Diffractive heavy pseudoscalar-meson productions by weak neutral currents

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Abstract. A first theoretical study for neutrino-induced diffractive productions of heavy pseudoscalar-mesons, $\eta_c$ and $\eta_b$, off a nucleon is performed based on factorization formalism in QCD. We evaluate the forward diffractive production cross section in perturbative QCD in terms of the light-cone wave functions of Z boson and $\eta_{c,b}$ mesons, and the gluon distribution of the nucleon. The diffractive production of $\eta_c$ is governed by the axial vector coupling of the longitudinally polarized Z boson to $Q\bar{Q}$ pair, and the resulting $\eta_c$ production cross section is larger than the $J/\psi$ one by one order of magnitude. The bottomonium $\eta_b$ production, which shows up for higher beam energy, is also discussed.

Exclusive diffractive lepton productions of neutral vector mesons provide unique insight into an interplay between nonperturbative and perturbative effects in QCD. The diffractive processes are mediated by the exchange of a Pomeron with the vacuum quantum numbers, whose QCD description is directly related to the gluon distributions inside the nucleons for small Bjorken-$x$ [1, 2, 3]. The processes also allow us to probe the light-cone wave functions (WFs) of the vector mesons. Relating to the latter point, however, the applicability is apparently limited to probing the neutral vector mesons due to the vector nature of the electromagnetic current.

In this talk, we propose the exclusive diffractive productions of mesons in terms of the neutrino beam. The weak currents allow us to observe both neutral and charged mesons by Z and W boson exchanges, and these mesons can be not only vector but also other types of mesons including pseudoscalar mesons. Thus, such processes may reveal structure of various kinds of mesons, the coupling of the QCD Pomeron to quark-antiquark pair with various spin-flavor quantum numbers, and also information on the CKM matrix elements. There already exist some experimental data for $\pi$, $\rho$, $D_s^\pm$, $D_s^+$ [4, 5], $D_s^{++}$ [6], and $J/\psi$ production [4], although the amount of the data is not enough. On the other hand, there are only a few theoretical calculations, e.g., for the $J/\psi$ production in a vector meson dominance model [7] and for $D_s^-$ production with the generalized parton density [8]. Here, our interest will be directed to diffractive productions of heavy pseudoscalar mesons, $\eta_c$ and $\eta_b$. So far $\eta_c$ has been observed via the decays of $J/\psi$ or $B$ mesons produced by $p\bar{p}$ and $e^+e^-$ reactions, while $\eta_b$ has not been observed. The diffractive productions via the weak neutral current will give a direct access to $\eta_c$ as well as a new experimental method to identify $\eta_b$ by e.g. measuring the two photon decay.
We treat the $\eta_c$ and $\eta_b$ productions by generalizing the approach in the leading logarithmic order of perturbative QCD, which has been developed successfully for the vector meson electroproductions [1, 2, 3]. We consider the near-forward diffractive productions $Z^*(q) + N(P) \rightarrow \eta_Q(q+\Delta) + N'(P-\Delta)$, where $Q = c, b$, and each momentum is labeled in Fig. 1. Here the total center-of-mass energy $W = \sqrt{(P+q)^2}$ is much larger than any other mass scales involved, i.e., $W^2 \gg Q^2(= -q^2), -t(-\Delta^2), m_Q^2, \Lambda_{QCD}^2$, etc., with $m_Q$ the heavy-quark mass. $-t \ll m_Q^2$ and $m_Q^2 \gg \Lambda_{QCD}^2$ are also assumed. The crucial point is that at such high $W$ the scattering of the $Q\bar{Q}$ pair on the nucleon occurs over a much shorter timescale than the $Z^* \rightarrow Q\bar{Q}$ fluctuation or the $\eta_Q$ formation times (see Fig. 1). As a result, the production amplitudes obey factorization in terms of the $Z$ and $\eta_Q$ light-cone WFs. Also, a hard scale $m_Q$ ensures the application of perturbative QCD for the $Z$ WFs even for $Q^2 = 0$. The $Q\bar{Q}$-N elastic scattering amplitude, sandwiched between the $Z$ and $\eta_Q$ WFs, is further factorized into the $Q\bar{Q}$-gluon hard scattering amplitude and the nucleon matrix element corresponding to the (unintegrated) gluon density distribution [1, 2, 3]. The participation of the “new players” $Z$ and $\eta_Q$ requires an extension of the previous works [1, 2, 3] by introducing the corresponding light-cone WFs.

\[
\begin{align*}
\text{FIGURE 1.} & \text{ A typical diagram for the exclusive diffractive } \eta_Q \ (Q = c, b) \text{ productions induced by neutrino (v) through the Z boson exchange. There are other diagrams by interchanging the vertices on the heavy-quark lines.} \\
\end{align*}
\]

First of all, we discuss the extension due to the participation of the $Z$ boson. The $ZQ\bar{Q}$ weak vertex of Fig. 1 is given by $(g_W/2\cos\theta_W)\gamma^\mu(c_V - c_A\gamma_5)$, where $(g_W/2\cos\theta_W)^2 = \sqrt{2}G_F M_Z^2$ with $G_F$ the Fermi constant and $M_Z$ the $Z$ mass. $c_V = 1/2 - (4/3)\sin^2\theta_W, c_A = 1/2$ for the $c$-quark and similarly for the $b$-quark. As usual, we introduce the two light-like vectors $q'$ and $p'$ by the relations $q = q' - (Q^2/s)p', P = p' + (M_N^2/s)q', s = 2q' \cdot p' + M_N$ the nucleon mass, and the Sudakov decomposition of all momenta, e.g., $k = \alpha q' + \beta p' + k_\perp (k_\perp^2 = -q_\perp^2)$. We also introduce the polarization vectors $\epsilon^{(\xi)}_\mu (\xi = 0, \pm 1)$ of the virtual $Z$ boson to satisfy $\sum_{\xi}(-1)^{\xi+1}\epsilon^{(\xi)}_\mu \epsilon^{(\xi)}_\nu = -g_{\mu\nu} + q_\mu q_\nu/M_Z^2$, which is the numerator of the propagator for the massive vector boson. Because the $q_\mu q_\nu/M_Z^2$ term vanishes when contracted with the neutral current, we conveniently choose as $\epsilon^{(0)} = q'/Q + p'/Q/s$ and $\epsilon^{(\pm 1)} = \epsilon^{(\pm 1)} = (0, 1, \pm i, 0)/\sqrt{2}$ for the longitudinal ($\xi = 0$) and transverse ($\xi = \pm 1$) polarizations respectively. The light-cone WFs for the virtual $Z$ boson can be obtained in analogy with the photon light-cone WF used in the $J/\psi$ electroproductions [1, 2] as

\[
\Psi_{\lambda'\lambda}^{Z(\xi)}(\alpha, k_\perp) = -\frac{g_W}{2\cos\theta_W} \frac{\sqrt{\alpha(1-\alpha)} \bar{u}_\lambda(k) \epsilon^{(\xi)}_\mu (c_V - c_A\gamma_5) v_{\lambda'}(q-k)}{\alpha(1-\alpha) Q^2 + k_\perp^2 + m_Q^2},
\]

(1)

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where \( u_\lambda(k) \) \((\nu_\lambda'(q - k))\) denote the on-shell spinor for the (anti)quark with helicity \( \lambda^{(i)} \).

Next we proceed to the light-cone WFs for the \( \eta_Q \) meson. It is convenient to exploit the correspondence with the case of the \( J/\psi \) electroproductions. The light-cone WFs for the heavy vector mesons \( V = J/\psi, \bar{Y} \) have been discussed in many works, but are still controversial in the treatment of subleading effects like the Fermi motion corrections \([2, 3, 9]\), corrections to ensure the pure \( S \)-wave \( Q\bar{Q} \) state \([10, 11]\), etc. Here we employ the vector-meson light-cone WFs given by

\[
\Psi_{\lambda\lambda'}^{V(\bar{\sigma})^*}(\alpha, k_\perp) = -\frac{\nu_\lambda'(q + \Delta - k)}{\sqrt{1 - \alpha}} \gamma_5 \mathcal{R} u_\lambda(k) \frac{\phi^*(\alpha, \vec{k}_\perp)}{M_V},
\]

(2)

where \( M_V \) and \( \epsilon^{(\bar{\sigma})(\sigma)} = (\bar{\sigma} = 0, \pm 1) \) are the mass and the polarization vector of the vector meson with \((q + \Delta)^2 = M_V^2, \quad \epsilon^{(\bar{\sigma})} \cdot (q + \Delta) = 0, \quad \epsilon^{(\bar{\sigma})} \cdot \epsilon^{(\sigma)} = -\delta_{\bar{\sigma}\sigma}. \) \( \mathcal{R}_\Delta \equiv [1 + (q + \Delta)/M_V]/2 \) denotes the projection operator, \( \mathcal{R}_\Delta^2 = \mathcal{R}_\Delta \) to ensure the \( S \)-wave \( Q\bar{Q} \) state in the heavy-quark limit \([9, 12]\). (This projection operator coincides with that discussed in Ref. \([10]\) up to the binding-energy effects of the quarkonia.) Note that eq. (2) reduces to the vector-meson WFs of Ref. \([1]\) by the replacement \( \mathcal{R}_\Delta \to 1 \). Essential difference of eq. (2) from the “perturbative” WFs (1) is that the scalar function \( \phi(\alpha, \vec{k}_\perp) \) contains nonperturbative dynamics between \( Q \) and \( \bar{Q} \). Now the light-cone WFs of the \( \eta_Q \) meson can be derived from eq. (2) utilizing spin symmetry, which is exact in the heavy-quark limit. This symmetry relates the \( S \)-wave states, \( \eta_Q \) and the three spin states of the vector meson. Namely, \( M_{\eta_Q} = M_V \), and the pseudoscalar state is related to the vector state with longitudinal polarization as \( |\eta_Q\rangle = 2\tilde{S}_Q^3|V(\bar{\sigma} = 0)\rangle \), where \( \tilde{S}_Q^3 \) is the third component of the hermitean spin operator \( \hat{S}_Q \) which acts on the spin of the heavy quark \( Q \) but does not act on \( \bar{Q} \). This implies that the \( \eta_Q \) WFs are given by the replacement \( u_\lambda \to 2S^3 u_\lambda \) in eq. (2), where \( S^3 \) is a matrix representation of \( \tilde{S}_Q^3 \) as \( S^3 = \gamma_5 (q + \Delta)\gamma^0/(2M_V) \) which is related to a spin matrix \( \sigma^{12}/2 = \gamma_5 \gamma^0 \gamma^3/2 \) in the meson rest frame by a Lorentz boost in the third direction:

\[
\Psi_{\lambda\lambda'}^{\eta_Q^*}(\alpha, k_\perp) = -\frac{\tilde{\nu}_\lambda'(q + \Delta - k)}{\sqrt{1 - \alpha}} \gamma_5 \mathcal{R} u_\lambda(k) \frac{\phi^*(\alpha, \vec{k}_\perp)}{M_{\eta_Q}}.
\]

(3)

This result shows that the \( \eta_Q \) is described by the same nonperturbative WF \( \phi(\alpha, \vec{k}_\perp) \) as the vector meson. We also note that, due to the presence of \( \mathcal{R}_\Delta \), the “\( Q\bar{Q}\eta_Q \) vertex” involves pseudovector as well as pseudoscalar coupling.

Combining our \( Z \) and \( \eta_Q \) WFs with the \( Q\bar{Q}-N \) elastic amplitude \([1]\), we get the total amplitude \( M^{(\xi)} \) for the polarization \( e^{(\xi)} \) of the virtual \( Z \) boson. We find \( M^{(\pm 1)} = 0 \), which reflects conservation of helicity in the high energy limit, and

\[
i M^{(0)} = -\sqrt{\frac{\pi^2 W^2}{N_c}} \frac{g_{Wm_Qc_A}}{M_{\eta_Q} Q \cos \theta_W} \alpha_s(Q^2_{\text{eff}}) \left[ 1 + i \frac{\pi}{2} \frac{\partial}{\partial \ln x} \right] x G(x, Q^2_{\text{eff}}) \int_0^1 \frac{d\alpha \hat{Q}}{\alpha (1 - \alpha)} \int_0^\infty db b^2 \phi^*(\alpha, b) K_1(b \hat{Q}),
\]

(4)
where \( x = (Q^2 + M_{\eta Q}^2)/s \), \( Q_{\text{eff}}^2 = (Q^2 + M_{\eta Q}^2)/4 \), \( \tilde{Q} = [\alpha(1 - \alpha)Q^2 + m_Q^2]^{1/2} \), \( K_1 \) is a modified Bessel function, and \( G(x, Q_{\text{eff}}^2) \) is the conventional gluon distribution. Eq. (4) is written in the “\( \vec{b} \)-space” conjugate to the \( \vec{k}_{\perp} \)-space via the Fourier transformation; \( \vec{b} \) denotes the transverse separation \( (\vec{b} \equiv |\vec{b}|) \) between \( Q \) and \( \tilde{Q} \). In derivation of eq. (4) we retain only the leading \( \ln(Q_{\text{eff}}^2/A_{\text{QCD}}^2) \) contribution, which corresponds to the “color-dipole picture” [2, 13]. As expected, the result (4) is proportional to \( c_A \) so that the \( \eta_Q \) meson is generated by the axial-vector part of the weak current. (For comparison, we also calculate the diffractive vector meson production via the weak neutral current. Using eq. (2), we find that \( M^{(0)} \) and \( M^{(\pm 1)} \) give the production of the longitudinally and transversely polarized vector mesons, respectively, and that all these amplitudes are proportional to the vector coupling \( c_V \).)

Combining eq. (4) with the \( Z \) boson propagator and the weak neutral current by a neutrino, the forward differential cross section for the \( \eta_Q \) production is given by

\[
\frac{d^3\sigma(\nu N \rightarrow \nu N' \eta_Q)}{dsdQ^2dt} \bigg|_{t=0} = \frac{1}{4(8\pi)^3E_{\nu}^3M_N^2\cos^2\theta_W} \frac{Q^2}{(Q^2 + M_Z^2)^2} \frac{H_{+}^2}{\alpha^2} \frac{\varepsilon}{1 - \varepsilon} |M^{(0)}|^2, \tag{5}
\]

where \( E_{\nu} \) is the neutrino beam energy in the lab system, and \( \varepsilon = [4(1 - y) - Q^2/E_{\nu}^2]/[2(1 + (1 - y)^2) + Q^2/E_{\nu}^2] \) with \( y = s/(2M_{NN}E_{\nu}) \). In order to evaluate the corresponding elastic \( \eta_Q \) production rate, we assume the \( t \)-dependence as \( d^3\sigma/dsdt = d^3\sigma/dsdt|_{t=0} \exp(B_{NN}t) \) with a constant diffractive slope, as in the case of the vector meson production. Integrating over \( t \), \( Q^2 \) and \( s \), we get the elastic production rate \( \sigma(\nu N \rightarrow \nu N' \eta_Q) \) as a function of \( E_{\nu} \).

![FIGURE 2](image)

**FIGURE 2.** The elastic production rates as functions of \( E_{\nu} \). Solid, dashed and dotted curves are for \( \eta_c \), \( J/\psi \) and \( \eta_b \), respectively. Note that the dotted curve shows the rate multiplied by \( 10^3 \).

For numerical computation of \( \sigma(\nu N \rightarrow \nu N' \eta_Q) \), we need explicit form of the nonperturbative WF \( \phi(\alpha, \vec{b}) \) of eq. (4). As mentioned above, we can use the corresponding nonperturbative part of the vector-meson WF and we use the one which was constructed in Ref. [13] based on non-relativistic Cornell potential model for heavy quarkonia with \( m_{c(b)} = 1.5(4.9) \) GeV. Also, we use the empirical values for the masses \( M_{\eta_c} = 2.98 \) GeV, \( M_N = 0.94 \) GeV and \( M_Z = 91.2 \) GeV. For \( \eta_b \), we use an estimate \( M_{\eta_b} = 9.45 \) GeV.
GeV [14]. Because the slope $B_{\eta'}$ introduced above is unknown, we assume that $B_{\eta'}$ has the same value as that for the corresponding vector meson: $B_{\eta_c} = 4.5 \text{ GeV}^{-2}$ and $B_{\eta_b} = 3.9 \text{ GeV}^{-2}$ (See Ref. [3]). For the gluon distribution function $G(x, Q^2_{\text{eff}})$ of eq. (4), we employ GRV95 NLO parameterization [15].

We show the elastic $\eta_c$ production rate, $\sigma(\nu N \to \nu' N' \eta_c)$, by the solid curve in Fig. 2. The result monotonically increases as a function of the beam energy $E_{\nu}$. Such behavior is similar with that observed in the $\pi$-production data by the neutral and charged currents [5]. For comparison, we show the elastic $J/\psi$ production rate $\sigma(\nu N \to \nu' N' J/\psi)$ by the dashed curve. The rate for $\eta_c$ production is much larger than that for $J/\psi$ by a factor $\sim 20$. This is mainly due to the relevant weak couplings, $c_A$ for $\eta_c$ and $c_V$ for $J/\psi$, as $(c_A/c_V)^2 \approx 7$. Also, most of the remaining factor arises from the different behavior of the $Z$ light-cone WFs (1) between the axial-vector and vector channels, resulting in a few times difference in the overlap integrals with the corresponding meson WFs. In Fig. 2, we also show the $\eta_b$ production rate $\sigma(\nu N \to \nu' N' \eta_b)$ by the dotted curve. Although the rate for $\eta_b$ is generally much smaller than that for $\eta_c$, the former increases more rapidly than the latter for increasing $E_{\nu}$. Therefore, the $\eta_b$ production rate could become comparable with the $J/\psi$ or $\eta_c$ productions for higher beam energy. It suggests a possibility to observe $\eta_b$ through the diffractive productions by high intensity neutrino beams available at ongoing or forthcoming neutrino factories.

In conclusion, we have computed the diffractive production cross sections of $\eta_c$ and $\eta_b$ mesons via the weak neutral current. Using the new results of the light-cone WFs for $Z$ and $\eta_{c,b}$, the production rates are obtained based on the factorization formalism in QCD. Our results demonstrate that neutrino-induced productions will open a new window to measure $\eta_{c,b}$.

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