

Role of Generalized Sullivan Processes in Semi-Inclusive Kaon Production

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We investigate, in deep inelastic scattering (DIS) by charged leptons, the numerical importance of semi-inclusive kaon production due to the generalized Sullivan process, in which the virtual photon may strike and smash the recoiling baryon upon meson emission of the target proton, as compared to the standard mechanism in the fragmentation picture of Field and Feynman. Our numerical results indicate that the standard fragmentation picture is modified only in a very limited kinematic region (characterized by very small x , very small y , and z close to unity, with x the Bjorken scaling variable, y the fraction of the energy transfer, and z the momentum fraction carried away by the kaon)

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I. Introduction

The parton model was proposed in 1969 by R. P. Feynman [1] for interpreting the results of the MIT-SLAC deep-inelastic scattering (DIS) experiments. Since then, the physics associated with the parton model has been a subject of great interest. In particular, we may mention that Drell, Levy, and Yan [2] in 1969-72 tried to understand the parton model using the field theory language, thereby developing a perturbation theory in the limit of infinite momentum, or better known as the light-cone perturbation theory. Also in 1969, Wilson [3] published the paper on the operator product expansion (OPE), which offers insights concerning the meaning of the parton distributions. Two years later, in 1971, Fritzsche, Gell-Mann, and Leutwyler [4] proposed lightcone algebra. In 1974 Kogut and Susskind [5] tried to connect a renormalized quantum field theory to the framework of the parton picture.

On a different front, asymptotic freedom for the candidate theory of strong interactions, namely quantum chromodynamics or QCD, was discovered in 1973 by Politzer, Gross and Wilczek [6]. Although it is not yet possible to understand in a quantitative manner the (observed) parton distributions from the first principle QCD, we may understand the Q^2 dependence of the parton distributions, at least for large Q^2 , from the Altarelli-Parisi

evolution equations [7], which give rise to the so-called "Scaling violations" as confirmed in detail experimentally in DIS.

In recent years, the efforts seem to indicate that the nonperturbative aspects of QCD are far from being trivially separable for many physical processes. To study properties of hadrons at low and intermediate energies, therefore, it has become an important exercise to prove if a factorization theorem of some sort, i.e. suitable factorization between the perturbative part of QCD and its nonperturbative aspect, can indeed be obtained.

One of the potentially very important nonperturbative aspects has to do with pions and other low-lying pseudoscalar mesons, often recognized as the Goldstone bosons arising naturally from the breaking of the $SU(N_f)_L \times SU(N_f)_R$ chiral symmetry in QCD for N_f flavors of almost massless quarks into the $SU(N_f)_V$ symmetry. These Goldstone particles couple to, e.g., nucleons nonperturbatively and it has become clear that such couplings cannot be reproduced in any existing perturbative QCD calculations. Thus, effects caused by the pion clouds associated with nucleons (or with other hadrons) are likely nonperturbative in nature and would have to be carefully identified in any effort of isolating nonperturbative aspects from QCD. One imperative question is whether it is possible to go over to high energy limits so as to avoid dealing directly with such nonperturbative pion exchange effect. As we shall see, we suspect that the answer to such important question may be highly nontrivial.

An important observation may have been made in 1972 by Sullivan [8], who considered how the presence of a pion cloud in a proton may influence extraction of the parton distributions. In Fig. 1, we illustrate one such process, in which the virtual photon strikes and smashes the pion in the cloud associated with the nucleon. Following Sullivan's derivation [8], we have

$$\delta F_{2N}^\pi(x, Q^2) = \int_x^1 dy f_\pi(y) F_{2\pi} \left(\frac{x}{y}, q^2 \right), \quad (1)$$

$$f_\pi(y) = \frac{3}{16\pi^2} \left(\frac{f_{\pi NN} 2m_N}{\mu} \right)^2 y \int_{-\infty}^{t^m} dt \frac{(-t) |F_\pi(t)|^2}{(-t + \mu^2)^2}, \quad (2)$$

with $t^m = -\frac{m_N^2 y^2}{(1-y)}$. Here m_N and μ are, respectively, the nucleon mass and the pion mass. $F_{2\pi}(x)$ is the pion structure function, as would be measured in DIS with the pion as the target. $\delta F_{2N}(x)$ is the correction to the nucleon structure function due to the given Sullivan process. $f_\pi(y)$ is the probability of finding a pion with momentum fraction y of the nucleon momentum. $F_\pi(t)$ is the form factor which characterizes the t -dependence of the πNN coupling.

The meson-exchange picture has been employed by many authors in a variety of contexts. For instance, Thomas and Frankfurt *et al.* [9] have attempted to attribute the observed asymmetry $\frac{1}{2}(\bar{u}(x) + \bar{d}(x)) - \bar{s}(x)$ to Sullivan processes and obtained a nucleon bag radius of about 0.8 fermi or a monopole πNN form factor of about 500 MeV (an unusually soft form factor). On the other hand, the so-called EMC effect, observed originally by the EMC collaboration [10] and later confirmed by others [11-13], may be explained by the pion-excess picture [14-18] which is based upon a treatment of the Sullivan process in

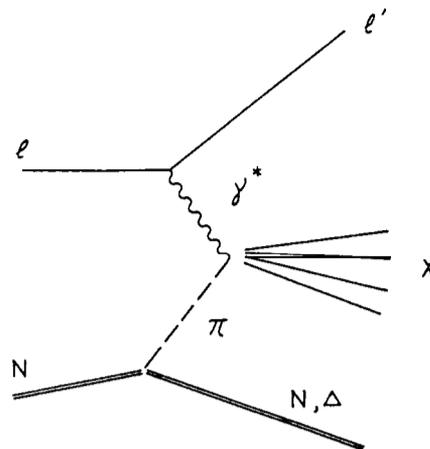


FIG.1. The original process is shown, in which the virtual photon strikes and smashes the pion in the cloud associated with the nucleon.

nuclear medium (as contrasted with that in free space). Very recently, Henley and Miller and several other authors [19-21] have related Sullivan processes to the violation of the Gottfried sum rule [22], as observed by the NMC group [23]. Furthermore, the possibility for understanding, as depolarization of the valence quark spins due to presence of Goldstone pions, the so-called "proton spin crisis", which is now established firmly in a series of recent experiments [24], was first pointed out in [25] and has been treated in a quantitative manner in [26].

In contrast with the idea of a parton model in which a hadron consists of partons of different kinds, the fragmentation hypothesis of Field and Feynman [27] states that a high-energy parton of a specific kind (a quark, an antiquark, or a gluon) has different probabilities of fragmenting into different hadrons. In such fragmentation scheme, the fragmentation function $D_q^h(z)$ is introduced for a given parton of flavor q and every hadron h , with z the fraction of the parton's longitudinal momentum carried away by the hadron. The fragmentation picture is to be used in many high energy processes, including semi-inclusive hadron production in deep inelastic scattering (DIS). An interesting example is semi-inclusive kaon production in DIS by charged leptons (e or μ):

$$l + p \longrightarrow l' + K^+ + X. \quad (3)$$

Here for the sake of definiteness we use semi-inclusive K^+ production with the proton target. The process may be depicted pictorially in the fragmentation picture by Fig. 2(a). However, we may consider a "generalized" Sullivan process as depicted by Fig. 2(b), in which the recoiling baryon A or Σ^0 is struck and smashed. It is clear that Fig. 2(b) should be considered as an additional mechanism for semi-inclusive production of K^+ in DIS. If the Sullivan process plays an essential role in connection with the parton model, it is a natural question to investigate whether the generalized Sullivan process as shown by Fig. 2(b) may

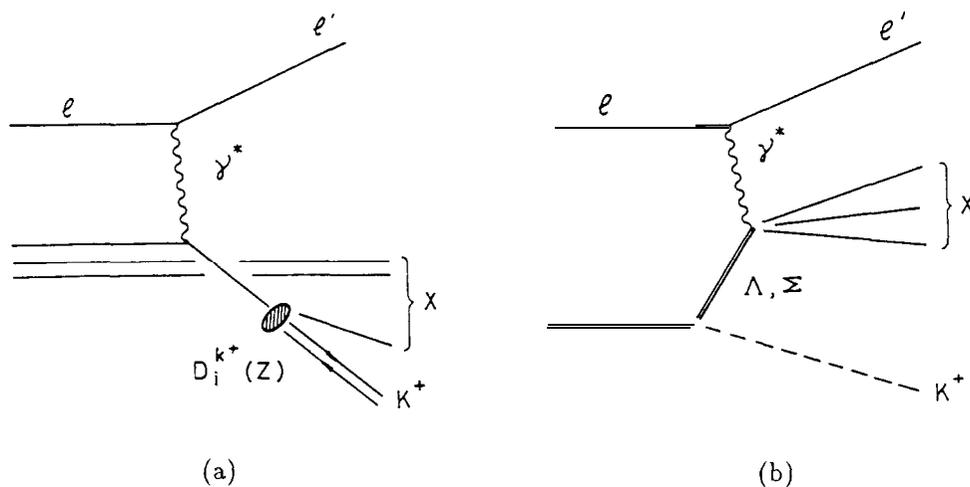


FIG. 2. (a)-(b). Semi-inclusive kaon production through (a) the fragmentation picture of Field and Feynman and (b) the generalized Sullivan processes.

give rise to some important contribution as compared to those given in the fragmentation picture [Fig. 2(a)]. This is the primary focus of the present paper.

At this juncture, it might be useful to consider some issues in relation to generalizations of Sullivan processes. The wave function for the target proton contains in principle (infinitely) many Fock components such as $|uud\rangle$ (the "bare" proton), $|N\pi\rangle$, $|\Delta\pi\rangle$, $|\Lambda K^+\rangle$, etc., each with certain amplitudes corresponding to specific vertices (in momentum space). The information regarding the offshell feature is in principle described by the form factor such as $F_\pi(t, p_i^2, p_f^2)$ in Eq. (2). Note that such form factor is a function of Lorentz invariants and it depends only on the variable t because we put the target and recoil baryons on shell. In configuration space, such information is represented by the deeply bound nature of a specific Fock component. Accordingly, there is no more off-shell information in relation to the wave function, or the parton substructure, of an individual hadron. On the other hand, once we probe the system at high enough energies with Q^2 considerably larger than the scale set by the binding energy or the off-shellness, it is clear that we need not concern ourselves with the question of how the off-shell nature modifies the parton distributions, or modifies the particle itself. As a result, it is the knowledge on the form factors such as $F_\pi(t, p_i^2, p_f^2)$, which is required in generalizing Sullivan processes.

It may also be useful to mention some recent investigations related to the present work. Korpa, Dieperink, and Scholten [28] have considered the pion contribution to semi-inclusive deep inelastic scattering with slow baryons (N' s and A' s) in the final state, while Hwang and Wen [29] have considered possible importance of generalized Sullivan processes in semi-inclusive Λ production. In these investigations the Sullivan process in which the mesons in the cloud are struck and smashed while in this work it is the recoiling baryon which is struck and smashed. Identification of the recoil baryons (especially non-strange baryons) could be a real challenge in avoiding possible backgrounds, but the effects might be more prominent [28, 29] than that proposed in the present work. In any event, it is clear that this investigation is complimentary to previous ones [28, 29] and may help to shed

light on directions for future research.

II. Formulation

We introduce the standard variables, with l , l' , and P the four-momenta of the initial lepton, the final lepton, and the target proton, respectively.

$q \equiv l - l'$: four-momentum transfer;

$Q^2 \equiv -q^2$: four-momentum transfer squared;

$\nu \equiv \frac{q \cdot P}{M_P} = E_l - E_{l'}$: energy transfer as seen in the proton rest system;

$x \equiv \frac{Q^2}{2q \cdot P} = \frac{Q^2}{2M_P \nu}$: Bjorken scaling variable;

$y = \frac{q \cdot P}{l \cdot P} = \frac{E_l - E_{l'}}{E_l}$.

The cross section for the generalized Sullivan process, as illustrated by Fig. 2(b), may be thought of as the cross section for $\ell + h \rightarrow \ell' + X$ convoluted with the probability of finding the hadron h inside the target proton. To write down the formula for the cross section, we introduce the light-cone variables:

$$\begin{aligned} P^+ &= \frac{1}{\sqrt{2}}(P^0 + P^3), \\ P^- &= \frac{1}{\sqrt{2}}(P^0 - P^3), \\ x_h &= \frac{p_q^+}{P_h^+}, \\ y_h &= \frac{P_h^+}{P^+}. \end{aligned} \tag{5}$$

Here p_q , P_h , and P are respectively the four-momenta of the struck quark q , the recoiling hadron h , and the target proton. In the infinite-momentum frame, we have $x = x_h y_h$.

The double differential cross section for $\ell + h \rightarrow \ell' + X$ is given by [30]

$$\frac{d^2\sigma_h}{dQ^2 d\nu} = \frac{E_{l'}}{E_l} \left(\frac{e^2}{q^2} \right)^2 \frac{1}{\pi} \{W_{2h} \cos^2 \theta/2 + 2W_{1h} \sin^2 \theta/2\}, \tag{6}$$

with

$$\begin{aligned} W_{2h} &= \int dx_h \sum_i f_{i/h}(x_h) \delta \left(x_h - \frac{Q^2}{2y_h P \cdot q} \right) \cdot x_h \frac{1}{4\nu} \cdot e_i^2, \\ 2W_{1h} &= \int dx_h \sum_i f_{i/h}(x_h) \delta \left(x_h - \frac{Q^2}{2y_h P \cdot q} \right) \frac{1}{4y_h M_P} e_i^2, \end{aligned}$$

We introduce

$$\begin{aligned} F_{1h}^T(x_h, Q^2) &= m_h W_{1h}^T, \\ F_{2h}^T(x_h, Q^2) &= \nu W_{2h}^T, \end{aligned} \quad (7)$$

SO that

$$\frac{d^2\sigma_h}{dQ^2 d\nu} = \frac{E E_l'}{E_l} \left(\frac{e^2}{q^2} \right)^2 \frac{1}{\pi} \left\{ \frac{1}{\nu} F_{2h}^T(x_h, Q^2) \cos^2 \theta/2 + \frac{1}{y_h M_P} 2F_{1h}^T(x_h, Q^2) \sin^2 \theta/2 \right\}.$$

Changing the variables from (Q^2, ν) to (x, y) , we obtain

$$\begin{aligned} \frac{d^2\sigma_h}{dx dy} &= \frac{d^2\sigma_h}{dQ^2 d\nu} \frac{Q^2}{xy} \nu \\ &= \left(\frac{e^2}{q^2} \right)^2 \frac{2E_l M_P}{\pi} \left\{ (1-y) F_{2h}^T(x_h, Q^2) + y^2 \frac{x}{y_h} F_{1h}^T(x_h, Q^2) \right\}. \end{aligned} \quad (8)$$

Using Eq. (7), we have

$$\begin{aligned} F_{2h}^T(x_h, Q^2) &= 2x_h F_{1h}^T(x_h, Q^2) \\ &= 2 \frac{x}{y_h} F_{1h}^T(x_h, Q^2). \end{aligned} \quad (9)$$

Combining Eqs. (8) and (9), we have [30]

$$\frac{d^2\sigma_h}{dx dy} = \left(\frac{e^2}{q^2} \right)^2 \frac{2E_l M_P}{\pi} \left\{ \left(1 - y + \frac{1}{2} y^2 \right) F_{2h}^T \left(\frac{x}{y_h}, Q^2 \right) \right\}. \quad (10)$$

We may now proceed to write down the cross section on the generalized Sullivan process as illustrated by Fig. 2(b).

$$\begin{aligned} \frac{d^2\sigma}{dx dy} &= \int dk_{\perp}^2 \int dy_h \frac{1}{16\pi^2} y_h (1-y_h) \frac{|V_{P \rightarrow \Lambda K}|^2}{E^2} \\ &\cdot \left(\frac{e^2}{q^2} \right)^2 \frac{2E_l M_P}{\pi} \left\{ \left(1 - y + \frac{1}{2} y^2 \right) F_{2h}^T \left(\frac{x}{y_h}, Q^2 \right) \right\}, \end{aligned} \quad (11)$$

with $|V_{P \rightarrow \Lambda K}|^2$ the absolute square of the $P \rightarrow \Lambda K^+$ vertex as given by

$$\begin{aligned} |V_{P \rightarrow \Lambda K}|^2 &= |-ig\bar{u}(p_{\Lambda})\gamma_5 u(p_P)|^2 \\ &= \frac{1}{2} \sum_{spin} |V_{P \rightarrow \Lambda K}|^2 \\ &= \frac{g^2}{2} \{4p_P \cdot p_{\Lambda} - 4M_P m_{\Lambda}\} \\ &= \frac{1}{y_h} \cdot g^2 \{y_h^2 M_P^2 + k_{\perp}^2 + m_{\Lambda}^2 - 2y_h M_P m_{\Lambda}\}. \end{aligned} \quad (12)$$

It is related to the probability of finding the hadron h in a ("dressed") physical proton:

$$f_{h/P}(y_h, k_{\perp}^2) dk_{\perp}^2 dy_h = \frac{dk_{\perp}^2 dy_h}{16\pi^2} y_h(1-y_h) \frac{|V_{P \rightarrow \Lambda K}|^2}{\tilde{E}^2}, \quad (13a)$$

with

$$\tilde{E} = \{M_P^2 y_h(1-y_h) - m_{\Lambda}^2(1-y_h) - m_K^2 y_h - k^2\}. \quad (13b)$$

Using the light-cone variables, we have, for a general vertex $A \rightarrow B + C$,

$$\begin{aligned} 1 - y_h &= \frac{P^+ - P_B^+}{P^+} \\ &= \frac{P_C^+}{P^+} \\ &= \frac{E_C^{lab} + P_{\epsilon 3}^{lab}}{M_A} \\ &= \frac{1}{M_A} \left(2E_C^L - \frac{k_{\perp}^2 + m_C^2}{2(1-y_h)M_A} \cdot 2 \right). \end{aligned} \quad (14)$$

Or, we have, with $z = E_h/\nu$,

$$\begin{aligned} k_{\perp}^2 &= 2z\nu(1-y_h)M_A - (1-y_h)^2 M_A^2 - m_C^2; \\ dk_{\perp}^2 &= 2\nu(1-y_h)M_A dz. \end{aligned} \quad (15)$$

Note that Eq. (15) gives rise to a constraint on z :

$$\begin{aligned} z &= \frac{1}{2\nu} \left\{ (1-y_h)M_A + \frac{k_{\perp}^2 + m_C^2}{(1-y_h)M_A} \right\}, \\ \frac{1}{2\nu} \left\{ (1-y_h)M_A + \frac{m_C^2}{(1-y_h)M_A} \right\} &< z < 1. \end{aligned} \quad (16)$$

Or, we have, for $z < \frac{M_A^2 + m_C^2}{2\nu M_A}$,

$$1 - \frac{2z\nu + \sqrt{(2z\nu)^2 - 4m_C^2}}{2M_A} < y_h < 1 - \frac{2z\nu - \sqrt{(2z\nu)^2 - 4m_C^2}}{2M_A} \quad (17)$$

To ensure that the factor inside the square root is non-negative, we have, for $z > \frac{M_A^2 + m_C^2}{2\nu M_A}$,

$$0 < y_h < 1 - \frac{2z\nu - \sqrt{(2z\nu)^2 - 4m_C^2}}{2M_A} \quad (18)$$

Finally, there is an additional constraint:

$$x < y_h < 1. \quad (19)$$

Changing the variables k^2, y_h to z, y_h in Eq. (11), we finally obtain

$$\frac{d^3\sigma}{\mathbf{d}\mathbf{x}\mathbf{d}\mathbf{y}\mathbf{d}\mathbf{z}} = \int_{y_h^{\min}}^{y_h^{\max}} dy_h \cdot 2\nu(1-y_h)M_P \cdot \frac{1}{16\pi^2} y_h(1-y_h) \frac{|V_{P \rightarrow \Lambda K}|^2}{\tilde{E}^2} \cdot \left(\frac{e^2}{q^2}\right)^2 \frac{2E_l M_P}{\pi} \left\{ \left(1-y+\frac{1}{2}y^2\right) F_{2h}^T\left(\frac{x}{y_h}, Q^2\right) \right\}. \quad (20)$$

The differential cross section as given by Eq. (20) is to be contrasted with what we may expect from the fragmentation picture of Field and Feynman [27]:

$$\frac{d^2\sigma(e+P \rightarrow h+X+e)}{\mathbf{d}\mathbf{x}\mathbf{d}\mathbf{y}} = \left(\frac{e^2}{q^2}\right)^2 \frac{2E_l M_P}{\pi} \left\{ \left(1-y+\frac{1}{2}y^2\right) \tilde{F}_2(x, z, Q^2) \right\}, \quad (21)$$

where

$$\tilde{F}_2(x, z, q^2) = \int dz \sum_i e_i^2 x q_i(x) D_i^h(z). \quad (22)$$

Here $D_i^h(z)$ is the fragmentation function which describes the probability of the quark i fragmenting into the hadron h , with z the momentum fraction of the hadron relative to the quark struck by the virtual photon. ($z \equiv \mathbf{P} \cdot \mathbf{P}_h / P \cdot q \equiv E_h / \nu$.) Differentiating Eq. (21) with respect to dz , we obtain the differential cross section as follows:

$$\frac{d^3\sigma(e+P \rightarrow \mathbf{h}+\mathbf{X}+\mathbf{e})}{\mathbf{d}\mathbf{x}\mathbf{d}\mathbf{y}\mathbf{d}\mathbf{z}} = \left(\frac{e^2}{q^2}\right)^2 \frac{2E_l M_P}{\pi} \left\{ \left(1-y+\frac{1}{2}y^2\right) \sum_i e_i^2 x q_i(x) D_i^h(z) \right\}. \quad (23)$$

In what follows, we shall consider numerically the ratio of the cross section given by Eq. (20) to that given by Eq. (23); such ratio is to be investigated as a function of the variables x , y , and z .

III. Sample numerical results

To investigate the importance of the generalized Sullivan process as shown by Fig. 2(b) as compared to the standard fragmentation picture of Fig. 2(a), we proceed to consider the ratio of the cross section given by Eq. (20) to that given by Eq. (23), treating the ratio as a function of the variables x , y , and z . To this end, we shall adopt, as the input for Eq. (20), the parameters (with the dipole form factor) used by Hwang, Speth, and Brown [20] and, as the input for Eq. (23), the fragmentation functions used by Owens, Reya, and Glück[31]. For the sake of our considerations, we may neglect possible Q^2 -dependence of the fragmentation functions [32]. Specifically, the fragmentation functions may be parametrized as follows:

$$\begin{aligned}
 zD_{K^+/\bar{s}} &= a\sqrt{z}(c-z) + \xi_K(1-z)^2, \\
 zD_{K^+/\bar{u}} &= b\sqrt{z}(c-z) + \xi_K(1-z)^2, \\
 zD_{K^+/\bar{v}} &= zD_{K^+/\bar{d}} = zD_{K^+/\bar{l}} = zD_{K^+/\bar{s}} \\
 &= \xi_K(1-z)^2,
 \end{aligned}
 \tag{24}$$

with $a = \frac{9}{40}$, $b = \frac{9}{80}$, $C = \frac{11}{9}$, and $\xi_K = 0.5 \cdot \frac{61}{125}$. [31]

In Fig. 3, we plot the ratio of the cross section given by Eq. (20) to that given by Eq. (23) as a function of y and z for a fixed value of x ($x = 0.05$). Here the Sullivan processes with both the intermediate A and Σ have been taken into account, although the contribution from Σ is fairly small by comparison (typically only a couple of percent of the A contribution). Fig. 3 indicates that, only in the kinematic region in which z is close to

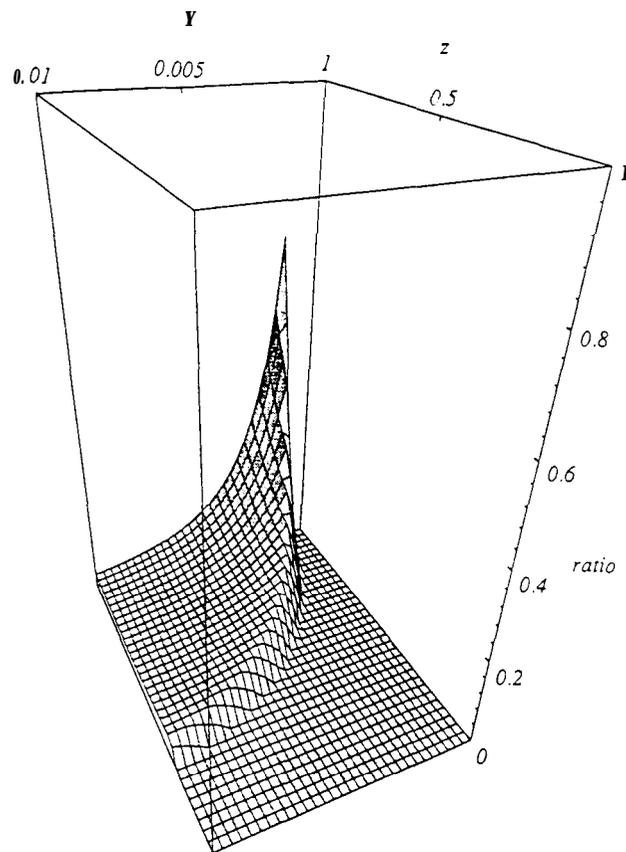


FIG. 3. The ratio of the cross section given by Eq. (20) to that given by Eq. (23) is plotted as a function of y and z for a fixed value of x ($x = 0.05$). Here the Sullivan processes with both the intermediate A and Σ are taken into account.

unity and y is rather small (say $y \leq 0.005$), the contribution due to the Sullivan process may be of some numerical importance. However, it is not clear that this rather limited kinematic region can be accessed in a present-day experiment.

In Fig. 4, we plot the ratio of the cross section given by Eq. (20) to that given by Eq. (23) as a function of z and x for $y = 0.05$. Here we show that the contribution due to the Sullivan process is of importance provided that z is close to unity and x is rather small (say $x \leq 0.01$).

In Fig. 5, we plot the ratio of the cross section given by Eq. (20) to that given by Eq. (23) as a function of x and y for $z = 1.0$. Here we find that the contribution due to the Sullivan process may be of numerical importance only in an extremely small kinematic region which may not be currently accessible in present-day experiments.

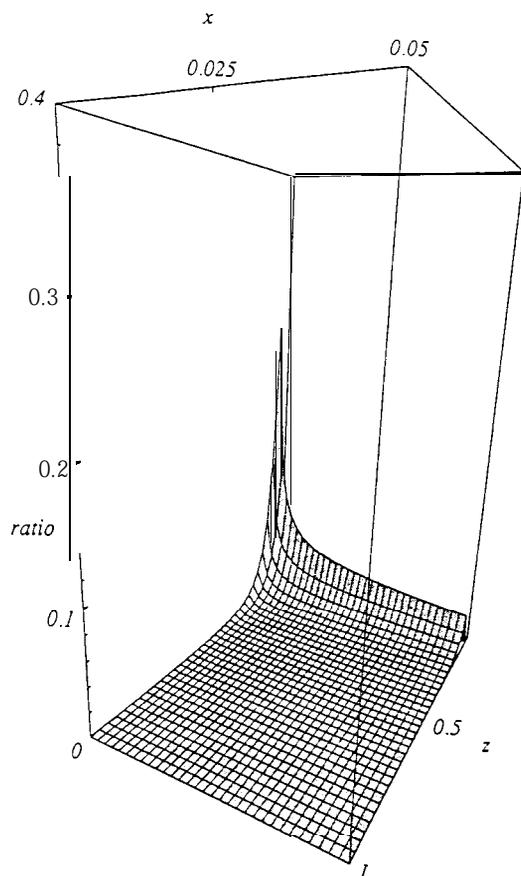


FIG. 4 The ratio of the cross section given by Eq. (20) to that given by Eq. (23) is plotted as a function of z and x for $y = 0.05$.

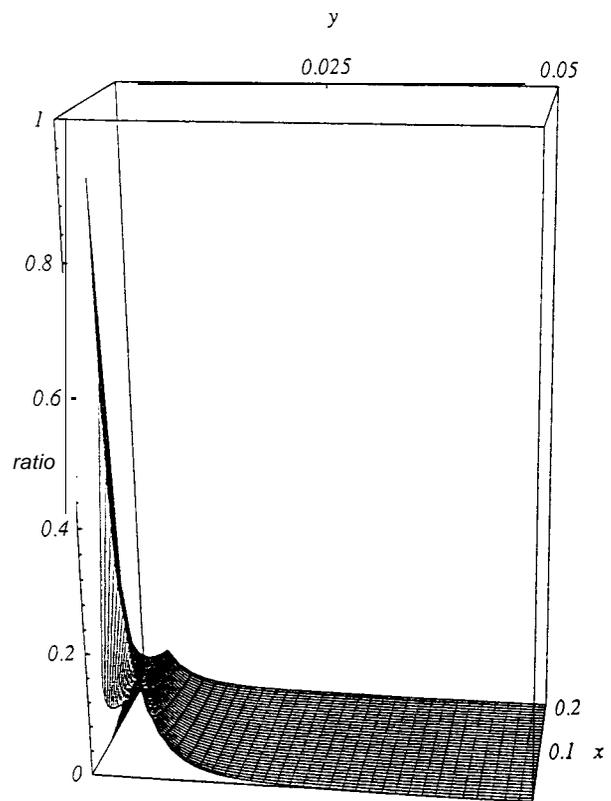


FIG. 5 The ratio of the cross section given by Eq. (20) to that given by Eq. (23) is plotted as a function of x and y for $z = 1.0$.

Figs. 3-5 indicate that, only in a very limited kinematic region as characterized generally by very small x and very small y , the contribution from the generalized Sullivan process may be of numerical importance as compared to that due to the standard fragmentation picture of Fig. 2(a). The kinematic region is limited that it might be only of limited experimental interest. As compared to the related investigations [28,29] where one finds that the Sullivan process is of some numerical importance, we may understand the present result by noting that effects are usually much smaller while the relevant kinematic variables x and y get squeezed to a region of very smaller values when the heavier final fragment (baryons rather than mesons) is struck and smashed.

IV. Summary

In light of the observation [20] that the generalized Sullivan process, in which the virtual photon may strike and smash the recoiling baryon upon meson emission of the target proton, may play an important role in deep inelastic scattering (DIS) by charged leptons and thus may affect significantly extraction of parton distributions from experiments, we investigate in this paper the numerical importance of semi-inclusive kaon production due to the generalized Sullivan process as compared to the standard mechanism in the fragmentation picture of Field and Feynman. Our numerical results, as illustrated in Figs. 3-5, indicate that the standard fragmentation picture is modified only in a very limited kinematic region (characterized by very small x , very small y , and z close to unity, with x the Bjorken scaling variable, y the fraction of the energy transfer, and z the momentum fraction carried away by the kaon).

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