{\bar{\nu} N}}{dxdy} \simeq \frac{G^2ME_\nu}{\pi} x(p(x) + n(x))(1 - y)^2. \quad (8)$$

To estimate the total cross section, using the form[4] $xp(x), xn(x) \sim (1 - x)^3$, we obtain

$$\sigma^{\nu N} \approx \sigma^{\nu N}_{\text{GIM}}\left[1 + \frac{(1 - y_n)^4}{3}\right], \quad y_n = \frac{E_{\text{thresh}}}{E_\nu}, \quad (9)$$

$$\sigma^{\bar{\nu} N} \approx \sigma^{\bar{\nu} N}_{\text{GIM}}. \quad (10)$$

We note that it follows from (7) that at very small values ($x \gtrsim 0.05$–0.1) the distribution in y in the scheme of[1,2] for the νN collisions takes the form

$$\frac{d\sigma^{\nu N}}{dy} = \frac{G^2ME_\nu}{\pi}\left[I_1 + I_2(1 - y)^2\right], \quad (11)$$

$$I_1 = \int_0^1 x(n(x) + p(x))dx, \quad I_2 = \int_0^{1 - \frac{y_n}{1-y}} x(n(x) + p(x))dx, \quad (12)$$

which differs qualitatively from the predictions of the GIM scheme

$$\frac{d\sigma^{\nu N}_{\text{GIM}}}{dy} \approx \frac{G^2ME_\nu}{\pi} I_1. \quad (13)$$

The total cross section of the νN interaction (9) exceeds by approximately 1/3 the value predicted in the GIM model in the region $E_\nu \gg E_{\text{thresh}}$. In the energy interval $E_\nu = E_{\text{thresh}}$ to $4E_{\text{thresh}}$ there is a smooth variation of the slope of the energy dependence of $\sigma^{\nu N}$. The table lists the energies E_ν starting with which we can expect an appreciable change in the cross section $\sigma^{\nu N}$, by 25%, as a function of the mass m_c of the produced charmed boson or fermion. Present-day experimental data do not permit definite conclusions to be drawn concerning the change in the character of the energy dependence. They do not exclude, however, the possibility of a 10–20% change in the slope of $\sigma^{\nu N}$. It is therefore quite important to carry out precision measurements of the energy dependence of $\sigma^{\nu N}$ or of the ratio $\sigma^{\bar{\nu} N}/\sigma^{\nu N}$ in the range of energies attainable with modern accelerators.

We note also that the requirement that the parton distributions in x be positive, the scheme of[1,2] imposes on the ratio of the $\nu_\mu N$ and $\bar{\nu}_\mu N$ interaction cross sections restrictions other than in the GIM model,

$$\frac{1}{4} \leqslant \frac{\sigma(\bar{\nu} N)}{\sigma(\nu N)} \leqslant 1. \quad (17)$$

In concluding the discussions of the consequences of the weak hadron-current scheme[1,2] as applied to high-energy neutrino physics, we present the Adler sum rule

$$I^p = \int dx(F_1^{\bar{\nu} p} - F_1^{\nu p}) = 0,$$
$$I^n = \int dx(F_1^{\bar{\nu} n} - F_1^{\nu n}) = -3. \quad (18)$$

In the GIM model the Adler sum rule for the nucleus F_e takes the form

$$I^{F_e} = -4, \quad (19)$$

whereas in the scheme of[1,2] we have

$$I^{F_e} = -90, \quad (20)$$

as follows from (18). A check on this rule is also very important for a final choice of the model of weak hadron interactions.

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