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Heavy quark currents in Ultra-High Energy Neutrino Interactions

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We discuss heavy quark contributions to the neutrino-nucleon total cross section at very high energies, well above the real top production threshold. The top-bottom weak current is found to generate strong left-right asymmetry of neutrino-nucleon interactions. We separate contributions of different helicity states and make use of the $\boldsymbol{\kappa}$-factorization to derive simple and practically useful formulas for the left-handed (F_L) and right-handed (F_R) components of the conventional structure function $2xF_3 = F_L - F_R$ in terms of the integrated gluon density. We show that $F_L \gg F_R$ and, consequently, $xF_3 \approx F_T$, where F_T is the transverse structure function. The conventional structure function $F_2 = F_S + F_T$ at $Q^2 \ll m_t^2$ appears to be dominated by its scalar (also known as longitudinal) component F_S and the hierarchy $F_S \gg F_L \gg F_R$ arises naturally. We evaluate the total neutrino-nucleon cross section at ultra-high energies within the color dipole BFKL-formalism.

The neutrino astrophysics [1] deals with neutrinos of ultra-high energy (UHE), well above the threshold of the top quark production. The role of heavy quark currents in UHE neutrino interactions has been explored extensively [2]. Different approaches to including the top quark effects in the UHE neutrino-nucleon cross section $\sigma^{\nu N}$ based on [2] are in use [3]. The goal of this communication is to show that a specific choice of relevant degrees of freedom greatly simplifies the analysis of $\sigma^{\nu N}$ in terms of the nucleon structure functions (SF). Making use of the $\boldsymbol{\kappa}$-factorization we isolate the leading contributions involving big log's of Q^2, m_t^2, and $1/x$ to derive simple and practically useful formulas for the UHE neutrino-nucleon SFs.

The UHE interactions correspond to neutrino energy above $E_\nu \sim 10^6$ GeV. The overall hardness scale of the process

$$\nu N \to \mu X \tag{1}$$

is determined by the gauge boson mass

$$Q^2 \sim m_W^2. \tag{2}$$

Therefore, the W-boson exchange probes the gluon density in the target nucleon at very small values of Bjorken x. Consequences of this observation for $\sigma^{\nu N}(E_\nu)$ are widely discussed [4] (for most recent publications see [5]).

At small x, to the Leading Log $(1/x)$ approximation, it is legitimate to consider the W-nucleon scattering in the laboratory frame in terms of interactions with the target of the $q\bar{q}\prime$-pair (color dipole) which the light-cone W-boson transforms into at large upstream distances. The dynamics of the $\log(1/x)$-evolution [6] of the dipole-nucleon cross section is described by the infrared regulated BFKL-equation with running coupling derived in [7, 8].

The differential cross section for the νN interactions is expressible in terms of the scalar (also known as longitudinal), left-handed and right-handed structure functions denoted by F_S, F_L, and F_R, respectively.

$$x\frac{d\sigma^{\nu N}}{dxdQ^2} = \frac{G_F^2}{4\pi}\left(\frac{m_W^2}{m_W^2+Q^2}\right)^2 \times \\ \times \left[2(1-y)F_S + F_L + (1-y)^2 F_R\right]. \tag{3}$$

The structure functions F_λ are related to the absorption cross sections for the W-boson in the polarization state $\lambda = S, L, R$,

$$F_\lambda(x, Q^2) = \frac{Q^2}{4\pi^2\alpha_W}\sigma_\lambda(x, Q^2), \tag{4}$$

where $\alpha_W = g^2/4\pi$, $g^2 = G_F m_W^2/\sqrt{2}$ and the $W \to t\bar{b}$ transition vertex is

$$g\bar{t}\gamma_\mu(g_V - g_A\gamma_5)b.$$

In terms of F_L, F_R the transverse structure function $F_T = F_2 - F_S$ reads $F_T = (F_L + F_R)/2$, and left-right antisymmetric structure function F_3 is as follows $2xF_3 = F_L - F_R$. Switching L to R in Eq. (3) yields the $\bar{\nu}N$ cross section.

Then, making use of the color dipole factorization [9] we arrive at the leading-log$(1/x)$ $\boldsymbol{\kappa}$-factorization representation for the absorption cross section for the W-boson in the polarization state λ

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$$\sigma_\lambda(x,Q^2)=\int_0^1 dz\int d^2\mathbf{r}|\Psi_\lambda(z,r)|^2\sigma(x,r)=$$
$$=\frac{\alpha_W}{\pi}\int_0^1 dz\int d^2\mathbf{k}\int\frac{d^2\boldsymbol{\kappa}}{\boldsymbol{\kappa}^4}\alpha_S(q^2)\mathcal{F}(x_g,\boldsymbol{\kappa}^2)\times$$
$$\times\left(S_\lambda\Phi_0^2+P_\lambda\mathbf{\Phi}_1^2\right). \qquad (5)$$

We address first the second line of Eq. (5)[2)] to derive some useful analytical expressions for $F_\lambda(x,Q^2)$. The first line of Eq. (5) will be used in our numerical studies of $\sigma^{\nu N}(E_\nu)$.

No restrictions on the quark transverse momentum, $\mathbf{k}$, and the gluon transverse momentum, $\boldsymbol{\kappa}$ are imposed in Eq. (5), where

$$\mathcal{F}(x_g,\boldsymbol{\kappa}^2)=\frac{\partial G(x_g,\boldsymbol{\kappa}^2)}{\partial\log\kappa^2},$$

and $G(x_g,\boldsymbol{\kappa}^2)=x_g g(x_g,\boldsymbol{\kappa}^2)$ is the integrated gluon structure function. It is taken at

$$x_g=\frac{Q^2+M^2}{W^2+Q^2} \qquad (6)$$

and the transverse mass of the $t\bar{b}$ pair is

$$M^2=\frac{m_t^2+\mathbf{k}^2}{z}+\frac{m_b^2+(\mathbf{k}-\boldsymbol{\kappa})^2}{1-z}, \qquad (7)$$

where z is the fraction of the W-boson light cone momentum carried by the t-quark.

The terms proportional to S_λ and P_λ describe interaction of the quark-antiquark states with the angular momentum $L=0$ (S-wave) and $L=1$ (P-wave), respectively. The S-wave and the P-wave factors in Eq. (5) for the right-handed and left-handed W-bosons are as follows [11]

$$S_R=\{g_V[(1-z)m_t+zm_b]+g_A[(1-z)m_t-zm_b]\}^2,$$
$$P_R=(g_V-g_A)^2z^2+(g_V+g_A)^2(1-z)^2,$$
$$S_L=\{g_V[(1-z)m_t+zm_b]-g_A[(1-z)m_t-zm_b]\}^2,$$
$$P_L=(g_V-g_A)^2(1-z)^2+(g_V+g_A)^2z^2, \qquad (8)$$

and for the scalar/longitudinal polarization [12, 11]

$$S_S=(g_V^2/Q^2)\times$$
$$\times\left\{2Q^2z(1-z)+(m_t-m_b)\left[(1-z)m_t-zm_b\right]\right\}^2+$$
$$+(g_A^2/Q^2)\times$$
$$\times\left\{2Q^2z(1-z)+(m_t+m_b)\left[(1-z)m_t+zm_b\right]\right\}^2,$$
$$P_S=(g_V^2/Q^2)(m_t-m_b)^2+(g_A^2/Q^2)(m_t+m_b)^2. \qquad (9)$$

[2)]To compare Eq. (5) with Eq. (2) of Ref. [10] one should make the substitution:

$$\frac{4\alpha_S}{\pi}\frac{V(\kappa)\kappa^4}{(\kappa^2+\mu_G^2)^2}\to\mathcal{F}(x_g,\boldsymbol{\kappa}^2).$$

The numerical factor 16 in Eq. (2) of Ref. [10] should be understood as 4.

In the charged current interactions $g_A=g_V=1$. In the neutral current neutrino interactions $m_q=m_{\bar{q}'}$ and corresponding vector and axial-vector couplings are given by the Standard Model.

In Eq. (5)

$$\Phi_0=\left[\frac{1}{\mathbf{k}^2+\varepsilon^2}-\frac{1}{(\mathbf{k}-\boldsymbol{\kappa})^2+\varepsilon^2}\right],$$
$$\mathbf{\Phi}_1=\left[\frac{\mathbf{k}}{\mathbf{k}^2+\varepsilon^2}-\frac{\mathbf{k}-\boldsymbol{\kappa}}{(\mathbf{k}-\boldsymbol{\kappa})^2+\varepsilon^2}\right], \qquad (10)$$

and

$$\varepsilon^2=z(1-z)Q^2+(1-z)m_t^2+zm_b^2. \qquad (11)$$

Then, making use of the technique developed in [13] for electro-production processes, separate the $\boldsymbol{\kappa}^2$-integration in (5) into the soft,

$$\boldsymbol{\kappa}^2\lesssim\overline{k^2}\equiv\varepsilon^2+\mathbf{k}^2, \qquad (12)$$

and hard,

$$\boldsymbol{\kappa}^2\gtrsim\overline{k^2}, \qquad (13)$$

regions of the gluon momentum. For soft gluons, in the DGLAP [14] region, upon the azimuthal integration we get

$$\int\frac{d\varphi}{2\pi}\mathbf{\Phi}_1^2\approx\boldsymbol{\kappa}^2\frac{\varepsilon^4+(\mathbf{k}^2)^2}{(\mathbf{k}^2+\varepsilon^2)^4},$$
$$\int\frac{d\varphi}{2\pi}\Phi_0^2\approx\frac{2\boldsymbol{\kappa}^2\mathbf{k}^2}{(\mathbf{k}^2+\varepsilon^2)^4}. \qquad (14)$$

In Eq. (5) the QCD running coupling $\alpha_S(q^2)$ enters the integrand at the largest relevant virtuality, $q^2=\max\{\overline{k^2},\boldsymbol{\kappa}^2\}$. For soft gluons one can take $q^2=\overline{k^2}$ and we arrive at the fully differential distribution of the t-quark in z and $\mathbf{k}$,

$$\frac{d\sigma_\lambda(x,Q^2)}{dzd^2\mathbf{k}}\approx\alpha_W\alpha_S(\overline{k^2})G(x_g,\overline{k^2})\times$$
$$\times\left[S_\lambda\frac{2\mathbf{k}^2}{(\mathbf{k}^2+\varepsilon^2)^4}+P_\lambda\frac{\varepsilon^4+(\mathbf{k}^2)^2}{(\mathbf{k}^2+\varepsilon^2)^4}\right], \qquad (15)$$

where $\overline{k^2}$ comes from Eq. (12). Integrating over $\mathbf{k}$ yields

$$\frac{d\sigma_\lambda(x,Q^2)}{dz}\approx\frac{2\pi\alpha_W}{3\varepsilon^2}\left(\frac{S_\lambda}{2\varepsilon^2}+P_\lambda\right)\alpha_S(\varepsilon^2)G(x_g,\varepsilon^2). \qquad (16)$$

The leading contribution to Eq. (16) comes from $\mathbf{k}^2\lesssim\varepsilon^2/3$. We factored out the slowly varying product $\alpha_S G$ at $\overline{k^2}\approx\varepsilon^2$. The $\mathbf{k}$- and $\boldsymbol{\kappa}$-dependence of x_g is given by Eqs. (6), (7).

Then, going from $d\sigma_L/dz$ to the structure function of the nucleon probed by the left-handed ($\lambda=L$) gauge boson, $F_L(x,Q^2)$ (see Eq. (4)), we find the soft gluon contribution to the P-wave component[3)] of

$$F_L = F_L^1 + F_L^0 \tag{17}$$

denoted by F_L^1

$$F_L^1(x,Q^2) \approx 2Q^2 \int_0^1 \frac{dz z^2}{\varepsilon^2} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2). \tag{18}$$

The leading contribution to F_L^1 comes from

$$z \sim 1 - \frac{m_b^2}{m_t^2 + Q^2}. \tag{19}$$

Therefore, Eq. (18) can be approximated by

$$F_L^1(x,Q^2) \approx \frac{2Q^2}{m_t^2+Q^2} \times \\ \times \int_{m_b^2}^{C(m_t^2+Q^2)} \frac{d\varepsilon^2}{\varepsilon^2} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2), \tag{20}$$

where $C \approx 0.25$. Certainly, $\sigma^{\nu N}$ is dominated by $Q^2 \ll m_t^2$. However, F_L^1 presented in the form (20) allows straightforward extension to the processes induced by the charm-strange current, where $Q^2 \gg m_c^2$ and $x_g \sim 2x$.

The S-wave component of F_L, denoted by F_L^0, is very small compared to F_L^1. Indeed, the z-integration in F_L^0 converges rapidly at $z \sim 1 - m_b^2/(m_t^2+Q^2)$ and

$$F_L^0(x,Q^2) \approx Q^2 m_b^2 \int_0^1 \frac{dz z^2}{\varepsilon^4} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2) \approx \\ \approx \frac{Q^2}{m_t^2+Q^2} \frac{\alpha_S(m_b^2)}{3\pi} G(x_g,m_b^2). \tag{21}$$

Therefore, $F_L^0 \ll F_L^1$ and $F_L \approx F_L^1$.

The S-wave component of the right-handed structure function

$$F_R = F_R^0 + F_R^1 \tag{22}$$

is as follows

$$F_R^0(x,Q^2) \approx Q^2 m_t^2 \int_0^1 \frac{dz(1-z)^2}{\varepsilon^4} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2) \approx \\ \approx \frac{Q^2}{m_t^2} \frac{\alpha_S(m_t^2/2)}{3\pi} G(x_g,m_t^2/2). \tag{23}$$

At $Q^2 \ll m_t^2$ the ratio $(1-z)^2/\varepsilon^4$ in Eq. (23) is flat in z and $\langle z \rangle \sim 1/2$.

[3)]Hereafter, the upper index in F_λ^L corresponds to the angular momentum $L = 0,1$ of the $q\bar{q}'$-state which the light-cone W-boson transforms into.

For $Q^2 \gg m_t^2$ the structure function F_R^0 is dominated by $z \sim m_t^2/Q^2$ and $x_g \approx 2x$. Hence,

$$F_R^0(x,Q^2) \approx \frac{\alpha_S(2m_t^2)}{3\pi} G(x_g, 2m_t^2). \tag{24}$$

The P-wave component of F_R reads

$$F_R^1(x,Q^2) \approx 2Q^2 \int_0^1 \frac{dz(1-z)^2}{\varepsilon^2} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2). \tag{25}$$

For $Q^2 \ll m_t^2$, $(1-z)^2/\varepsilon^2 \approx (1-z)/m_t^2$ and

$$F_R^1(x,Q^2) \approx \frac{Q^2}{m_t^2} \frac{\alpha_S(m_t^2)}{3\pi} G(x_g,m_t^2). \tag{26}$$

In the region of very high virtualities of the gauge boson, $CQ^2 \gg m_t^2$ (see Eq. (20)),

$$F_R^1(x,Q^2) \approx 2 \int_{m_t^2}^{CQ^2} \frac{d\varepsilon^2}{\varepsilon^2} \frac{\alpha_S(\varepsilon^2)}{3\pi} G(x_g,\varepsilon^2) \tag{27}$$

with $x_g \approx 2x$. Certainly, Eq. (27) is irrelevant to the problem of $\sigma^{\nu N}$: the latter is dominated by $Q^2 \ll m_t^2$. However, Eq. (27) is not entirely useless. It describes – with obvious substitutions – the dominant contribution to the F_R^1 in high-Q^2 processes induced by the charm-strange current (see Fig. 1).

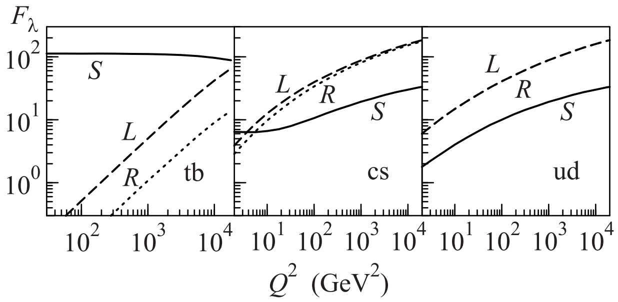

Fig. 1. The nucleon structure function F_λ for $\lambda = S, L, R$ as a function of Q^2 at $x = 5 \cdot 10^{-7}$ in the neutrino reactions induced by the top-bottom, charm-strange and up-down quark currents

From explicite expressions for F_L and F_R obtained above it is evident that for the top-bottom current

$$F_L \gg F_R \tag{28}$$

in a wide range of Q^2. Another observation is that the contribution of left-handed W-bosons to $\sigma^{\nu N}$ in processes induced by the top-bottom current is much smaller than that coming from the absorption of scalar W-bosons

$$F_S \gg F_L. \tag{29}$$

The scalar structure function F_S (under the name longitudinal) has been discussed in [15]. It was found that the higher twist corrections brought about by the non-conservation of the top-bottom current result in considerable enhancement of the P-wave component of

$$F_S = F_S^1 + F_S^0, \tag{30}$$

denoted by F_S^1

$$F_S^1(x,Q^2) \approx \frac{m_t^2}{m_t^2+Q^2}\int_{m_b^2}^{\varepsilon_m^2}\frac{d\varepsilon^2}{\varepsilon^2}\frac{\alpha_S(\varepsilon^2)}{3\pi}G(x_g,\varepsilon^2), \tag{31}$$

where $\varepsilon_m^2 = m_t^2$ if $Q^2 < m_t^2 - m_b^2$ and $\varepsilon_m^2 \approx (Q^2+m_t^2)/4$ if $Q^2 > m_t^2 - m_b^2$.

At $Q^2 \ll m_t^2$, F_S^0 is much smaller than F_S^1. Indeed, from Eq. (16) it follows that

$$F_S^0 \approx \frac{Q^2}{12\pi}\int_0^1 \frac{dz S_S}{\varepsilon^4}\alpha_S(\varepsilon^2)G(x_g,\varepsilon^2), \tag{32}$$

where S_S comes from Eq. (9) and ε^2 from Eq. (11). For $Q^2 \ll m_t^2$

$$\frac{Q^2 S_S}{\varepsilon^4} \approx 2\left[1+\frac{\delta^2}{(1-z+\delta^2)^2}\right] \tag{33}$$

with $\delta = m_b/m_t$. Therefore,

$$F_S^0(x,Q^2) \approx \frac{\alpha_S(m_t^2/2)}{6\pi}G(x_g,m_t^2/2) + \\ + \frac{\alpha_S(m_b^2)}{6\pi}G(x_g,m_b^2). \tag{34}$$

Our F_S survives the limit $m_t^2 \to \infty$[4]. The point is that the scalar/longitudinal W-boson interacts with the quark current $j_\mu = V_\mu - A_\mu$ which is not conserved. The vertex $W \to t\bar{b}$ which is $\propto \partial_\mu j_\mu \propto m_t \pm m_b$ gives the factor m_t^2 in the expression for F_S and cancels $1/m_t^2$ coming from the tb quark box. Evidently, there is no real clash between Eq. (34) and the decoupling theorem [16]. The latter is relevant to calculating observables, for instance cross sections in kinematical regions well below certain heavy quark thresholds. We are dealing with the real top production by neutrinos of ultra-high energy, though at $Q^2 \ll m_t^2$.

Two terms in Eq. (34) correspond to two very different kinematics. The first one describes the symmetric $t\bar{b}$-state with uniform z-distribution and $\langle z\rangle \approx 1/2$, the second one corresponds to the asymmetric configuration with $z \approx 1-\delta^2$. Both $t\bar{b}$-states have approximately the same invariant mass, $M^2 \sim 2m_t^2$, and, consequently, probe the gluon density at nearly the same x_g.

For $Q^2 \gg m_t^2$ – the substitution $t \to c$ suggests itself – the ratio S_S/ε^4 is flat in z and $Q^2 S_S/\varepsilon^4 \approx 8$ with $\langle z\rangle \approx 1/2$ and $x_g \approx 2x$. Consequently,

$$F_S^0(x,Q^2) \approx \frac{2\alpha_S(Q^2/4)}{3\pi}G(x_g,Q^2/4). \tag{35}$$

In the neutrino reactions induced by the charm-strange current the term F_S^1 is suppressed at high $Q^2 \gg m_c^2$ by the factor $\sim m_c^2/Q^2$ and the scalar/longitudinal structure function F_S is dominated by its S-wave component, F_S^0 [17, 18]. Specific features of the structure functions F_λ are illustrated by Fig. 1. From Fig. 1 and Eqs. (3), (2) it follows that the left-right asymmetry generated by the top-bottom current is much stronger than that generated by the charm-strange current but the top-bottom contribution to the total neutrino nucleon cross section is dominated by the absorption of the scalar gauge boson.

In our numerical studies of F_λ and $\sigma^{\nu N}$ we rely upon the color dipole factorization represented by the first line of Eq. (5), where the total cross section of interaction of the $t\bar{b}$ color dipole of the transverse size $\mathbf{r}$ with the nucleon target is related to the differential density of gluons $\mathcal{F}(x_g,\kappa^2)$ by the equation [9]

$$\sigma(x,r) = \\ = \frac{\pi r^2}{N_c}\int\frac{d^2\boldsymbol{\kappa}}{\kappa^2}\frac{4[1-\exp(i\boldsymbol{\kappa}\mathbf{r})]}{\kappa^2 r^2}\alpha_S(\kappa^2)\mathcal{F}(x_g,\kappa^2). \tag{36}$$

The light cone density of the $t\bar{b}$ Fock states with the transverse size $\mathbf{r}$ and the fraction z of the W-boson light cone momentum carried by the t-quark is [11, 12]

$$|\Psi_\lambda(z,r)|^2 = \frac{2\alpha_W N_c}{(2\pi)^2}\left[S_\lambda K_0^2(\varepsilon r) + P_\lambda \varepsilon^2 K_1^2(\varepsilon r)\right], \tag{37}$$

where S_λ and P_λ come from Eqs. (8), (9). In (37) $K_{0,1}(y)$ is the modified Bessel function. The terms proportional to $K_0^2(\varepsilon r)$ and $K_1^2(\varepsilon r)$ describe the quark-antiquark states with the angular momentum $L=0$ (S-wave) and $L=1$ (P-wave), respectively. The $\log(1/x)$-evolution of $\sigma(x,r)$ is determined by the infrared regulated BFKL-equation with running coupling [7, 8]. The preferred choice of the infrared regularization gives the intercept of the pomeron trajectory, $\alpha_{\mathbf{IP}}(t)$, in the angular momentum plane $\Delta_{\mathbf{IP}} = \alpha_{\mathbf{IP}}(0) - 1 = 0.4$ and leads to a very good description of the data on the proton structure functions at small x [19]. In particular, this

[4] The limit $m_t^2 \to \infty$ implies the limit of the infinite neutrino energy, $1/x_g \to \infty$. The finite energy effects, built in the gluon density function , suppress the structure function in Eq. (34). Specifically, for $x_g \to 1$ the gluon density G is known to vanish as $(1-x_g)^n$, where $n \approx 5$ and x_g comes from Eqs. (6), (7).

means that the gluon density and all structure functions grow fast with growing $1/x$, $F_\lambda \propto x^{-\Delta_{\mathbf{IP}}}$. The analysis of the unitarity effects will be published elsewhere. The structure functions of the isoscalar nucleon probed by W-bosons of different helicity are presented in Fig. 1.

The charged current neutrino-nucleon cross section, $\sigma^{\nu N}$ as a function of the neutrino energy, E_ν and its decomposition into components of different origin is shown in Fig. 2.

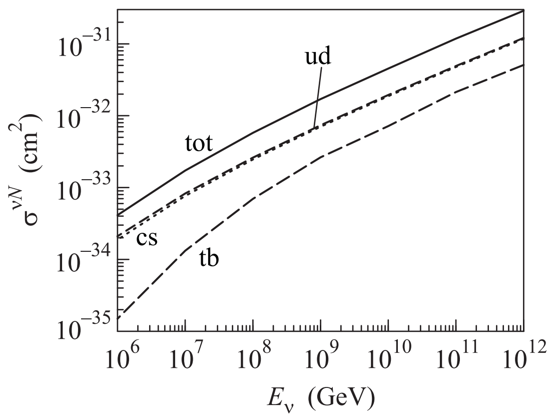

Fig. 2. The total νN charged current cross section (solid line) decomposed into components of different origin as a function of the laboratory neutrino energy

Summarizing, we presented the $\boldsymbol{\kappa}$-factorization formulas for the differential cross section $d\sigma_\lambda/dzd^2\mathbf{k}$ which describes the absorption of scalar, left-handed and right-handed W-bosons. We isolated leading contributions to the related structure functions $F_\lambda(x, Q^2)$ at low x and high Q^2 and obtained simple and numerically accurate estimates for F_λ in processes induced by massive quark currents. It was shown that the non-conservation of the top-bottom current leads to the hierarchy of structure functions $F_S \gg F_L \gg F_R$. Making use of the color dipole BFKL approach to the $\log(1/x)$ QCD-evolution we evaluated the charged current νN cross sections for ultra-high energy neutrino beams.

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