Left-right asymmetry of nuclear shadowing in charged current DIS

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Abstract

We study the shadowing effect in highly asymmetric diffractive interactions of left-handed and right-handed W-bosons with atomic nuclei. The target nucleus is found to be quite transparent for the charmed-strange Fock component of the light-cone $W^+$ in the helicity state $\lambda = +1$ and rather opaque for the $c\bar{s}$ dipole with $\lambda = -1$. The shadowing correction to the structure function $\Delta xF_3 = xF_{3N}^\nu - xF_{3N}^{\bar{\nu}}$ extracted from $\nu Fe$ and $\bar{\nu} Fe$ data is shown to make up about 20% in the kinematical range of CCFR/NuTeV.
1 Introduction

There are several well-established facts about nuclear shadowing in photo- and electro-production, the phenomenon expected since long ago [1, 2, 3, 4]:

- the phenomenon does exist [5],
- it scales [6],
- it is dominated by large, hadronic size color $q\bar{q}$-dipoles, $r \sim m_q^{-1}$, that the virtual photon transforms into at a large distance upstream the target [6].

The latter point implies that shadowing in photo- and electro-production is $\propto m_q^{-2}$ and is rather small in the charm structure function of nuclei.

The situation is quite different in the charged current (CC) neutrino deep-inelastic scattering (DIS). At small values of Bjorken $x$ the charm production in CC DIS is driven by the $W^+$-gluon fusion,

$$W^+g \rightarrow c\bar{s}. \tag{1}$$

In the process (1) charm is inseparable from strangeness and the size of the relativistic heavy-light color dipole $c\bar{s}$ depends strongly on the momentum partition in the light-cone $W^+$-decay. The latter is determined by the light-cone wave function (LCWF) of the Fock state $|c\bar{s}\rangle$ and depends on the helicity of $W^+$-boson [7]. In this communication we demonstrate that for the left-handed $W^+$ shadowing is $\propto 1/\mu^2$ and for the right-handed $W^+$ shadowing is $\propto 1/m^2$, where $\mu$ and $m$ are the strange and charmed quark masses, respectively. Therefore, in spite of the presence of heavy quark in the color singlet state propagating through the target nucleus, one can expect considerable shadowing corrections to the nucleon structure function

$$xF_3 \propto \sigma_L - \sigma_R$$

extracted from nuclear data. Here $\sigma_L$ and $\sigma_R$ are the absorption cross sections for the left-handed and right-handed $W^+$-bosons, respectively. This structure function is known to be determined by the process (1) if $x$ is small enough. The latter is associated with the t-channel vacuum exchange or the sea-quark contribution to $xF_3$. However, the data available were
taken at moderately small-$x$ and the valence term in $xF_3$ cannot be neglected [7]. Free of the non-vacuum contributions is the combination of $\nu N$ and $\bar{\nu} N$ structure functions

$$\Delta x F_3 = x F_3^{\nu N} - x F_3^{\bar{\nu} N}. \quad (2)$$

Hereafter, $N$ stands for an iso-scalar nucleon. Our finding is that the shadowing correction to $\Delta x F_3$ extracted from $\nu Fe$ and $\bar{\nu} Fe$ data amounts to $20 - 25\%$ in the kinematical range of CCFR/NuTeV experiment [8, 9, 10].

Different approaches to nuclear shadowing in the neutrino DIS have been discussed previously (see [11, 12]). In this communication we develop the color dipole description of the phenomenon with particular emphasis on the left-right asymmetry effect specific for the charged current DIS.

## 2 Color dipole description of CC DIS off nuclei

The interaction of high-energy neutrino with the target nucleus can be treated as mediated by the interaction of the quark-antiquark dipole that the virtual $W^+$ transforms into. At small values of Bjorken $x$ this transition takes place at a large distance $l$, upstream the target:

$$l \sim \frac{1}{k_L} \sim \frac{1}{x m_N}. \quad (3)$$

Here $k_L$ is the longitudinal momentum transfer in the transition (1), $x = Q^2/2pq$, $p$ and $q$ are the target nucleon and the virtual $W^+$ four-momenta, respectively, $q = (\nu, 0, 0, \sqrt{\nu^2 + Q^2})$, $Q^2 = -q^2$ and $2pq = 2m_N\nu$.

Interaction of the color dipole of size $r$ with the target nucleon is described by the beam- and flavor-independent color dipole cross section $\sigma(x, r)$ [6, 13, 14]. At small $x$ the dipole size $r$ is a conserved quantum number. Therefore, the contribution of the excitation of open charm/strangeness to the nuclear absorption cross section for left-handed, $(\lambda = L = -1)$ and right-handed, $(\lambda = R = +1)$, $W^+$-boson of virtuality $Q^2$ is given by the color dipole factorization formula [15, 16]

$$\sigma^A_\lambda(x, Q^2) = \langle \Psi_\lambda | \sigma^A(x, r) | \Psi_\lambda \rangle$$

$$= \int dz d^2r \sum_{\lambda_1, \lambda_2} |\Psi^{\lambda_1, \lambda_2}_\lambda (z, r)|^2 \sigma^A(x, r), \quad (4)$$

\[ \]
where \[17\]
\[
\sigma^A(x, r) = 2 \int d^2b \left\{ 1 - \exp \left[ -\frac{1}{2} \sigma(x, r) T(b) \right] \right\}.
\]
(5)

Here \(T(b)\) is the the optical thickness of a nucleus,
\[
T(b) = \int_{-\infty}^{+\infty} dz \, n(z) (\sqrt{z^2 + b^2}),
\]
(6)
b is the impact parameter and \(n(r)\) is the nuclear matter density normalized as follows:
\[
\int d^3r n(r) = A.
\]
(7)

It is assumed that \(A \gg 1\). One can expand the exponential in Eq. (5) to separate the impulse approximation term and the shadowing correction, \(\delta\sigma^A\), in (4),
\[
\sigma^A = A\sigma - \delta\sigma^A.
\]
(8)

Here
\[
\sigma = \langle \Psi_\lambda | \sigma(x, r) | \Psi_\lambda \rangle
\]
\[
= \int dz d^2r \sum_{\lambda_1, \lambda_2} |\Psi_{\lambda_1, \lambda_2}^z(z, r)|^2 \sigma(x, r).
\]
(9)

To the lowest order in \(\sigma T\) the shadowing term reads
\[
\delta\sigma^A \approx \frac{\pi}{4} \langle \sigma^2 \rangle S^2_A(k_L) \int db^2 T^2(b),
\]
(10)

where
\[
\langle \sigma^2 \rangle = \langle \Psi_\lambda | \sigma(x, r)^2 | \Psi_\lambda \rangle
\]
\[
= \int dz d^2r \sum_{\lambda_1, \lambda_2} |\Psi_{\lambda_1, \lambda_2}^z(z, r)|^2 \sigma^2(x, r).
\]
(11)

The longitudinal nuclear form factor \(S_A(k_L)\) in Eq. (10) takes care about the coherency constraint,
\[
l \gg R_A.
\]
(12)

The approximation (10) represents the driving term of shadowing, the double-scattering term. It is reduced by the higher-order rescatterings by about 30% for iron and 50% for lead nuclei.
This accuracy is quite sufficient for order-of-magnitude estimates. The numerical calculations presented below are done for the full Glauber series (5),

$$\delta \sigma_\lambda^A = \pi T_A^2(k_L) \sum_{n=2}^{\infty} \frac{(-1)^n \langle \sigma_\lambda^n \rangle}{n! 2^{n-1}} \int db^2 T^n(b),$$

(13)

where the effect of finite coherence length is modeled by the factor $S_A^2(k_L)$ in rhs. A consistent description of the latter effect in electro-production was obtained in Ref. [18] based on the light-cone path integral technique of Ref. [19]. For related numerical studies of nuclear shadowing in electro-production see [20].

The LCWF, $\Psi_{\lambda_1,\lambda_2}(z, r)$, in Eqs. (4), (9) and (11) describes the Fock state $|c\bar{s}\rangle$ with the $c$-quark carrying the fraction $z$ of the $W^+$ light-cone momentum and the $\bar{s}$-quark with momentum fraction $1 - z$. The $c$- and $\bar{s}$-quark helicities are $\lambda_1 = \pm 1/2$ and $\lambda_2 = \pm 1/2$, respectively. This wave function derived in Ref. [7] is found in Appendix.

The diagonal elements of the density matrix

$$\rho_{\lambda\lambda'} = \sum_{\lambda_1,\lambda_2} \Psi_{\lambda_1,\lambda_2}^\lambda \left( \Psi_{\lambda_1,\lambda_2}^{\lambda'} \right)^*$$

(14)

for $\lambda = \lambda' = L, R$ entering Eqs. (4), (9) and (11) are as follows:

$$\rho_{RR}(z, r) = \left| \Psi_R^{1/2,1/2} \right|^2 + \left| \Psi_R^{-1/2,1/2} \right|^2$$

$$= \frac{8 \alpha_W N_c}{(2\pi)^2} (1 - z)^2 \left[ m^2 K_0^2(\varepsilon r) + \varepsilon^2 K_1^2(\varepsilon r) \right]$$

(15)

and

$$\rho_{LL}(z, r) = \left| \Psi_L^{-1/2,-1/2} \right|^2 + \left| \Psi_L^{-1/2,1/2} \right|^2$$

$$= \frac{8 \alpha_W N_c}{(2\pi)^2} z^2 \left[ \mu^2 K_0^2(\varepsilon r) + \varepsilon^2 K_1^2(\varepsilon r) \right],$$

(16)

where

$$\varepsilon^2 = z(1 - z)Q^2 + (1 - z)m^2 + z\mu^2$$

(17)

and $K_\nu(x)$ is the modified Bessel function. The $c$-quark and $\bar{s}$-quark masses are $m$ and $\mu$, respectively.

Consequences of the striking momentum partition asymmetry of both $\rho_{LL}$ and $\rho_{RR}$ were studied in Ref. [7]. This asymmetry was found to determine the vacuum exchange contribution.
to the structure function $xF_3$ of CC DIS:

$$2xF_3(x, Q^2) = \frac{Q^2}{4\pi^2\alpha_W} \left[ \sigma_L(x, Q^2) - \sigma_R(x, Q^2) \right].$$ (18)

### 3 Left and right: different scales - different dynamics

The CCFR/NuTeV structure functions $xF^\nu\nu_N$ and $xF^\bar{\nu}\nu_N$ are extracted from the $\nu Fe$ and $\bar{\nu} Fe$ data [8, 9, 10]. To estimate the strength of the nuclear shadowing effect in $xF_3$ at high $Q^2$ such that

$$\frac{m^2}{Q^2} \ll 1, \quad \frac{\mu^2}{Q^2} \ll 1$$ (19)

one should take into account that the dipole cross section $\sigma(x, r)$ in Eqs. (10) and (11) is related to the un-integrated gluon structure function $F(x, \kappa^2) = \frac{\partial G(x, \kappa^2)}{\partial \log \kappa^2}$, as follows [21]:

$$\sigma(x, r) = \frac{\pi^2}{N_c} r^2 \alpha_S(r^2) \int \frac{d\kappa^2 \kappa^2}{(\kappa^2 + \mu_G^2)^2} \frac{4[1 - J_0(\kappa r)]}{\kappa^2 r^2} F(x, \kappa^2).$$ (20)

In the Double Leading Logarithm Approximation (DLLA), i.e. for small dipoles, we have

$$\sigma(x, r) \approx \frac{\pi^2}{N_c} r^2 \alpha_S(r^2) G(x, C/r^2),$$ (21)

where $\mu_G = 1/R_c$ is the inverse correlation radius of perturbative gluons and $C \simeq 10$ comes from properties of the Bessel function $J_0(y)$. Because of the scaling violation $G(x, Q^2)$ rises with $Q^2$, but the product $\alpha_S(r^2)G(x, C/r^2)$ is approximately flat in $r^2$. Let us estimate first the contribution to $\langle \sigma^2 \rangle$ coming from the P-wave term, $\varepsilon^2 K_1(\varepsilon r)^2$, in Eqs. (15) and (16). The asymptotic behavior of the Bessel function, $K_1(x) \simeq \exp(-x)/\sqrt{2\pi x}$ makes the $r$-integration rapidly convergent at $\varepsilon r > 1$. Integration over $r$ in Eq. (11) yields

$$\langle \sigma^2_L \rangle \propto \int_0^1 dz \frac{z^2}{\varepsilon^4} \propto \frac{1}{Q^2 \mu^2}$$ (22)

and similarly

$$\langle \sigma^2_R \rangle \propto \int_0^1 dz \frac{(1 - z)^2}{\varepsilon^4} \propto \frac{1}{Q^2 \mu^2}.$$ (23)

We are not surprised to see that shadowing is the scaling, rather than the higher twist, $1/Q^2$, effect. Obviously, the integral (22) is dominated by $z \gtrsim 1 - \mu^2/Q^2$ i.e., by $\varepsilon^2 \sim \mu^2$ and,
consequently, by $r^2 \sim 1/\varepsilon^2 \sim 1/\mu^2$. A comparable contribution to the integral (22) comes from the S-wave term $\propto \mu^2 K_0(\varepsilon r)^2$ in $\rho_{LL}$. In Eq. (23) the integral is dominated by $z \lesssim m^2/Q^2$, corresponding to $\varepsilon^2 \sim m^2$. Therefore, $r^2 \sim 1/\varepsilon^2 \sim 1/m^2$. Thus, we conclude that the typical dipole sizes which dominate $\sigma_\lambda$ and $\langle \sigma_\lambda^2 \rangle$ are very different. In Ref. [7] basing on the color dipole approach we found the scaling cross sections $\sigma_L$ and $\sigma_R \propto 1/Q^2$ times the Leading-Log scaling violation factors $\propto \log Q^2/\mu^2$ and $\propto \log Q^2/m^2$, respectively. The scaling violations were found to be (logarithmically) dominated by

$$r^2 \sim 1/Q^2. \quad (24)$$

On the contrary, the contribution of small-size dipoles, $\sim 1/Q^2$, to $\langle \sigma_\lambda^2 \rangle$, defined in Eq. (11), proved to be negligible. At $\lambda = -1$ $\langle \sigma_\lambda^2 \rangle$ is dominated by large hadronic size $c\bar{s}$-dipoles, $r \sim 1/\mu$. Consequently,

$$\delta \sigma_L^A \propto 1/\mu^2. \quad (25)$$

At $\lambda = +1$ a typical $c\bar{s}$-dipole is rather small, $r \sim 1/m$, and $\delta \sigma_R^A$ is small as well:

$$\delta \sigma_R^A \propto 1/m^2. \quad (26)$$

Thus, there is a sort of filtering phenomenon, the target nucleus absorbs the $c\bar{s}$ Fock component of $W^+$ with $\lambda = -1$, but is nearly transparent for $c\bar{s}$ states with opposite helicity, $\lambda = +1$.

In a region of very small $Q^2$ a considerable asymmetry of the shadowing effect is expected also because at $Q^2 \to 0$

$$\langle \sigma_L^2 \rangle \propto \frac{1}{m^2 \mu^2} \gg \langle \sigma_R^2 \rangle \propto \frac{1}{m^4}. \quad (27)$$

However, at currently available $x \gtrsim 0.01$ both the mass threshold effect and nuclear form factor suppress $\delta \sigma_{L,R}^A$ at $Q^2 \lesssim (m + \mu)^2$.

4 Shadowing in $\Delta x F_3$: the magnitude of effect

It should be repeated that we focus on the vacuum exchange contribution to $xF_3$ corresponding to the excitation of the $c\bar{s}$-pair in the process (1). Therefore, the structure function $xF_3$
differs from zero only due to the strong left-right asymmetry of the light-cone Fock state \(|cs\rangle\). This contribution to \(xF_3\) can be reinterpreted in terms of the sea-quark densities of the target nucleon/nucleus. At currently available \(x \gtrsim 0.01\), in addition to the sea, the valence contribution must be taken into account. The valence term, \(xV\) is the same for both \(\nu N\) and \(\bar{\nu} N\) structure functions of an iso-scalar nucleon. The sea-quark term in \(xF_3^{\nu N}\) denoted by \(xS\) has opposite sign for \(xF_3^{\bar{\nu} N}\). Indeed, the substitution \(m \leftrightarrow \mu\) in Eqs. (15) and (16) entails \(\sigma_L \leftrightarrow \sigma_R\). Hence,

\[ xF_3^{\nu N} = xV + xS \]  
\[ \text{and} \]
\[ xF_3^{\bar{\nu} N} = xV - xS. \]

One can combine the \(\nu N\) and \(\bar{\nu} N\) structure functions to isolate the Pomeron exchange term. Indeed, from Eqs. (2), (28) and (29) it follows that

\[ \Delta xF_3 = 2xS. \]

The extraction of \(\Delta xF_3\) from CCFR \(\nu_\mu Fe\) and \(\bar{\nu}_\mu Fe\) differential cross section in a model-independent way has been reported in Ref. [10]. From Eq.(8) it follows that the shadowing correction to nucleonic \(\Delta xF_3\) extracted from nuclear data is

\[ \delta(\Delta xF_3) = \frac{Q^2}{4\pi^2\alpha_W} \frac{1}{A} \left( \delta\sigma^A_L - \delta\sigma^A_R \right). \]

Obviously, the shadowing correction to \(\Delta xF_3\) is related to \(\delta xF_3\),

\[ \delta xF_3 = \frac{1}{2} \delta(\Delta xF_3). \]

To give an idea of the magnitude of the shadowing effect we evaluate the ratio of the nuclear shadowing correction, \(\delta(\Delta xF_3)\), to the nuclear structure function of the impulse approximation, \(A\Delta xF_3\),

\[ R = \frac{\delta(\Delta xF_3)}{A\Delta xF_3} = \frac{\delta\sigma^A_L - \delta\sigma^A_R}{A\sigma_L - A\sigma_R}. \]

We calculate \(R\) as a function of \(Q^2\) for several values of Bjorken \(x\) in the kinematical range of CCFR/NuTeV experiment. Our results obtained for realistic nuclear densities of Ref.[22] are presented in Figure 1. Shown is the ratio \(R(Q^2)\) for different nuclear targets including \(^{56}\text{Fe}\).
Figure 1: The shadowing ratio $R$ as a function of $Q^2$ for several values of $x$ calculated from the nuclear charge densities of Ref.[22] for some sample nuclei.

We evaluate the ratio $R$ making use of the color dipole factorization as described above. The differential gluon density $\mathcal{F}(x_g, \kappa^2)$ in Eq. (20) was determined in Ref. [23]. It is a function of the gluon momentum fraction, $x_g$. For our purposes it suffices to use the approximation

$$x_g = 2x \left( 1 + \frac{M^2}{2Q^2} \right),$$

where $M^2 = (m + \mu)^2$. The constituent $u$, $d$, $s$- and $c$-quark masses we use are 0.2, 0.2, 0.35 and 1.3 GeV, respectively [7]. At $Q^2 \gg M^2$ Eq. (34) reduces to $x_g = 2x$ corresponding to the collinear DLLA. The longitudinal momentum transfer entering the nuclear form factor $S_A$ is

$$k_L = x_g m_N.$$  

One more simplification is that instead of the numerical Fourier transform of the realistic nuclear density, for $S_A$ we take a Gaussian parameterization normalized to the correct nuclear charge radius $R_A$.

At small $x$ and high $Q^2$ the shadowing correction scales, $\delta\sigma_{L,R} \propto 1/Q^2$. The absorption cross section $\sigma_{L,R}$ scales as well. The ratio $\delta\sigma_{L,R}/\sigma_{L,R}$ slowly decreases with growing $Q^2$ because of the logarithmic scaling violation in $\sigma_{L,R}$. Toward the region of $x > 0.01$, both the nuclear form factor and the mass threshold effect suppress $R$ at $Q^2 \lesssim M^2$ (see Fig.1).
5 Summary

In this paper we have presented the color dipole analysis of nuclear effects in charge current DIS. The emphasis was put on the pronounced effect of left-right asymmetry of shadowing in neutrino-nucleus DIS at small values of Bjorken $x$. The L-R asymmetry of $\nu N$ interactions is quantified in terms of the conventional structure function $xF_3$. The latter is dominated by the diffractive excitation of highly asymmetric $c\bar{s}$ component of the light-cone $W^+$. We predicted strikingly different scaling behavior of nuclear shadowing for the left-handed and right-handed $W^+$. Large, about $20 - 25\%$, shadowing in the Fe structure functions is predicted, which is important for a precise determination of the nucleon structure functions $xF_3$ and $\Delta xF_3$.

Appendix. The $W^+ \to c\bar{s}$-transition vertex is

$$gU_{cs}\bar{c}\gamma_{\mu}(1 - \gamma_5)s,$$

where $U_{cs}$ is an element of the CKM-matrix and the weak charge $g$ is related to the Fermi coupling constant $G_F$ through the equation

$$\frac{G_F}{\sqrt{2}} = \frac{g^2}{m_W^2}. \quad (36)$$

The polarization states of W-boson carrying the laboratory frame four-momentum

$$q = (\nu, 0, 0, \sqrt{\nu^2 + Q^2}) \quad (37)$$

are described by the four-vectors $e_{\lambda}$, with

$$e_0 = \frac{1}{Q}(\sqrt{\nu^2 + Q^2}, 0, 0, \nu),$$
$$e_{\pm} = \mp \frac{1}{\sqrt{2}}(0, 1, \pm i, 0), \quad (38)$$

the unit vectors $\vec{e}_x$ and $\vec{e}_y$ being in $q_x$- and $q_y$-direction, respectively. Then, the vector ($V$) and axial-vector ($A$) components of the light-cone wave function

$$\Psi_{\lambda_1,\lambda_2}^{\lambda_1,\lambda_2}(z, \mathbf{r}) = V_{\lambda_1,\lambda_2}(z, \mathbf{r}) - A_{\lambda_1,\lambda_2}(z, \mathbf{r}), \quad (39)$$

are as follows [7]:

$$\Psi_{\lambda_1,\lambda_2}^{\lambda_1,\lambda_2}(z, \mathbf{r}) = V_{\lambda_1,\lambda_2}(z, \mathbf{r}) - A_{\lambda_1,\lambda_2}(z, \mathbf{r}), \quad (39)$$
\[ V^{\lambda_1,\lambda_2}_{0}(z, r) = \frac{\sqrt{\alpha_W N_c}}{2\pi Q} \left\{ \delta_{\lambda_1,-\lambda_2} \left[ 2Q^2 z(1 - z) + (m - \mu)[(1 - z)m - z\mu] \right] K_0(\varepsilon r) - i\delta_{\lambda_1,\lambda_2}(2\lambda_1) e^{-i2\lambda_1\phi}(m - \mu)\varepsilon K_1(\varepsilon r) \right\} , \]  

(40)

\[ A^{\lambda_1,\lambda_2}_{0}(z, r) = \frac{\sqrt{\alpha_W N_c}}{2\pi Q} \left\{ \delta_{\lambda_1,-\lambda_2}(2\lambda_1) \left[ 2Q^2 z(1 - z) + (m + \mu)[(1 - z)m + z\mu] \right] K_0(\varepsilon r) + i\delta_{\lambda_1,\lambda_2} e^{-i2\lambda_1\phi}(m + \mu)\varepsilon K_1(\varepsilon r) \right\} . \]  

(41)

The Eqs. (40,41) describe scalar and axial quark-antiquark excitations of \( W^+ \). For the right- and left-handed \( W^+ \) corresponding to \( \lambda = \pm 1 \) we obtain

\[ V^{\lambda_1,\lambda_2}_{\lambda}(z, r) = -\frac{\sqrt{2\alpha_W N_c}}{2\pi Q} \left\{ \delta_{\lambda_1,-\lambda_2} \delta_{\lambda,2\lambda_1} [(1 - z)m + z\mu] K_0(\varepsilon r) - i(2\lambda_1)\delta_{\lambda_1,-\lambda_2} e^{i\lambda\phi} [(1 - z)\delta_{\lambda,-2\lambda_1} + z\delta_{\lambda,2\lambda_1}] \varepsilon K_1(\varepsilon r) \right\} \]  

(42)

and

\[ A^{\lambda_1,\lambda_2}_{\lambda}(z, r) = \frac{\sqrt{2\alpha_W N_c}}{2\pi Q} \left\{ \delta_{\lambda_1,-\lambda_2} \delta_{\lambda,2\lambda_1} (2\lambda_1) [(1 - z)m - z\mu] K_0(\varepsilon r) + i\delta_{\lambda_1,\lambda_2} e^{i\lambda\phi} [(1 - z)\delta_{\lambda,-2\lambda_1} + z\delta_{\lambda,2\lambda_1}] \varepsilon K_1(\varepsilon r) \right\} , \]  

(43)

where

\[ \varepsilon^2 = z(1 - z)Q^2 + (1 - z)m^2 + z\mu^2 \]  

(44)

and \( K_\nu(x) \) is the modified Bessel function. We consider only Cabibbo-favored transitions and

\[ \alpha_W = g^2/4\pi. \]

The quark and anti-quark masses are \( m \) and \( \mu \), respectively. The azimuthal angle of \( r \) is denoted by \( \phi \). To switch \( W^+ \rightarrow W^- \) one should perform the replacement \( m \leftrightarrow \mu \) in the equations above. The light-cone description of the neutral current (NC) interactions mediated by the \( Z \)-boson transition, \( Z \rightarrow q\bar{q} \), is an obvious extension of Eqs. (40,41,42,43), where one should equate quark masses, \( m = \mu \), and multiply vector and axial-vector components of \( \Psi^{\lambda_1,\lambda_2}_\lambda \) by corresponding NC coupling constants.
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