The influence of direct $D$-meson production to the determination on the nucleon strangeness asymmetry via dimuon events in neutrino experiments

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Abstract. Experimentally, the production of oppositely charged dimuon events by neutrino and anti-neutrino deep inelastic scattering (DIS) is used to determine the strangeness asymmetry inside a nucleon. Here we point out that the direct production of $D$-meson in DIS may make substantial influence to the measurement of nucleon strange distributions. The direct $D$-meson production is via the heavy quark recombination (HQR) and via the light quark fragmentation from perturbative QCD (LQF-P). To see the influence precisely, we compute the direct $D$-meson productions via HQR and LQF-P quantitatively and estimate their corrections to the analysis of the strangeness asymmetry. The results show that HQR has stronger effect than LQF-P does, and the former may influence the experimental determination of the nucleon strangeness asymmetry.

1 Introduction

Studying strange and anti-strange quark distributions of a nucleon is an important part in the study of the nucleon structure. An asymmetric strange distribution, i.e., the parton distribution function (PDF) of strange quark being not equal to that of anti-strange quark inside a nucleon, is naturally predicted by some non-perturbative models [1, 2, 9]. Further clear check of the strangeness asymmetry is not only important for the study of the nucleon structure itself, but also for understanding relevant phenomenon in some experiments. For example, the so-called NuTeV anomaly phenomenon [4, 5] can be explained by nonzero strangeness asymmetry $\kappa_0^{s,\bar{s}}=0.00196\pm0.00046(\text{stat})\pm0.00045(\text{syst})^{+0.00148}_{-0.00107}$ (external) [17], which is consistent with the global analysis, although at early stage the analysis of CCFR and NuTeV dimuon events at LO even NLO do not support the strangeness asymmetry $\kappa_0^{s,\bar{s}}=0.0002$.

Because of the smallness of the strange and anti-strange components in a nucleon, to measure the strangeness asymmetry is a challenging job indeed. The most sensitive reaction to measure the strange and anti-strange distributions is the production of dimuon events in neutrino and anti-neutrino nucleon deep inelastic scattering (DIS). To leading order (LO) of the dimuon production, the events are caused by the charged-current (CC) charm production subprocesses $\nu_\mu + s(d) \to \mu^- + c$ or $\bar{\nu}_\mu + \bar{s}(\bar{d}) \to \mu^+ + \bar{c}$ and a cascade decay $c \to \mu^+ + \cdots$ or $\bar{c} \to \mu^- + \cdots$. The relevant transition $\nu_\mu(\bar{\nu}_\mu) + d(\bar{d}) \to \mu^- (\mu^+) + c(\bar{c})$ is Cabibbo suppressed whereas $\nu_\mu(\bar{\nu}_\mu) + s(\bar{s}) \to \mu^- (\mu^+) + \bar{c}(\bar{s})$ is Cabibbo favored. Thus the dimuon events are sensitive to the strange and anti-strange distributions of the target nucleon. In the literature, measurements of the strangeness asymmetry via oppositely signed dimuon are carried out by CCFR and NuTeV [12, 13, 14, 15, 16, 17] experiments.

The results of global analysis [18, 6, 19] indicate the strangeness asymmetry $S^{-} = \int \xi \bar{\xi} s(\xi) - \bar{s}(\xi) d\xi$ likely to be positive, e.g., in Ref. [19] $-0.001 < S^{-} < 0.005$ is obtained. The recent NuTeV reanalysis up-to next-to-leading order (NLO) of perturbative QCD (pQCD) with improved method supports the positive strangeness asymmetry $S^{-} = 0.00196\pm0.00046(\text{stat})\pm0.00045(\text{syst})^{+0.00148}_{-0.00107}$ (external) [17], which is consistent with the global analysis, although at early stage the analysis of CCFR and NuTeV dimuon events at LO and even NLO do not support the strangeness asymmetry $\kappa_0^{s,\bar{s}}=0.0002$.

In this work, we take a systematic study on the influence of direct $D$-meson production at order $\alpha_s^2$ to the determination of the nucleon strangeness asymmetry via dimuon events in neutrino experiments. With consideration of the experimental kinematic cuts in CCFR and NuTeV, we point out that there are two kinds of direct $D$-meson production: heavy quark recombination process (HQR) and light quark fragmentation in the pQCD picture (LQF-P), which can contribute to the cross section difference between neutrino and anti-neutrino induced CC DIS. The direct $D$-meson production may influence the strangeness measurement (determination) that depends on their magnitude, so we further calculate the production quantitatively and investigate their influence to the strangeness determination. Although our preliminary result on HQR process has been briefly reported in Ref. [20],...
in this paper, we would present it in more detail and with some improvements. From final results we find that the influence of the direct \( D \)-meson production to the measurement of the strangeness asymmetry could not be negligible.

The rest of the paper is organized as follows. In section II, we discuss CC charm production to the dimuon cross sections at LO and NLO. And then we present the two kinds of direct \( D \)-meson production: HQR and LQF-P, which in fact are of high order ones, and show how the two processes can affect the extraction of the nucleon strangeness asymmetry. In section III, we show the numerical calculation about HQR direct production of strange distribution asymmetry. In section II, we discuss CC charm production to the dimuon and estimate the influence due to the direct production. In section IV, we present the cross sections at LO and NLO. And then we present the obtain results.

### 2 Dimuon events and direct \( D \)-meson production in neutrino DIS

Experimentally with the CCFR and NuTeV detectors, the oppositely charged dimuon signal induced by CC charm production in \( \nu_\mu \) (\( \overline{\nu}_\mu \)) DIS has a distinct feature and is not very difficult to be detected. The first muon of the dimuon is from the \( \nu_\mu(\overline{\nu}_\mu) \) vertex, and the second muon is from a little delayed muonic decay of the produced charm. The life time of \( \pi \) and \( K \) is much longer than charmed, so those muons from \( \pi \) or \( K \) meson decay can be largely eliminated, i.e., they will not contribute to the dimuon events concerned here.

According to pQCD factorization theorem, for \( \nu_\mu \)-proton DIS, to LO the differential cross section for dimuon production induced by CC production of charmed hadron \( H \) can be expressed as \cite{21}:

\[
\frac{d^2\sigma_{\nu_\mu p \rightarrow \mu^- \mu^+ X}}{d\xi dQ^2} \propto f_c \left[ d(\xi, Q^2) \left\{ V_{cd}^2 + s(\xi, Q^2) |V_{cs}|^2 \right\} + \sum_H D^H_c(z) Br_H \right].
\]

where \( d(\xi, Q^2) \) and \( s(\xi, Q^2) \) are the parton distribution functions (PDFs) of \( d \) and \( s \) quarks in the proton, and \( \xi \), relating to the Bjorken scaling variable \( x \) through \( \xi = x(1 + m^2_c/Q^2) \), is the light-cone momentum fraction of the struck quark; \( Q^2 = -(p - k)^2 \) is the minus squared invariant momentum transfer with \( p \) and \( k \) being the momentum of the incident \( \nu_\mu \) and the scattered \( \mu^- \) respectively. \( r_w \equiv 1 + Q^2/M_{W}^2 \) and \( f_c \equiv 1 - m^2_c/2S \) with \( M_W \) being the W-boson mass and \( S \) being the squared C.M. energy of the neutrino-proton system; \( D^H_c(z) \) is the fragmentation function for a charm quark to the charmed hadron \( H \), and \( Br_H \) is the branching ratio of muonic decay for \( H \). Carrying out the integration over \( z \) and the summation over \( H \), with the definition \( Br_{cf} \equiv \int dz \sum_H D^H_c(z) Br_H \), for target nucleus with proton number \( P \) and neutron number \( N \), the differential cross section can be expressed as

\[
\frac{d^2\sigma_{\nu_\mu A \rightarrow \mu^- \mu^+ X}}{d\xi dQ^2} = \frac{G_F^2}{\pi r_w^2} f_c \left[ \frac{P d(\xi, Q^2) + N u(\xi, Q^2)}{P + N} |V_{cd}|^2 + s(\xi, Q^2) |V_{cs}|^2 \right] Br_{cf}.
\]

where the PDFs in the neutron is related to PDFs of proton by \( d_n(\xi) = u(\xi) \), \( s_n(\xi) = s(\xi) \) etc.

Since \( |V_{cd}|^2 \sim 0.05 \) and \( |V_{cs}|^2 \sim 0.9 \) \cite{22}, the \( \nu_\mu \) induced dimuon cross section is sensitive to the strange distribution \( s(\xi, Q^2) \) in the target nucleus. Similarly, \( \overline{\nu}_\mu \) induced dimuon cross section is sensitive to anti-strange distribution \( \overline{s}(\xi, Q^2) \). The difference between dimuon cross sections induced by \( \nu_\mu \) and \( \overline{\nu}_\mu \) is directly related to the strange distribution asymmetry. The difference to LO can be expressed as:

\[
\frac{d^2\sigma_{\nu_\mu N \rightarrow \mu^- \mu^+ X}}{d\xi dQ^2} - \frac{d^2\sigma_{\overline{\nu}_\mu N \rightarrow \mu^- \mu^+ X}}{d\xi dQ^2} = \frac{G_F^2}{\pi r_w^2} f_c \left\{ [s(\xi, Q^2) - \overline{s}(\xi, Q^2)] |V_{cs}|^2 \right\} = \frac{1}{P + N} \left[ P d_v(\xi, Q^2) + N u_v(\xi, Q^2) \right] |V_{cd}|^2.
\]

\( d_v(\xi, Q^2) = d(\xi, Q^2) - \overline{s}(\xi, Q^2) \) and \( u_v(\xi, Q^2) = u(\xi, Q^2) - \overline{V}_{cs}(\xi, Q^2) \) are valence distributions of proton.

The Feynman diagrams for CC charm production at NLO from subprocess of \( \nu_\mu N \) DIS are shown in FIG. 1 \cite{16}. In fact, the last two diagrams of FIG. 1 \cite{16} are involved in the leading logarithm (LL) evolution for the parton flavor-singlet components of PDFs, and the flavor singlet components of the PDFs contribute to the dimuon cross sections symmetrically for the \( \nu_\mu \) and \( \overline{\nu}_\mu \)-induced DIS. Thus even up-to NLO, the difference between \( \nu_\mu \) and \( \overline{\nu}_\mu \)-induced dimuon cross section is still proportional to \( [s(\xi, Q^2) - \overline{s}(\xi, Q^2)] |V_{cs}|^2 + \frac{1}{P + N} \left[ P d_v(\xi, Q^2) + N u_v(\xi, Q^2) \right] |V_{cd}|^2 \), i.e., we need consider neither the last two diagrams of FIG. 1 nor the flavor singlet components of PDFs in the dimuon cross section difference.

However, the direct \( D \)-meson production that convolutes to the valence components in the nucleon, as we will discuss in the following, can raise the rate of dimuon production after experimental kinematic cuts, therefore determinations of strangeness asymmetry can be distorted by the direct \( D \)-meson production in certain degree.

Now let us focus the light on the contributions from the direct \( D \)-meson production, although according to pQCD the lowest order Feynman diagrams of the direct production are of the order \( a_s^2 \).

\(^1\) The \( D^* \) meson, the excited (1\(\overline{\eta}\)) bound state in \( ^3S_1 \), has very great cross section in production to compare with the \( D \)-meson production, and decays into \( D \) meson via strong and/or electromagnetic interaction with almost 100% branching ratio, thus the consequence of the \( D^* \)-meson production will be as direct \( D \)-meson production in accounting the muons in the dimuon events. Therefore throughout the paper “\( D \)-meson production” always mean the production of \( D \) and \( D^* \).
One of the direct $D$-meson production mechanisms, the so-called heavy quark recombination (HQR) process, at LO level is described by the diagrams in FIG. 2. It is stimulated from the heavy quark recombination mechanism \cite{23, 24, 25}, which combines a heavy quark and a light anti-quark of similar velocity to form a meson. Refs. \cite{24}, \cite{25} employ simple pQCD pictures and explain the charm photoproduction asymmetry and the leading particle effect \cite{26} successfully.

Namely the difference between $\nu_{\mu}^-$- induced and $\bar{\nu}_{\mu}$- induced dimuon cross sections caused by the HQR production of direct $D$-meson can be computed:

$$
\left[ \frac{d^2\sigma_{\nu_{\mu}^\pm \rightarrow \mu^\pm X}}{dxdQ^2} - \frac{d^2\sigma_{\bar{\nu}_{\mu}^\pm \rightarrow \mu^\pm X}}{dxdQ^2} \right]_{HQR} = \sum_{q,D} \int dx [q(x, \mu^2) - q(x, \mu^2)] \frac{d^2\sigma_{D}}{dxdQ^2} Br_{D},
$$

where $q$ denotes a possible light quark in the target and $D$ denotes a produced $D$-meson. $d\sigma_{D}$ denotes the differential cross section of the subprocess

$$
\nu_{\mu} + q \rightarrow \mu^+ + s \rightarrow d + D,
$$

which, in terms of $CP$ transformation, is known to be equal to that of the subprocess

$$
\bar{\nu}_{\mu} + q \rightarrow \mu^+ + s \rightarrow d + \bar{D}.
$$

And $Br_{D}$ denotes the muonic decay rate of the $D$-meson. In fact, from Eq. (5), it is easy to realize that only the $u - d$ valance components of the PDFs contribute to the cross section difference ($s - \bar{s}$ contribution comparatively negligible), thus $q$ could be either $u$ or $d$. Such a contribution may reduce the strange distribution asymmetry $S^{-}(\xi, Q^2) \equiv \xi[s(\xi, Q^2) - \bar{s}(\xi, Q^2)]$ value obtained by the analysis where the HQR process has been ignored. Thus, roughly speaking, the true value of the strange distribution inside a nucleon $S^{-}(\xi, Q^2)$ may be changed from the existent analysis result which has ignored the HQR effects by a positive correction factor $\delta S_{HR}^{HQR}(\xi, Q^2)$

$$
S_{real}^{-}(\xi, Q^2) = S_{analy}^{-}(\xi, Q^2) + \delta S_{HR}^{HQR}(\xi, Q^2).
$$

For a quantitative estimate of the HQR correction to the measured strangeness asymmetry, we suppose that the Cabibbo suppressed valence contribution in Eq. (4) has been deduced in LO (or NLO) analysis of dimuon events though present knowledge of the PDFs,

$$
f_{LO}|_{Vex}|^2 [s(\xi, Q^2) - \bar{s}(\xi, Q^2)]_{LO}^{\text{analy}} = \left[ \frac{d^2\sigma_{\nu_{\mu}}}{d\xi dQ^2} - \frac{d^2\sigma_{\bar{\nu}_{\mu}}}{d\xi dQ^2} \right]_{\text{ex}} = \frac{f_{LO}|_{Vex}|^2}{P + N} \frac{2P \bar{d}d}{\bar{N} u \bar{d} d}(\xi, Q^2),
$$

where $f_{LO} = G_f^2 f_{Vex}/\pi r^2$, is the coefficient in Eq. (4) for LO cross section, and $d^2\sigma/d\xi dQ^2$ is the differential cross section for $\nu_{\mu}(\bar{\nu}_{\mu})$-induced dimuon production. While in fact, the measured cross section also includes the contribution from HQR as Eq. (4), which should be deducted to obtain the real strange distribution asymmetry

$$
f_{LO}|_{Vex}|^2 [s(\xi, Q^2) - \bar{s}(\xi, Q^2)]_{LO}^{\text{real}} = \left[ \frac{d^2\sigma_{\nu_{\mu}}}{d\xi dQ^2} - \frac{d^2\sigma_{\bar{\nu}_{\mu}}}{d\xi dQ^2} \right]_{\text{ex}} = \left[ \frac{d^2\sigma_{\nu_{\mu}}}{d\xi dQ^2} - \frac{d^2\sigma_{\bar{\nu}_{\mu}}}{d\xi dQ^2} \right]_{\text{HQR}}.
$$
Fig. 3. Diagrams for the LQF-P process in $\nu_\mu$-induced CC DIS. The produced $D$ meson will decay partially into $\mu^+$ to form the second muon. Thick lines denote that of heavy quarks, and shaded blobs denote the $D$ meson produced directly.

$$-\frac{1}{P+N}f_{LO}|V_{cd}|^2[Pd_{c}(\xi,Q^2) + Nu_{c}(\xi,Q^2)]$$

$= f_{LO}|V_{cs}|^2[s(\xi,Q^2) - \bar{\pi}(\xi,Q^2)]_{LO}^{analy} - \left[\frac{d^2\sigma_{\nu_\mu}}{d\xi dQ^2} - \frac{d^2\sigma_{\nu_\mu}}{d\xi dQ^2}\right]_{HQR}$

Thus, the real strange distribution asymmetry $S^{-}_c(\xi,Q^2) \equiv \xi[s(\xi,Q^2) - \bar{\pi}(\xi,Q^2)]$ can be deduced.

$$S^{-}_{real,c}(\xi,Q^2) = S^{-}_{LO,c}(\xi,Q^2) - \frac{\xi}{f_{LO}|V_{cs}|^2} \left[\frac{d^2\sigma_{\nu_\mu}}{d\xi dQ^2} - \frac{d^2\sigma_{\nu_\mu}}{d\xi dQ^2}\right]_{HQR}$$

LQF-P production can be expressed as

$$\left[\frac{d^2\sigma_{\nu_\mu,N\rightarrow\mu^-\mu^+X}}{d\xi dQ^2}\right]_{\text{LQF-P}}$$

$$= \sum_D \int dx \frac{P\Phi(x) + N \Phi(x)}{P + N} \frac{d\sigma_{\nu_\mu}}{d\xi dQ^2} B_{RD}.$$ (12)

Where $d\sigma_{\nu_\mu}$ denotes the cross section of subprocess $\nu_\mu + \pi \rightarrow \mu^- + D + \pi$ (diagrams in Fig.3). The $\nu_\mu$-induced dimuon cross section from LQF-P production can be obtained by $CP$ transformation to that of $\nu$, and thus the $\nu_\mu$ and $\nu_\mu$-induced dimuon cross section difference from LQF-P process can be expressed as

$$\left[\frac{d^2\sigma_{\nu_\mu,N\rightarrow\mu^-\mu^+X}}{d\xi dQ^2} - \frac{d^2\sigma_{\nu_\mu,N\rightarrow\mu^-\mu^+X}}{d\xi dQ^2}\right]_{\text{LQF-P}}$$

$$= - \int dx \frac{P u_c(x) + N d_c(x)}{P + N} \sum_D \frac{d\sigma_{\nu_D}}{d\xi dQ^2} B_{RD}.$$ (13)

Note here that to compare with the valence components, the other components are tiny, so in Eq. (13), except the valence components, the other components are ignored.

Numerical calculation of the LQF-P process will be presented in section IV.

3 Calculation of the HQR process

As pointed out in the above section, the direct $D$ meson production from HQR process (diagrams in Fig. 2) may influence the measurement of the nucleon strangeness asymmetry. Our calculation of the HQR process follows the method in Refs. [23,24]. Namely we ‘factorize’ the production into the two steps, one of them (the one we start with) is the production of the relevant $1S_0$ or $3S_1$ state of $(c\bar{c})$, the other one (the following one) is the combined quark pair to evolve into the relevant $D$- or $D^+$-meson under certain possibility, that is determined by other experiments.

When a charm $c$ quark and a $\bar{q}$ light anti-quark with momentum $p_c$ and $p_\bar{q}$ respectively are produced with $p_\bar{q}$ to be small in the $c$-quark rest frame, the $c$ and $\bar{q}$ can be constructed into a $1S_0$ color-singlet state $(c\bar{c})$ if $q$- and $c$-quark have the same color, then we have the following substitution in amplitude for the combination of the partons:

$$v_j(u_q)\bar{u}_i(c) \rightarrow x_d \frac{\delta_{ij}}{N_c} m_c (p_c - m_c) \gamma_5,$$ (14)

where $p_q = x_q p_c$ and limit $x_q \rightarrow 0$ may be taken. Accordingly, for the color-singlet $3S_1$ state $(c\bar{c})$, $\gamma_5$-matrix in Eq. (13) should be replaced by $\gamma$, where $\epsilon^\mu(P,\lambda)$ is the polarization vector of the $3S_1$ $(c\bar{c})$ state with total momentum $P$ and polarization $\lambda$ of the quark pair. $\delta_{ij}$ in Eq. (13) is the color factor for color-singlet state with quark colors $i, j = 1, 2, 3$. Corresponding to the diagrams in Fig. 2(a)(b), the amplitude for the production of the
(c\bar{q}) color-singlet $^{1}S_{0}$ state induced by $\nu_{\mu}$ may be written down directly:

$$M = \frac{16\pi G_{F}\alpha_{s} m_{c} \delta_{mn}}{9\sqrt{2}m_{c}(2l \cdot p_{c})} V_{c\bar{q}}(l)\gamma_{\nu}(1 - \gamma_{5})u(p)$$

$$\times \overline{u}(l)\gamma^{\nu}\left[\gamma_{5}\frac{\not{p} - \not{k} - \not{k_{s}} + m_{c}}{(p - k - l)^{2} - m_{c}^{2}}\gamma_{\mu}(1 - \gamma_{5}) + \gamma_{\mu}(1 - \gamma_{5})\frac{\not{l} - \not{k_{s}}}{(l - k)^{2}}\gamma_{\nu}\right]v(k_{s}).$$

(15)

Here $p, k, l, k_{s}$ denote the momenta of the $\nu_{\mu}$, the initial incoming $\bar{q}$, the produced $\mu^{-}$ and the produced $\pi$ or $d$ respectively. $V_{c\bar{q}}$ is the CKM matrix element with $f = s, d$. Color factor $T_{mn}^{b} = \frac{4}{3}\delta_{mn}$ has been included in the amplitude of Eq.(15). Squaring the amplitude, averaging over the spin and color of the particles in the initial state, and summing up the spin, color and possible flavors ($f = s, d$) of the particles in the final state, the ‘averaged and summed’ squared amplitude is

$$|\overline{M}|^{2}[(c\bar{q})^{1}_{S_{0}}] = \frac{4\pi^{2}16^{3}G_{F}^{2}\alpha_{s}^{2}m_{c}^{2}}{81r_{c}^{2}B(s - A - B)^{3}}$$

$$\times \left\{C^{2}[(s - A)m_{c}^{2} - (s - A - B)Q]^{2}
+ m_{c}C[(s - A)m_{c}^{2} - (s - A - B)(s - Q^{2})]
+ m_{c}^{2}A[s(Q^{2} + m_{c}^{2}) - BQ^{2}]\right\}.$$ (16)

Here $s \equiv (p + l)^{2} = \sqrt{S}x$ is the squared C.M. energy of the subprocess. The variables $A, B, C$ are defined as $A = \frac{2}{3}(l \cdot k), B = 2(l \cdot k_{s}), C = 2(p \cdot k_{s})$.

For the color-singlet $^{3}S_{1}$ $(c\bar{q})$-production state, the ‘averaged and summed’ squared amplitude is

$$|\overline{M}|^{2}[(c\bar{q})^{3}_{S_{1}}] = \frac{4\pi^{2}16^{3}G_{F}^{2}\alpha_{s}^{2}m_{c}^{2}}{81r_{c}^{2}B(s - A - B)^{3}}$$

$$\times \left\{3C^{2}[(s - A)m_{c}^{2} + (s - A - B)Q]^{2}
+ C[(3Q^{2} - 5s + 2A)(s - A - B)m_{c}^{2}
+ 3(s - A)m_{c}^{2} + 4(B - S)(s - A - B)Q^{2}]
+ [s(Q^{2} + m_{c}^{2}) - BQ^{2}]\right\}.$$ (17)

For the color-octet $(c\bar{q})$ either in $^{1}S_{0}$ or $^{3}S_{1}$ state, the amplitude differs from the color-singlet ones only in the color factor, e.g., the $\delta_{mn}$ in Eq.(15) should be replaced by $\sqrt{\frac{3}{2}}\delta_{mn}$ with $a = 1, 2, \cdots, 8$.

Then the produced $(c\bar{q})$-states will evolve into either $D$ or $D^{*}$ meson with certain probabilities. The cross section for the subprocess of a $D$-meson production $\nu_{\mu} + \pi \rightarrow \mu^{-} + D + \pi(n\bar{d})$ can be expressed as

$$d\sigma_{D} = \sum_{c,s,q} d\hat{\sigma}(\hat{c}\bar{q}) \rho((\hat{c}\bar{q})^{c}_{s} \rightarrow D).$$

(18)

Where $c$ and $s$ denote the color and angular momentum quantum numbers of the $(c\bar{q})$ state, and $\rho((\hat{c}\bar{q})^{c}_{s} \rightarrow D)$ is the non-perturbative parameter characterizing the probability for the $(\hat{c}\bar{q})^{c}_{s}$ to evolve into a state including the $D$ meson.

For $(\hat{c}\bar{q})^{c}_{s}$ with various antiquark $\bar{q}$, spin $(s)$ and color $(c)$, there may be a number of $\rho$ parameters. However, the number can be greatly reduced in terms of symmetries and some approximations. First of all, according to the color-factor of the processes, the cross sections for the production via color-octet and color-singlet $(c\bar{q})$ differ by a single factor $\frac{1}{8}$, therefore we can express the cross section for $D$-meson production as follows:

$$d\hat{\sigma}_{D} = \sum_{c,s,q} d\hat{\sigma}((\hat{c}\bar{q})^{c}_{s} \rightarrow D) \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D).$$

(19)

with the definition:

$$\rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D) = \rho((\hat{c}\bar{q})^{c}_{s} \rightarrow D) + \frac{1}{8}\rho((\hat{c}\bar{q})^{c}_{s} \rightarrow D).$$

If the produced $D$-meson has different flavor from that of the quark pair $(c\bar{q})$, that means the quark pair $(c\bar{q})$ must emit a flavored object (such as a pion etc) in the meantime forming the $D$ meson, e.g., $(c\bar{q}) \rightarrow D^{+} + \pi^{-}$, then the relevant $\rho_{\text{eff}}$ will be relatively suppressed in the large $N_{c}$ limit. As in Refs. [24][25], we neglect such transitions from the quark pair $(c\bar{q})$ to a different flavored $D$-meson. Furthermore, SU(3) light quark flavor symmetry indicates $\rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{0}) \simeq \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*})$. As discussed in Ref. [24], heavy quark spin symmetry implies

$$\rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*}(c\bar{q})) = \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*}(c\bar{q})).$$

(20)

Thus, only two independent parameters are left:

$$\rho_{\text{sm}} \equiv \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{+}) = \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*+})$$

$$= \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{0}) = \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*0}),$$

(21)

$$\rho_{\text{sf}} \equiv \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*+}) = \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{+})$$

$$= \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*0}) = \rho_{\text{eff}}((\hat{c}\bar{q})^{c}_{s} \rightarrow D^{0}).$$

(22)

Now $\delta S_{\text{HR}}^{\pm}(\xi, Q^{2})$ caused by HQR, according to Eq.(11) can then be evaluated by means of Eqs.(16,17) with the auxiliary parameters $\rho_{\text{sm}}$ and $\rho_{\text{sf}}$ which are determined from relevant experiments. For the target to be proton, the HQR correction of Eq.(11) can be expressed as

$$\delta S_{\text{HR}}^{\pm}(\xi, Q^{2}) = \frac{\pi^{2}2\xi}{G_{F}^{2}F_{\pi}B_{\pi x}[V_{c\bar{q}}^{2}]
\times \int dx \left\{[d_{c}(x, \mu^{2})B_{D^{+}} + u_{c}(x, \mu^{2})B_{D^{0}}]
\times \left[\frac{d^{2}\hat{\sigma}(\hat{c}\bar{q})^{c}_{s} \rightarrow D^{+}}{d\xi dQ^{2}}\rho_{\text{sm}} + \frac{d^{2}\hat{\sigma}(\hat{c}\bar{q})^{c}_{s} \rightarrow D^{0}}{d\xi dQ^{2}}\rho_{\text{sf}}\right]
+ [d_{c}(x, \mu^{2})B_{D^{*+}} + u_{c}(x, \mu^{2})B_{D^{*0}}]
\times \left[\frac{d^{2}\hat{\sigma}(\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*+}}{d\xi dQ^{2}}\rho_{\text{sf}} + \frac{d^{2}\hat{\sigma}(\hat{c}\bar{q})^{c}_{s} \rightarrow D^{*0}}{d\xi dQ^{2}}\rho_{\text{sm}}\right]\right\}.$$ (23)
where the subprocess cross sections $d\hat{\sigma}(\varpi_j)_{s_0}$ and $d\hat{\sigma}(\varpi_j)_{s_1}$ are related to the ‘averaged and summed’ squared amplitudes of Eq. [10] and are independent of quark flavor $q$. And for the nucleus target with proton number $P$ and neutron number $N$, the HQR correction can be expressed as

$$\delta S_{HR}^{-}(\xi, Q^2) = \frac{\pi r_w^2 \xi}{G_{F}^{2} f_{j} B_{bf} |V_{cs}|^2} \int dx \left\{ \left[ \frac{P d_{v}(x, \mu^2) + N u_{v}(x, \mu^2)}{P + N} B_{fD}^{+} + \frac{P u_{v}(x, \mu^2) + N d_{v}(x, \mu^2)}{P + N} B_{fD}^{-} \right] \right. $$

$$\times \left( \frac{d^2\hat{\sigma}(\varpi_j)_{s_0}}{d\xi dQ^2} \rho_{sm} + \frac{d^2\hat{\sigma}(\varpi_j)_{s_1}}{d\xi dQ^2} \rho_{sf} \right) $$

$$\left. + \left[ \frac{P d_{v}(x, \mu^2) + N u_{v}(x, \mu^2)}{P + N} B_{fD}^{+} + \frac{P u_{v}(x, \mu^2) + N d_{v}(x, \mu^2)}{P + N} B_{fD}^{-} \right] \right. $$

$$\times \left( \frac{d^2\hat{\sigma}(\varpi_j)_{s_0}}{d\xi dQ^2} \rho_{sf} + \frac{d^2\hat{\sigma}(\varpi_j)_{s_1}}{d\xi dQ^2} \rho_{sm} \right) \right\} (24)$$

For a quantitative estimate, in the following, we calculate the $\delta S_{HR}^{-}(\xi, Q^2)$ for isoscalar target ($P = N$), with the parameters $\rho_{sm} = 0.15$, $\rho_{sf} = 0$, which are taken from the extraction from experimental charm photoproduction asymmetry by Ref. [22]. In this case, the HQR correction of Eq. (24) can be simplified.

$$\delta S_{HR}^{-}(\xi, Q^2) = \frac{\pi r_w^2 \xi \rho_{sm}}{G_{F}^{2} f_{j} B_{bf} |V_{cs}|^2} \int dx \frac{d_{v}(x, \mu^2) + u_{v}(x, \mu^2)}{2} $$

$$\times \left[ (Br_{fD}^{+} + Br_{fD}^{-}) \frac{d^2\hat{\sigma}(\varpi_j)_{s_0}}{d\xi dQ^2} + (Br_{fD}^{++} + Br_{fD}^{--}) \frac{d^2\hat{\sigma}(\varpi_j)_{s_1}}{d\xi dQ^2} \right] . (25)$$

In fact, for calculating Eq. (25) we need to carry out a three-dimensional integration numerically. In the C.M. frame of the subprocess, $\xi$ and $Q^2$ are related to the energy $k_0$ and the angle $\theta_1$ of the outgoing muon (from $\nu_\mu$ or $\bar{\nu}_\mu$ vertex) relative to the incident direction:

$$\xi = \frac{x_\perp^2 k_0 (1 - \cos \theta_1) + m_{c}^2}{s - x_\perp k_0 (1 + \cos \theta_1)} (26)$$

$$Q^2 = \sqrt{s k_0 (1 - \cos \theta_1)} , (27)$$

Bearing the NuTeV dimuon experiment in mind, the incident energy of neutrino or anti-neutrino is taken to be $E_\nu = 160$ GeV, which is approximately averaged value of the experiment. Furthermore, for our numerical calculation, charm mass $m_{c}$ is fixed with the value 1.5 GeV, the coupling constant $\alpha_{S}(\mu)$ is running as specified in CTEQ6L [28] and the parton distribution functions of the nucleon are taken from CTEQ6L [28] too. The branching ratio for the muonic decay of $D$ meson and the CKM matrix elements are taken to be the central values in Ref. [22], and the $B_{jf}$ is taken to be the central value of $B_{c \rightarrow \mu^{+} + X}$ in Ref. [22]. Since the two opposite charged muons in NuTeV experiment are required to have energy greater than 5 GeV, so we make a cut for the produced $\mu$ and the $D$ meson accordingly.

The obtained result of $\delta S_{HR}^{-}$ as a function of $\xi$ for $Q^2 = 10$ GeV$^2$ (solid lines), $Q^2 = 20$ GeV$^2$ (dash-dotted lines), and $Q^2 = 30$ GeV$^2$ (dashed lines) is shown in FIG.4. Since the calculation is of LO for the direct $D$-meson production, there are theoretical uncertainties, such as that from the energy scale $\mu$ of perturbative QCD, so to see the uncertainty, we calculate $\delta S_{HR}^{-}$ with two types of choices about $\mu$. The thick lines in the figure present the results where the factorization scale $\mu$ is taken to be $\mu = \sqrt{p_{\perp}^2 + m_{c}^2} \equiv m_{c}$, with $p_{\perp}$ being the transverse momentum of the produced $D$ meson to the direction of the $W$ boson in the nucleon rest frame. It is an analogous choice to that in Ref. [23], where its relevant charm photoproduction is calculated under the factorization scale $\mu = \sqrt{p_{\perp}^2 + m_{c}^2}$ with $p_{\perp}$ being the transverse momentum of the produced $D$ meson to the incident photon direction. The results in an alternative scale choice $\mu = \sqrt{Q^2} = Q$, are also shown by the thin lines in FIG.4. One fact that should be noted here is that the cross section from the HQR process decreases very slowly with the increase of the energy cut of the produced $D$ meson $E_{cut}$, namely, the ‘recombination’ is not suppressed by the cut taken in experiments very much. That can be understood by the fact that the difference of the direct production of $D$-meson by the HQR is related to the subprocess with valence quark inside a nucleon (Eq. (11)) so that the $D$-meson relevant to the difference can carry comparatively high energy (momentum) that escapes from the cut quite a lot.

From FIG.4 one can see that at fixed $Q^2$ in each case, $\delta S_{HR}^{-}(\xi)$ peaks in the region $\xi = 0.1 - 0.2$, over the peak

![Fig. 4. $\delta S_{HR}^{-}$ as a function of $\xi$ for $Q^2 = 10$ GeV$^2$ (solid lines), $Q^2 = 20$ GeV$^2$ (dash-dotted lines), and $Q^2 = 30$ GeV$^2$ (dashed lines). Thick lines are results for $\mu = m_{c} \equiv \sqrt{p_{\perp}^2 + m_{c}^2}$ and thin lines are for $\mu = Q$.](image-url)
$\delta S_{HR}$ as a function of $Q^2$ for $\xi = 0.06$ (solid lines), $\xi = 0.15$ (dash-dotted lines), and $\xi = 0.3$ (dashed lines). Thick lines are results for $\mu = \mu_0 = \sqrt{p_{T\perp} + m_s^2}$ and thin lines are for $\mu = Q$.

$\delta S_{HR}$ decreases with $Q^2$ increases, and the results with factorization scale $\mu = Q$ are smaller than those with $\mu = \mu_0$. The uncertainty from the choice of the factorization scale $\mu$ can also be seen when the scale $\mu = \mu_0$ is varied by a factor of 2: the results become nearly trebles when $\mu = \mu_0/2$, and the results reduce nearly by half when $\mu = 2\mu_0$. Generally the uncertainty may be suppressed by NLO calculation, but we leave the study beyond the present calculations.

The behavior of $\delta S_{HR}$ as functions of $Q^2$ is shown in FIG.5: the solid lines are those for $\xi = 0.06$, the dash-dotted lines are those for $\xi = 0.15$ and dashed lines are those for $\xi = 0.3$. The thick lines are results for factorization scale $\mu = \mu_0$ and the thin lines are those for $\mu = Q$.

In our calculation, a collinear singularity may arise from the strange quark propagator in diagrams FIG.2(b) and FIG.2(d) when the strange quark mass $m_s$ is set to be zero. In the limit $x_q \to 0$, the denominator of the propagator is $2k_s \cdot l$, which can reach zero when $k_s$ and $l$ are in the same direction. To avoid this singularity, we have taken the strange quark mass $m_s$ equal to its current mass 90 MeV in our above numerical calculations.

There are some uncertainties in the treatment of this singularity. To see this, in FIG.6, we show the results of $\delta S_{HR}$ as a function of $m_s$ for $Q^2 = 20$ GeV$^2$ at two $\xi$ values for two choices of $\mu$. Generally, the results decrease with the increase of $m_s$, and in the favored range 70-120 MeV for $m_s$, the results may vary within 15%. Lower $m_s$ range shows greater $m_s$ dependence. There is another way to avoid the singularity, i.e., to keep the $x_q$ suppressed terms in the s quark propagator and take $x_q$ to be reasonable finite instead of zero. When taking $x_q = 1/6$ (approx light constituent mass over D meson mass), the numerical results reduce nearly by half.

As indicated by the results above, when measuring the strange distribution asymmetry inside a nucleon, one should consider the correction $\delta S_{HR}(\xi, Q^2)$ caused by HQR, which is comparable to the existent measured value. For example, at $Q^2 = 20$ GeV$^2$ (about averaged value in NuTeV experiment) and for $\mu = \mu_0$, the HQR correction to strangeness asymmetry by integrating $\delta S_{HR}(\xi)$ over $\xi$ can be 0.002 approximately. To be comparison, the recent NLO analysis of the NuTeV dimuon data [17] and the global analysis [19] present the central value of the strangeness asymmetry $S^- \approx 0.002$. Thus, the HQR could not be negligible in the extraction of the strangeness asymmetry.

The HQR correction may enhance the strangeness asymmetry by a larger positive value, and large positive strangeness asymmetry can help to explain the NuTeV anomaly [8][9][10][11].

The value of the parameter $\rho_{sm}$ still has some uncertainty as discussed in Refs.[24][25], whereas, the magnitude order of $\rho_{sm}$ can not be changed and the influence from HQR to the measurement of the nucleon strangeness asymmetry could not be negligible. Moreover, more accurate $\rho$ parameters are needed not only for a better understanding of the HQR effect but also for better determination of the strangeness asymmetry inside a nucleon.

4 Calculation of the LQF-P process

The diagrams for the LQF-P process in $\nu_\mu$-induced CC DIS are shown in FIG.3. The subprocess can contribute to the $\nu_\mu$- and $\bar{\nu}_\mu$-induced dimuon cross section difference as shown in Eq.(13), and thus may also influence the measurement of the strangeness asymmetry.
The calculation of the cross section $d\hat{\sigma}_D$ for LQF-P process can be factorized into the convolution of the LO subprocess cross section for light quark $\bar{q}$ production and the fragmentation function $D_D^D(z)$ of the light quark $\bar{q}$ into a $D$ meson:

$$d\hat{\sigma}_D = \sum_q \int_0^1 dz \, d\hat{\sigma}_{q\pi^+ \to \bar{q} \gamma} D_D^D(z),$$

(28)

where $\bar{q}$ could be $\bar{d}$ or $\bar{u}$, with $D$ being $D^+$ or $D_s^+$ or $D_{s}^{(*)}$. With $CP$ transformation and SU(3) flavor symmetry, only two independent light-quark fragmentation functions remain:

$$D_q(z) \equiv D_{D^+}^D(z) = D_{D^{(*)}}^D(z),$$

$$D_s^*(z) \equiv D_{D_{s}^{(*)}}^D(z) = D_{D_{s}^{(*)}}^{D^{(*)}}(z).$$

(29)

The fragmentation functions are calculable with pQCD [30][31][32][33]. For $\bar{q} \rightarrow D_{D^+}(\bar{q} \gamma)$, the fragmentation function can be expressed as [32]

$$D_D^D(z) = \frac{1}{16\pi^2} \times \int ds \, \theta(s - \frac{(m_q + m_c)^2}{z} - \frac{m_c^2}{1 - z}) \lim_{q' \to \infty} \frac{|M|^2}{|M_0|^2}$$

(30)

where $M$ and $M_0$ are amplitudes for the $D$ production and the LO on shell $\bar{q}\gamma$ production respectively; Let $P_D = p_q + p_c$ and $k_c$ denote the momenta of the produced $D$ meson and the $\bar{q}$ quark respectively. Then $q = P_D + k_c$ is their total momentum, and $s = q^2$. The variable $z$ is defined by $z = (P_0^2 + P_D^2)/(q^2 + q'^2)$ in a frame, where $q = (q^1, 0, 0, q^4)$. In axial gauge, the amplitude corresponding to the diagram Fig.3(b) is suppressed and can be neglected [33]. Thus only diagram Fig.3(a) contributes. The form of the bound state can be described by the $B$-S wave functions, for the production of $^1S_0$ state color-singlet $D$ meson,

$$\chi_B(1S_0) = \frac{R(0)}{3\sqrt{3}\pi M_D} \gamma_5 (P_D + M_D),$$

(31)

and for the production of $^3S_1$ state color-singlet $D$ meson,

$$\chi_B(3S_1) = \frac{R(0)}{3\sqrt{3}\pi M_D} \frac{\epsilon \cdot P}{P_D + M_D},$$

(32)

which should appear in the amplitude $M$ of Eq. (30). Here $R(0)$ is the non-relativistic radial wave function at the origin for the $D$ meson, and $\epsilon$ is the polarization vector of the $^3S_1$ state $D$ meson.

The fragmentation function for light quark into $^1S_0$ state $D$ meson is given by [32]

$$D_q(z) = \frac{2\alpha_s(2m_c)^2 |R(0)|^2}{81\pi^2 m_c^2} \frac{rz(1 - z)^2}{|1 - (1 - r)z|^6} \times [6 - 18(1 - 2r)z + (21 - 74r + 68r^2)z^2 - 2(1 - r)(6 - 19r + 18r^2)z^3 + 3(1 - r)^2(1 - 2r + 2r^2)z^4],$$

(33)

where $r = m_c/M_D$. And the fragmentation function for light quark into $^3S_1$ state $D$ meson is given by [32]

$$D_s^*(z) = \frac{2\alpha_s(2m_c)^2 |R(0)|^2}{27\pi m_c^3} \frac{rz(1 - z)^2}{|1 - (1 - r)z|^6} \times [2 - 2(3 - 2r)z + 3(3 - 2r + 4r^2)z^2 - 2(1 - r)(4 - r + 2r^2)z^3 + (1 - r)^2(3 - 2r + 2r^2)z^4].$$

(34)

The value of $R(0)$ can be estimated from the pseudoscalar meson decay constant $f_\pi$ through the relation

$$R(0) = \sqrt{\frac{\pi M_D}{3}} f_\pi.$$  

(35)

We take the central values from Ref. [22] for $M_{D^+}$ and $f_{D^+}$ in calculating $R(0)$, and obtain $R(0) = 0.31$ GeV$^{-2}$. We take one-loop $\alpha_s$ with $\Lambda = 326$ MeV for 4 flavors as in CTEQ6L, and obtain $\alpha_s(2m_c) = 0.255$ for $m_c = 1.5$ GeV. The $r$ value in Eqs. (33,34) is evaluated by taking $M_D = 1.87$ GeV and $M_{D^+} = 2.01$ GeV [22].

The fragmentation functions $D_q(z)$ and $D_s^*(z)$ are shown in Fig.7. Integrating $D_q(z)$ and $D_s^*(z)$ over $z$, then $D_q \equiv \int D_q(z)dz = 2.01 \times 10^{-5}$ and $D_s^* \equiv \int D_s^*(z)dz = 1.77 \times 10^{-5}$ are obtained.

With the formulas for factorization and the fragmentation functions Eqs. (33,34), for isoscalar target, the ‘correction’ $\delta S^{LQF-P}_{\pi\rightarrow D_{D^+}}$ from LQF-P to measured strangeness asymmetry can be computed:

$$\delta S^{LQF-P}_{\pi\rightarrow D_{D^+}} \approx \int dz dz' \frac{\pi r^2 x}{G_F |V_{ct}|^2 f_c B_{ct}} \frac{u_c(x) + d_c(x)}{2}$$

$$\times \left\{ \frac{d\hat{\sigma}_{q\pi^+ \to \bar{q} \gamma}}{dQ^2} [D_q(z) Br_{D^+} + D_s^*(z) Br_{D_{s}^{(*)}}] \right\}$$

$$\approx \frac{1}{2 |V_{ct}|^2 B_{ct}} \{ Y(D_q Br_{D^+} + D_s^* Br_{D_{s}^{(*)}}) \}. \ (36)$$
Here \( Y = \int [u(x) + d(x)](1 - y^2)dx \) with \( y = Q^2 / xS < 1 \). \( Y = 0.16 \) is evaluated from the CTEQ6L parton distributions at \( Q^2 = 20 \text{ GeV}^2 \) and \( E_p = 160 \text{ GeV} \), and the muonic decay rates are taken to be the central value from Ref. [22]. With the parameters given above, the final result of \( \delta S_{LQF-P}^\text{−} \) is obtained:

\[
\delta S_{LQF-P}^\text{−} \approx 0.53 \times 10^{-5}.
\] (37)

Such a ‘correction’ to the strangeness asymmetry from the LQF-P process is much smaller than that measured. Thus, LQF-P gives little influence in the extraction of the nucleon strangeness asymmetry.

We should note here that as pointed out in Section II, LQF-P in neutrino and anti-neutrino DIS may generate not only the oppositely charged dimuon events, but also the same charged dimuon events and trimuon events (not from direct \( D \)-meson production), thus the fact that in neutrino and anti-neutrino DIS experiments either the same charged dimuon events or trimuon events are very rare is consistent with the small value of \( \delta S_{LQF-P}^\text{−} \) as shown in Eq. (37).

5 conclusions

The measurement of the nucleon strangeness asymmetry is important for the study of nucleon structure and certain related phenomenon. The cross section difference between the dimuon production from neutrino and anti-neutrino DIS is sensitive observable to the strange distribution asymmetry. Whereas in this work, we point out two types of direct charmed meson production, i.e., HQR and LQF-P at order \( \alpha_s^2 \), which also contribute to oppositely charged dimuon production. These processes are not included in the experimental analysis, therefore we further study their influence to the measurements of the nucleon strangeness asymmetry. With quantitative calculations in terms of pQCD, we find that HQR affects the extraction of the strange distribution asymmetry with a positive ‘correction’ \( \delta S_{HR}^\text{−}(\xi, Q^2) \), and the ‘correction’ can be so large as \( \delta S_{HR}^\text{−} \sim 10^{-3} \). For the other one, LQF-P provides a small ‘correction’ to the measurement, that is of the order \( \lesssim 10^{-5} \), so that it can be neglected. The influence of HQR to the measurement of the nucleon strangeness asymmetry from neutrino and anti-neutrino DIS can not be negligible, that may provide a positive correction to the present value of the strangeness asymmetry, and is also helpful to explain the NuTeV anomaly. We think that a reanalysis of the strange distribution asymmetry with consideration of the direct-\( D \) production from HQR process is needed to improve the value of the nucleon strangeness asymmetry.

\( D \)-meson directly produced in neutrino and anti-neutrino deep inelastic scattering (DIS) contributes to oppositely charged dimuon events, which has not been considered in experimental analysis so far. Hence we conclude that in determining the strangeness asymmetry in a nucleon via measuring the dimuon production in neutrino and anti-neutrino DIS, the contribution from HQR production of the direct \( D \)-meson should be deducted precisely.

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