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QUASI-ELASTIC NEUTRINO PRODUCTION OF CHARMED BARYONS FROM THE POINT OF VIEW OF LOCAL DUALITY

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The reactions

$$\nu + n \longrightarrow \mu^- + \Lambda_c(2285)^+ \quad (1)$$
$$\nu + n \longrightarrow \mu^- + \Sigma_c(2455)^+ \quad (2)$$
$$\nu + p \longrightarrow \mu^- + \Sigma_c(2455)^{++} \quad (3)$$

provide an important information on weak form factors (FF) of nucleon transitions to charmed baryons. Investigation of these processes can improve or change usually accepted notions of the structure of baryons with a charmed quark. A key issue in planning the experimental investigations is the evaluation of cross sections for the considered processes. Usually it is pointed out that the discrepancy about one order in the predictions of the known models, makes it rather difficult to foresee the experimental statistics.

To our mind, the discrepancies revealed by formal comparison of the predictions of these models are not so crucial. Now it is clear that if SU_4-symmetry breaking is taken into account correctly, then the discrepancies disappear. Problems with the predictions of these models come from another source. We mean the dipole parametrization of Q^2- dependence of the nucleon FF in reactions (1)-(3). The dipole parametrization involves the empirical formulae obtained from the experimental data on the electromagnetic and weak elastic nucleon FF. It has no direct theoretical interpretation. Thus the parameters of these formulae, which are the vector and axial masses cannot be a priori found from theory. In this situation the predictive abilities of the models mentioned essentially weaken.

The aim of the present paper is to evaluate the cross sections of processes (1)-(3) in the framework of the approach [8-10] based on the QCD version [12,13] of Bloom-Gilman (B-G) duality [11] applied to the νN-scattering. We consider this approach as an alternative to the dipole parametrization. The phenomenological parameters of the approach have a clear physical interpretation and can be fixed or bound in values rather reliably. As a result we are able to obtain the upper bound for the cross section of processes (1)-(3).

The QCD version of B-G duality connects the smooth nonsinglet structure function (SF) $F^{th}(x,Q^2)$ of deep inelastic scattering, calculated in the QCD perturbation theory with the observable SF $F^{ph}(x,Q^2)$ possessing the baryon resonance bumps at small Q^2. These two functions coincide with a good precision when averaged both over the whole kinematic interval $0 \le x \le 1$ (global duality):

$$\int_0^1 (F^{ph}(x,Q^2) - F^{th}(x,Q^2))dx \simeq 0 \qquad (4)$$

and over some vicinity of a certain resonance $x_{min} \le x \le x_{max}$ (local duality):

$$\int_{x_{min}}^{x_{max}} (F^{ph}(x,Q^2) - F^{th}(x,Q^2))dx \simeq 0. \qquad (5)$$

The nonsinglet SF $F^{th}(x,Q^2)$ describes the nucleon transitions to hadronic final states with the definite electric charge and SU_4-quantum numbers but with non fixed quantum numbers of the space-time symmetries. In its turn SF $F^{ph}(x,Q^2)$ contains only the contributions of the resonances with the internal quantum numbers of these final states.

At the moderate and low values of Q^2 relation (5) is strongly violated because of kinematic and non-perturbative power corrections to $F^{th}(x,Q^2)$. But if we replace the simple Bjorken variable $x=Q^2/2M\nu$ by a more complex one

$$\bar{\xi}(\nu,Q^2) = \xi(\nu,Q^2)\ [1 + \frac{M_o^2}{Q^2}\ (1 + \frac{M_o^2}{Q^2+ M_o^2})], \qquad (6)$$

which takes into account[8] a part of the dominating power corrections to $F^{th}(x,Q^2)$, we can essentially improve the precision of relation (5). In formula (6) the parameter M_o is the scale of non-perturbative (twist-) power corrections like $(M_o^2/Q^2)^n$; ξ is the well known ξ-scaling variable[12,13]

$$\xi(\nu,Q^2) = \frac{Q^2/\ M}{\nu + \sqrt{\nu^2+ Q^2}} \qquad (7)$$

absorbing the kinematic power corrections like $(M^2/Q^2)^n$ with the nucleon mass M as a scale.

When a zero width approximation for the resonance contribution to $F^{ph}(x,Q^2)$ is used, then in relation (5) this function should be replaced by $\mathcal{Z}^R\ F^{ph}(x,Q^2)$. Here $\mathcal{Z}^R$ are constants different for different resonances and not equal to 1 in general case. This replacement was explained in[10] from the point of view of Wilson operator product expansions using the results and methods of the works[11,12].

We shall not dwell on substantiation of the approach . These is the subject of the papers[8-10]. Here we only derive the final formulae for the cross sections in the zero width approximation for the resonance contribution. To do this let's substitute SF $F^{ph}(x,Q^2)$ to relation (5) in the form

$$F_1^{ph}(x,Q^2) = \bar{W}_1(Q^2)\delta(\nu - \nu_R)$$
$$F_{2,3}^{ph}(x,Q^2) = \frac{\nu}{2\ M}\ \bar{W}_{2,3}(Q^2)\delta(\nu - \nu_R)\ , \qquad (8)$$

using a new variable (6) and taking into account the normalization factor $\mathcal{Z}^R$. Then we obtain the resonance functions $\bar{W}_k$ in terms of smooth function $F^{th}(x,Q^2)$ which is well controlled by the QCD perturbation theory and the modified QCD parton model. Substituting SF $F^{ph}(x,Q^2)$ in this form to the known formula for the νN-scattering cross section

$$\frac{d\ \sigma}{d\nu\ dQ^2} = \frac{G^2}{2\pi}\ [\frac{Q^2}{4ME^2}\ F_1^{ph} + \frac{4E(E - \nu) + Q^2}{4E^2\nu}\ F_2^{ph} - \frac{Q^2(2E-\nu)}{4M\nu E^2}\ F_3^{ph}\ ] \qquad (9)$$

and integrating over the variable ν we finally obtain:

$$\frac{d\ \sigma^R}{d\ Q^2} = \frac{G^2}{2\pi}\ \mathcal{Z}^R\ \mathcal{D}^R(Q^2)\ \ [1 - \frac{\nu_R}{E} + \frac{Q^2}{4E^2} + \frac{Q^2}{2ME\bar{\xi}_R}\ ]\ \frac{\sqrt{\nu_R^2+ Q^2}}{\nu_R\ \bar{\xi}_R}\ \theta(E - E^R_{sr}), \qquad (10)$$

where

$$\mathcal{D}^R(Q^2) = \int_{\bar{\xi}_+(Q^2)}^{\bar{\xi}_-(Q^2)} t\ f(t,Q^2)dt \qquad (11)$$

$$\nu_R = \frac{M_R^2 - M^2+ Q^2}{2M}\ ; \qquad \nu_R^{\pm} = \frac{[(M_R\pm\ \Delta_R^{\pm}\ )^2 - M^2+ Q^2]}{2M}$$
$$\bar{\xi}_R(Q^2) = \bar{\xi}(\nu_R,Q^2), \qquad \bar{\xi}_{\pm}(Q^2) = \bar{\xi}(\nu_R^{\pm},Q^2), \qquad (12)$$
$$E^R_{sr} = \frac{(M_R+ m_\mu)^2- M^2}{2M}, \qquad G^2/(2\pi) \simeq 0.8\times10^{-38}sM^2\ GeV^{-2}\ ;$$

E is the energy of the initial neutrino; M, m_μ are the masses of the nucleon and the μ-meson; E^R_{sr} is the threshold energy of the resonance $R = (\Lambda_c^+, \Sigma_c^+, \Sigma_c^{++})$ production; M_R is the mass of this resonance; $\Delta_R^\pm$ are the mass parameters which define the duality interval $\xi \le [\xi_+, \xi_-]$, where as it follows from relation (5) the considered resonance is averaged by the smooth function $F^{th}(x,Q^2)$.

The function $f(x,Q^2)$ is a distribution function (DF) of the valence quark in an initial nucleon. An external current induces a transition of this quark to the final one forming a considered baryon with two other quarks-spectators. This point is explained in fig.1. For reactions (1)-(3) $f(x,Q^2)$ is a DF of the valence d-quark. It obeys QCD evolution equations and can be found if the initial condition $f(x,Q_o^2)$ at fixed $Q^2 = Q_o^2$ is introduced. However, for the estimations that we wish to obtain here it is enough to accept approximately $f(x,Q^2)=f(x,Q_o^2)$ and don't take into account logarithmic Q^2- dependence of DF $f(x,Q^2)$ that corresponds to the QCD corrections. Let's take this function in the form:

$$f(x) = \frac{x^{-1/2}(1-x)^\tau}{B(1/2,\tau+1)} . \tag{13}$$

We use the $\tau=4$ which is generally accepted now and follows from the different experimental data analyses[14-15].

To fix the normalization constant Z^R we use the predictions of SU_4-symmetry of strong interactions violated by a large mass difference of charmed and non-charmed quarks. The phenomenological form factors $F_{1,2}(Q^2)$, $F_A(Q^2)$ (FF) describing the nucleon transitions $N \longrightarrow \Lambda_c^+, \Sigma_c^+, \Sigma_c^{++}$ obey the following relations at $Q^2=0$[16]:

- reaction (1)

$$\begin{aligned} &F_1(0) = \sqrt{3/2}\; F_1^p(0); \quad F_2(0) = \sqrt{3/2}\; \mu_p F_2^p(0); \\ &F_A(0) = \sqrt{3/2}\; [\, f_A(0) + 1/3\, d_A(0)]; \end{aligned} \tag{14}$$

- reaction (2)

$$\begin{aligned} &F_1(0) = -\frac{F_1^p(0)}{\sqrt{2}} \; ; \quad F_A(0) = -\frac{1}{\sqrt{2}}\, [\, f_A(0) - d_A(0)\,]; \\ &F_2(0) = -\frac{1}{\sqrt{2}}\, [\, \mu_p F_2^p(0) + 2\, \mu_n F_2^n(0)\,] . \end{aligned} \tag{15}$$

For reaction (3) the right hand side of relations (15) have to be multiplied by $\sqrt{2}$. Here $F_1^{p,n}$; $F_2^{p,n}$ are the electromagnetic FF of the proton and the neutron. $f_A(Q^2)$, $d_A(Q^2)$ are the FF of the axial-vector f,d-coupling of the nucleon. The values of the constants introduced

are

$F_{1,2}^{p,n}(0) = 1$; $\mu_p = 1{,}79$; $\mu_n = -1{,}91$;

$f_A(0) = 0{,}3$; $d_A(0) = 0{,}95$.

Writing down the cross section at $Q^2=0$ with above-mentioned FF and relations (14)-(15)

$$\left.\frac{d\sigma}{dQ^2}\right|_{Q^2=0} = \frac{G^2 \sin^2\theta_c\, (2ME + M^2 - M_R^2)}{4\pi\, ME}\, [F_1^2(0) + F_A^2(0)] \tag{16}$$

then comparing it with (10), we obtain:

$$Z^R = \frac{2M_o^2 \sin^2\theta_c\, [F_1^2(0) + F_A^2(0)]}{D^R(0)(\, M_R^2 - M^2)} . \tag{17}$$

Now we are ready to consider quantitative predictions of the dual approach for cross sections of reactions (1)-(3). The approach contains three phenomenological parameters M_o, $\Delta_R^\pm$. M_o is a proper scale of internal nucleon dynamics which we consider to be equal for different processes. It can be fixed from the analysis[9] of the data on cross section $\sigma(\nu n \longrightarrow \mu p)$ and electromagnetic FF. The result is:

$$M_o = 0{,}08 \pm 0{,}02 \text{ GeV}. \tag{18}$$

The parameters $\Delta_R^\pm$ obey the constraints $\Delta_R^\pm \le |(M_{R\pm} - M_R)|$, where $M_{R-} \le M_R \le M_{R+}$ are the masses of the neighboring resonances $R^\pm$ which are the closest to R in masses and possess the same the electric charge, SU_4-quantum numbers as R has. For the considered baryons Λ_c^+, Σ_c^+, Σ_c^{++} we have:

$$\begin{aligned} &\Lambda_c^+:\ \ \xi_- = 1, \quad \Delta^+ \le M_{B_o^1} - M_{\Lambda_c}; \quad \Sigma_c^+:\ \xi_- = 1, \quad \Delta^+ \le M_{C_1^*} - M_{\Sigma_c}; \\ &\Sigma_c^{++}:\ \xi_- = 1, \quad \Delta^+ \le M_{C_1^*} - M_{\Sigma_c} . \end{aligned} \tag{19}$$

Here B_o^{1+} is the first excited charmed baryon state with the orbital momentum L=1, C_1^{*+}, C_1^{*++} are the L=0 ground-state charmed baryons with the spin S = 3/2. We estimate their masses following the method of the work[4], but using modern values of the masses of the baryons Λ_c^+, Σ_c^+, Σ_c^{++}[17]. The result is:

$M_{B_o^1} \simeq 2{,}84$ GeV; $\qquad M_{C_1^*} \simeq 2{,}64$ GeV.

When the scattering channel with quantum numbers of one of the baryons Λ_c^+, Σ_c^+ or Σ_c^{++} is considered, there is only one resonance bump inside the interval $\xi_+ \le \xi \le 1$, corresponding to the contribution of

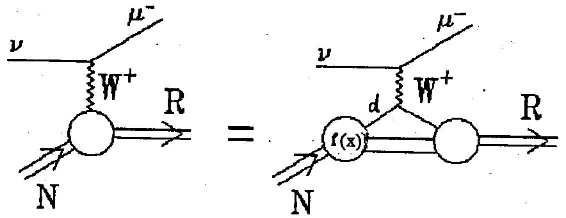

Fig.1. Quark diagram for charmed baryon neutrino production.

this baryon. The reason is that this baryon has the lowest mass in this channel . That's why we accept $\bar{\xi}=1$ in (19) as an upper limit of the duality interval.

With the help of (18),(19), the estimation of the upper bounds for the cross sections of reactions (1)-(3) can be found. To make this we use the following parameter values

$$M_0 = 0{,}1 \text{ GeV} ; \quad \Delta^+_{\Lambda_c} = 0{,}56 \text{ GeV}; \quad \Delta^+_{\Sigma_c} = 0{,}2 \text{ GeV}. \qquad (20)$$

The final results of our calculations for the upper bounds of the differential and total cross sections are presented in fig.2,3.

In conclusion we note that the obtained predictions of the dual approach are consistent with the results of works /1-5/, but lie essentially lower than the results of works /2-4/.

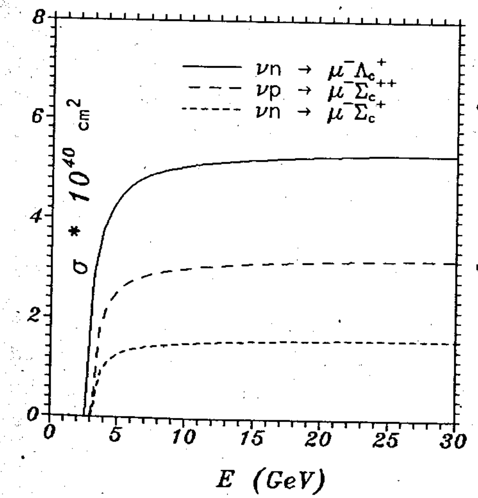

Fig.2. Upper bounds for total cross sections of charmed baryons neutrino production.

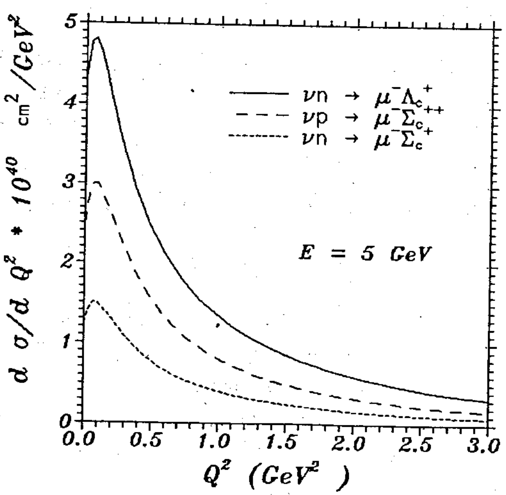

Fig.3 Upper bounds for differential cross sections of charmed baryons neutrino production.

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