WEAK HYPERON PRODUCTION IN $ep$ SCATTERING*

XuemIn Jin † and R. L. Jaffe ‡
Center for Theoretical Physics
Laboratory for Nuclear Science
and Department of Physics
Massachusetts Institute of Technology
Cambridge, Massachusetts 02139
(MIT-CTP-2594, hep-th/9612316. December 1996)

Abstract

We study the kinematics and cross section of the scattering process $ep \rightarrow e\Sigma^+$. The cross section is expressed in terms of complex form factors characterizing the hadron vertices. We estimate the cross section for small momentum transfer using known experimental information. To first order in the momentum transfer, we obtain a model independent result for the photon-exchange part of the cross section, which is completely determined by the decay width $\Gamma(\Sigma^+ \rightarrow p\gamma)$. For the kinematics of the parity violation experiment at MAMI, this first order result gives rise to a ratio of $(d\sigma/d\Omega)_{ep \rightarrow e\Sigma^+} / (d\sigma/d\Omega)_{ep \rightarrow ep} \simeq 4.0 \times 10^{-15}$. The $Z^0$-exchange and interference parts give much smaller contributions due to the suppression of the flavor changing weak neutral current in the standard model. Feasibility of the experimental measurement is briefly discussed.

Submitted to: Physical Review D

†Email address: jin@ctp02.mit.edu

‡Email address: jaffe@mitlns.mit.edu

*This work is supported in part by funds provided by the U.S. Department of Energy (D.O.E.) under cooperative research agreement #DF-FC02-94ER40818.
I. INTRODUCTION

Currently, experimentalists at MAMI in Mainz are considering the possibility of measuring the scattering processes $ep \rightarrow e\Sigma^+$ or $ed \rightarrow ep\Lambda$ [1]. To our knowledge, there have been no theoretical investigations of such processes. It is thus important to explore the physics and the feasibility of the experimental measurement. Here we discuss the kinematics and cross section for $ep \rightarrow e\Sigma^+$. With minor alterations our analysis applies as well to $en \rightarrow e\Lambda$.

In lowest order the scattering process $ep \rightarrow e\Sigma^+$ proceeds via the exchange of one photon or one $Z^0$ boson (see Fig. 1). Thus, the cross section consists of a pure $\gamma$, pure $Z^0$ and interference parts. The last two depend on the physics at the $pZ^0\Sigma^+$ vertex, a classic flavor changing weak neutral current, which is severely suppressed in the standard model. On the other hand, the photon-exchange part, $p\gamma\Sigma^+$, includes all gauge interactions of standard model: strong, weak, and electromagnetic. The same vertex appears in the weak radiative decay $\Sigma^+ \rightarrow p\gamma$, which has been well measured experimentally [2]. Despite substantial theoretical effort [3], hyperon weak radiative decays remain poorly understood. Measurement of $ep \rightarrow e\Sigma^+$ might provide more information about the vertex $p\gamma\Sigma^+$ and hence constrain theoretical models.

We express the cross section in terms of various form factors, which reflect the structure of the hadron vertices. We then estimate the cross section at small four-momentum transfer using known experimental information on the weak radiative decay, $\Sigma^+ \rightarrow p\gamma$, and the flavor changing weak neutral current. To first order in the momentum transfer, we obtain a model independent result for the photon-exchange part, which is completely determined by the weak radiative decay width $\Gamma(\Sigma^+ \rightarrow p\gamma)$. For the kinematics of the parity violation experiment at MAMI ($\theta \simeq 35^0$ and $q^2 \simeq -0.237$ GeV$^2$ for the elastic scattering $ep \rightarrow ep$, with $\theta$ the scattering angle and $q^2$ the squared four-momentum transfer) [4], this first order result leads to a suppression factor of $4 \times 10^{-15}$ relative to the $ep$ elastic scattering. The physical reasons for the suppression are the factor of $G_F^2$ from the weak hamiltonian and a factor of $q^2$ from electromagnetic gauge invariance that suppresses the cross section at small momentum transfer. Using the experimental result for the branching ratio of $K^0_L \rightarrow \mu^+\mu^-$, we estimate that the $Z^0$-exchange and interference parts give negligible contributions. We shall discuss briefly the feasibility of experimental measurements at available facilities.

This paper is organized as follows: In Sec. II we discuss the kinematics and derive the cross section. Sec. III gives an estimate of the cross section. Sec. IV is devoted to summary and conclusion.

II. KINEMATICS AND CROSS SECTION

The kinematics for the process $ep \rightarrow e\Sigma^+$ is illustrated in Fig. 1. The four momentum of the initial and final states are denoted by $k = \{E,k\}$ for the initial electron, $P = \{E_p,p\}$ for the target (proton) ($P = \{M_p,0\}$ in the target rest frame), $k' = \{E',k'\}$ for the outgoing electron, and $P' = \{E'_p,p'\}$ for the outgoing hyperon ($\Sigma^+$). One can then define the two invariants

---

1. The text is extracted from a scientific paper discussing the possibility of measuring scattering processes involving hyperons at MAMI in Mainz. It emphasizes the importance of exploring the physics behind such measurements and the feasibility of the experimental setup.

2. Key points include:
   - The consideration of $ep \rightarrow e\Sigma^+$ and $ed \rightarrow ep\Lambda$ processes.
   - The lack of theoretical investigations for these processes.
   - The significance of exploring the physics and feasibility of experimental measurements.
   - The use of form factors to describe the structure of hadron vertices.
   - The estimation of the cross section at small momentum transfer using experimental information on weak radiative decays.
   - The suppression factor of $4 \times 10^{-15}$ relative to the elastic $ep$ scattering.
   - The physical reasons for the suppression due to $G_F^2$ from the weak hamiltonian and $q^2$ from electromagnetic gauge invariance.
   - The estimation of negligible contributions from $Z^0$-exchange and interference parts.

3. The paper is organized into sections discussing kinematics, cross section estimation, and feasibility of experimental measurements.
\[ q^2 \equiv (k-k')^2 = (P' - P)^2 = -4EE' \sin^2(\theta/2) = -Q^2 < 0 \quad (2.1) \]

\[ \nu \equiv P \cdot q = M_p(E - E') \quad (2.2) \]

where the electron mass has been neglected (and will be henceforth). The scattering angle \( \theta \) has been indicated in Fig. 1. Unless otherwise noted, the target rest frame will be adopted.

Since the final \( \Sigma^+ \) state is on-shell, one finds

\[ 2\nu + q^2 = M_{\Sigma}^2 - M_p^2 \quad (2.3) \]

\[ E' = \frac{M_pE - \frac{1}{2}(M_{\Sigma}^2 - M_p^2)}{M_p + 2E\sin^2(\theta/2)} \quad (2.4) \]

Note that there are only two independent variables, \( E \) and \( \theta \), and all the other quantities can be expressed in terms of them. The kinematic domain is thus given by the two conditions \( q^2 < 0 \) (\( Q^2 > 0 \)) and \( 2\nu = M_{\Sigma}^2 - M_p^2 - q^2 \). Since \( E' > 0 \), one has

\[ E > \frac{1}{2M_p} \left( M_{\Sigma}^2 - M_p^2 \right) \quad (2.5) \]

which implies a minimum initial electron energy of \( E_{\text{min}} \approx 285 \text{ MeV} \).

In the discussions to follow, we shall consider the scattering process with unpolarized beam and target and with final spins unobserved. The extension to other situations is straightforward. The differential cross section in this case can be written as:

\[ d\sigma = \frac{e^4}{4EQ^4} \int \frac{d^3k'}{(2\pi)^3} \frac{d^3p'}{(2\pi)^3} \frac{M_{\Sigma}}{E_p} \frac{1}{E'} \delta^4(P' + k' - P - k) \sum_{i=\gamma,\gamma,Z,Z} \eta^{ij}_{\mu\nu} W^j_{\mu\nu} \quad (2.6) \]

\[ \begin{array}{c}
  e(k) \\
  \theta \\
  \gamma, Z^0(q) \\
  p(P) \\
  \Sigma^+(P') \end{array} \]

FIG. 1. Kinematics of the scattering \( ep \rightarrow e\Sigma^+ \) in the proton rest frame.
where
\[ \eta^\gamma = 1, \quad \eta^{\gamma Z} = \frac{1}{\sin^2 2\theta_W} \frac{Q^2}{Q^2 + M_Z^2}, \quad \eta^Z = (\eta^{\gamma Z})^2, \] (2.7)
with \( \theta_W \) the weak angle and \( M_Z \) the \( Z^0 \) mass. Here we use the normalization conventions of Itzykson and Zuber [3]. The leptonic tensor \( l^{\mu\nu}_l \) is simply given by
\[ l^{\mu\nu}_l = k^\mu k'^{\nu} + k'^\mu k^\nu - k \cdot k' g^{\mu\nu}, \] (2.8)
\[ l^{\mu\nu}_Z = 2g^e_\gamma (k^\mu k'^{\nu} + k'^\mu k^\nu - k \cdot k' g^{\mu\nu}) - 2g^e_\alpha i\epsilon^{\mu\nu\rho\sigma} k^\rho k_\sigma, \] (2.9)
\[ l^{\mu\nu}_Z = [(g^e_\gamma)^2 + (g^e_\alpha)^2] (k^\mu k'^{\nu} + k'^\mu k^\nu - k \cdot k' g^{\mu\nu}) - 2g^e_\gamma g^e_\alpha i\epsilon^{\mu\nu\rho\sigma} k^\rho k_\sigma, \] (2.10)
where
\[ g^e_\gamma = -\frac{1}{2} + 2 \sin^2 \theta_W, \quad g^e_\alpha = -\frac{1}{2}. \] (2.11)

On the other hand, the hadronic tensor \( W^{\mu\nu}_h \) describes the complicated structure of the vertices \( p\gamma \Sigma^+ \) and \( pZ^0 \Sigma^+ \),
\[ W^{\gamma\mu\nu}_h = \sum_{\text{spins}} \langle \Sigma^+ | J^\gamma_\mu | p \rangle \langle \Sigma^+ | J^\gamma_\nu | p \rangle^*, \] (2.12)
\[ W^{\gamma Z\mu\nu}_h = \sum_{\text{spins}} \left[ \langle \Sigma^+ | J^\gamma_\mu | p \rangle \langle \Sigma^+ | J^Z_\nu | p \rangle^* + \langle \Sigma^+ | J^Z_\mu | p \rangle \langle \Sigma^+ | J^\gamma_\nu | p \rangle^* \right], \] (2.13)
\[ W^{Z\mu\nu}_h = \sum_{\text{spins}} \langle \Sigma^+ | J^Z_\mu | p \rangle \langle \Sigma^+ | J^Z_\nu | p \rangle^*, \] (2.14)
where \( J^\gamma_\mu \) and \( J^Z_\mu \) denote the electromagnetic and weak neutral currents, respectively. The matrix element \( \langle \Sigma^+ | J^j_\mu | p \rangle \{ j = \gamma, Z \} \), upon the use of Lorentz covariance, takes the following form:
\[ \langle \Sigma^+ | J^j_\mu | p \rangle \equiv \bar{U}(P') \left\{ \gamma^j \gamma F^j_1(q^2) + q_\mu B^j(q^2) + \frac{i\sigma_{\mu\nu} q^\nu}{M_p + M_\Sigma} F^j_2(q^2) \right. \]
\[ + \gamma_5 \gamma_\mu F^j_3(q^2) + \gamma_5 q_\mu B^j_3(q^2) + \frac{i\gamma_5 \sigma_{\mu\nu} q^\nu}{M_p - M_\Sigma} F^j_4(q^2) \left\} U(P), \] (2.15)
where \( U(P') \) and \( U(P) \) are the Dirac spinors of \( \Sigma^+ \) and proton, respectively. The form factors are in general complex. While \( F^j_1, F^j_2, \) and \( B^j \) are parity-conserving, \( F^j_3, F^j_4, \) and \( B^j_3 \) are parity-violating.

From electromagnetic current conservation, one obtains the following relations:
\[ (M_p - M_\Sigma) F^\gamma_1 = q^2 B^\gamma, \quad (M_p + M_\Sigma) F^\gamma_3 = q^2 B^\gamma_3. \] (2.16)
Since there is no zero-mass particle involved except the photon and the photon propagator has been included explicitly, Eq. (2.16) implies that
Note that the condition $F_1^j(0) = 0$ is different from the usual $F_1^j(0) = 1$ (= electric charge of the proton) seen in the proton electromagnetic form factors. This difference arises from gauge invariance. While the proton couples minimally to the photon field $A_\mu$, the $p\gamma\Sigma^+$ vertex must be proportional to $F_{\nu\mu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. Thus, there are only four independent form factors at the vertex $p\gamma\Sigma^+$, which we choose as $F_1^j, F_2^j, F_3^j$, and $F_4^j$. Since $J_\mu^Z$ is not conserved, the six form factors describing the vertex $pZ^0\Sigma^+$ are in general independent.

Lorentz covariance and electromagnetic current conservation allows one to separate $W_{\mu\nu}$ into three distinct structures:

$$
W_{\mu\nu}^\gamma \equiv \left( -g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \right) \frac{M_p}{M_\Sigma} W_1^\gamma(q^2) + \left[ \left( P_\mu - \frac{P \cdot q}{q^2} q_\mu \right) \left( P_\nu - \frac{P \cdot q}{q^2} q_\nu \right) - i\epsilon_{\mu
u\rho\sigma} P^\rho q^\sigma \frac{W_3^\gamma(q^2)}{2M_p M_\Sigma} \right] \frac{W_2^\gamma(q^2)}{M_p M_\Sigma},
$$

(2.18)

Here the introduction of the factors $M_p/M_\Sigma, 1/M_p M_\Sigma$, and $1/2M_p M_\Sigma$ makes the $W_{\mu\nu}^\gamma$'s dimensionless. On the other hand, $W_{\mu\nu}^i$ for $\{i = \gamma Z, Z\}$ can be decomposed into six structures:

$$
W_{\mu\nu}^i \equiv -g_{\mu\nu} \frac{M_p}{M_\Sigma} W_1^i(q^2) + P_\mu P_\nu \frac{W_2^i(q^2)}{M_p M_\Sigma} - i\epsilon_{\mu
u\rho\sigma} P^\rho q^\sigma \frac{W_3^i(q^2)}{2M_p M_\Sigma} + \mathcal{O}(q_\mu \text{ or } q_\nu).
$$

(2.19)

Here terms proportional to $q_\mu$ or $q_\nu$ vanish upon contraction with $l_{\mu\nu}^i$. Carrying out the integral in Eq. (2.17), we obtain the following expressions for the three parts of the cross section:

$$
\left( \frac{d\sigma}{d\Omega} \right)_\gamma = \left( \frac{d\sigma}{d\Omega} \right)_0 \eta^\gamma \left[ W_2^\gamma(q^2) + 2 \tan^2 (\theta/2) W_1^\gamma(q^2) \right],
$$

(2.20)

$$
\left( \frac{d\sigma}{d\Omega} \right)_{\gamma Z} = \left( \frac{d\sigma}{d\Omega} \right)_0 \eta^{\gamma Z} \left\{ 2g^e_\nu \left[ W_2^{\gamma Z}(q^2) + 2 \tan^2 (\theta/2) W_1^{\gamma Z}(q^2) \right] - 2g^e_\alpha \frac{E + E'}{M_p} \tan^2 (\theta/2) W_3^{\gamma Z}(q^2) \right\},
$$

(2.21)

$$
\left( \frac{d\sigma}{d\Omega} \right)_Z = \left( \frac{d\sigma}{d\Omega} \right)_0 \eta^Z \left\{ \left[ (g^e_\nu)^2 + (g^e_\alpha)^2 \right] \left[ W_2^Z(q^2) + 2 \tan^2 (\theta/2) W_1^Z(q^2) \right] \right.
$$

$$
- \left. 2g^e_\alpha g^e_\nu \frac{E + E'}{M_p} \tan^2 (\theta/2) W_3^Z(q^2) \right\},
$$

(2.22)

where

$$
\left( \frac{d\sigma}{d\Omega} \right)_0 = \frac{\alpha_E^2}{4E^2 \sin^4 (\theta/2)} \left\{ \frac{\cos^2 (\theta/2)}{2[1 + 2(E/M_p) \sin^2 (\theta/2)]]} \right\},
$$

(2.23)

with $\alpha_E$ the electromagnetic fine structure constant. Therefore, the cross section is determined by various $W_{k j}^i$, which, in terms of the form factors, can be expressed as:
\[ W_i^j = \frac{1}{2M_p} \left\{ |F_1^j|^2 \left[ Q^2 + (M_p + M_{\Sigma})^2 \right] + |F_2^j|^2 \left[ Q^2 + (M_p + M_{\Sigma})^2 \right] \right\}, \quad (2.24) \]

\[ W_2^j = 2 \left\{ |F_2^j|^2 + |F_3^j|^2 + Q^2 \left[ \frac{|F_2^j|^2}{(M_p + M_{\Sigma})^2} + \frac{|F_3^j|^2}{(M_p - M_{\Sigma})^2} \right] \right\}, \quad (2.25) \]

\[ W_3^j = -4 \text{Re} \left[ (F_1^j + F_2^j)(F_3^j + F_4^j)^* \right], \quad (2.26) \]

for \( \{j = \gamma, Z\} \), where \( \text{Re} \) denotes the real part. The expressions for \( W_k^{\gamma Z} \), which contain the product of \( F_k^\gamma \) and \( F_k^Z \), have not been listed.

### III. ESTIMATE OF THE CROSS SECTION

The hadron vertices \( p\gamma\Sigma^+ \) and \( pZ\Sigma^+ \) are very complicated, containing the interplay among the strong, weak, and (for \( p\gamma\Sigma^+ \)) electromagnetic interactions. Obviously, it is difficult to calculate the form factors and cross section directly from the standard model. At this stage, the best one could do is to calculate the form factors in models. For the experiment conditions of interest we find it possible to calculate the cross section without invoking any explicit effective model, by making use of experimental information on the weak radiative decay \( \Sigma^+ \rightarrow p\gamma \) and the flavor changing weak neutral current.

Let us first consider the photon-exchange part. It is easy to show that in our notation, the weak radiative decay width of the \( \Sigma^+ \) can be written as

\[ \Gamma(\Sigma^+ \rightarrow p\gamma) = \frac{e^2}{\pi} \left( \frac{M_{\Sigma}^2 - M_p^2}{2M_{\Sigma}} \right)^3 \left[ \frac{|F_2^\gamma(0)|^2}{(M_p + M_{\Sigma})^2} + \frac{|F_4^\gamma(0)|^2}{(M_p - M_{\Sigma})^2} \right], \quad (3.1) \]

and the asymmetry parameter as

\[ \alpha_{\gamma} = 2 \left( \frac{M_p + M_{\Sigma}}{M_p - M_{\Sigma}} \right) \left\{ \frac{\text{Re} \left[ F_2^\gamma(0)F_4^{\gamma*}(0) \right]}{|F_2^\gamma(0)|^2 + [(M_p + M_{\Sigma})(M_p - M_{\Sigma})] |F_4^\gamma(0)|^2} \right\}. \quad (3.2) \]

Experimentally, the branching ratio of \( \Sigma^+ \rightarrow p\gamma \) is \( (1.25 \pm 0.07) \times 10^{-3} \), and \( \alpha_{\gamma} = -0.76 \pm 0.08 [2] \). Note that \( \tan^2(\theta/2) \sim q^2 \) [Eq. (2.1)], and \( F_1^\gamma(q^2) \sim q^2, F_2^\gamma(q^2) \sim q^2 \) [Eq. (2.17)] (at most) for small \( q^2 \). To first order in \( q^2 \), we have

\[ W_k^\gamma(q^2) + 2 \tan^2 \left( \frac{\theta}{2} \right) W_1^\gamma(q^2) = -2q^2 \left[ \frac{|F_2^\gamma(0)|^2}{(M_p + M_{\Sigma})^2} + \frac{|F_4^\gamma(0)|^2}{(M_p - M_{\Sigma})^2} \right] \]

\[ + 2 \tan^2 \left( \frac{\theta}{2} \right) \left[ \frac{(M_p^2 - M_{\Sigma}^2)^2}{2M_p^2} \right] \left[ \frac{|F_2^\gamma(0)|^2}{(M_p + M_{\Sigma})^2} + \frac{|F_4^\gamma(0)|^2}{(M_p - M_{\Sigma})^2} \right] + O[(q^2)^2] \]

\[ = -(2q^2) \times 3.11 \times 10^{-14}(\text{GeV})^{-2} + 2 \tan^2 \left( \frac{\theta}{2} \right) \times 5.05 \times 10^{-15} + O[(q^2)^2]. \quad (3.3) \]

Here in the last step, we have used the experimental values for \( \Gamma(\Sigma^+ \rightarrow p\gamma) \) and the lifetime of \( \Sigma^+ \). Therefore, the above first order result is model independent. To be concrete, we consider
the kinematics of the parity violation experiment at MAMI [4], \( \theta = 35^0 \) and \( q^2 \simeq -0.237 \text{ GeV}^2 \) for the elastic scattering, which implies \( q^2 \simeq -0.16 \text{ GeV}^2 \) for the process under consideration. Using the above estimate, we arrive at the following result:

\[
\frac{\left( \frac{d\sigma}{d\Omega} \right)_\gamma (ep \to e\Sigma^+)}{\frac{d\sigma}{d\Omega} (ep \to ep)} \simeq 4 \times 10^{-15} .
\]

(3.4)

This indicates a severe suppression relative to elastic scattering. This suppression arises from the weak interaction, gauge invariance, and kinematics.

We observe from Eqs. (2.24) and (2.25) that \( W_1^\gamma \) and \( W_2^\gamma \) may increase as \( Q^2 \) gets larger, implying, perhaps, a larger value for the ratio of \( (d\sigma/d\Omega)_{ep \to e\Sigma^+} / (d\sigma/d\Omega)_{ep \to ep} \). On the other hand, as \( Q^2 \) goes to infinity, one expects all the form factors go to zero. Thus, there may be a chance that the ratio has a maximum at a non-zero momentum transfer, which, however, is unlikely to alter the result of Eq. (3.4) qualitatively.

The \( Z^0 \)-exchange part and the interference part involve the flavor changing weak neutral current, which is suppressed (at the tree level) in the standard model. At the quark level, the relevant vertex is \( sZ^0d \). For the purposes at hand, we can parameterize this vertex in terms of a vector coupling \( g^Z_V \) and an axial vector coupling \( g^Z_A \). Note that the same vertex is also responsible for the decay \( K^0_L \to \mu^+\mu^- \). Although the typical momentum scales in the two processes (\( ep \to e\Sigma^+ \) and \( K^0_L \to \mu^+\mu^- \)) differ by multiples of some typical hadronic scale, we will not make a significant error by treating \( g^Z_V \) and \( g^Z_A \) as constants. Since for electron and muon, \( g^e,\mu_v \sim 0 \), only \( g^Z_A \) enters.

Thus we can make use of the experimental information on the decay \( K^0_L \to \mu^+\mu^- \) to estimate \( g^Z_A \). Since the decay width of \( K^0_L \to \mu^+\mu^- \) agrees with the standard model estimate (via \( K^+ \to \gamma\gamma \to \mu^+\mu^- \)), we can safely assume \( g^Z_A \) contribution to \( K^0_L \) decay to be below the limit of the experimental errors on the \( K^0_L \to \mu^+\mu^- \) branching ratio. This gives:

\[
|g^Z_A| \leq \frac{1}{2} \sin \theta_c \cos \theta_w \left[ \frac{\Gamma(K^0_L \to \mu^+\mu^-)}{\Gamma(K^+ \to \mu^+\nu\mu)} \right]^{1/2} \simeq 2 \times 10^{-6} ,
\]

(3.5)

where \( \theta_c \) is the Cabibbo angle. Here we have neglected phase space difference, which is expected to be small.

With the neglect of QCD binding effects, we can apply the above estimate directly to the vertex \( pZ^0\Sigma^+ \), where only the axial vector coupling contributes at \( Q^2 \sim 0 \). We then obtain

\[
\frac{\left( \frac{d\sigma}{d\Omega} \right)_Z (ep \to e\Sigma^+)}{\frac{d\sigma}{d\Omega} (ep \to ep)} \leq 7 \times 10^{-22} .
\]

(3.6)

Therefore, the cross section for the process \( ep \to e\Sigma^+ \) is dominated by the photon exchange part. The \( Z^0 \)-exchange and the interference give negligible contributions because of the suppression of the flavor changing weak neutral current.
IV. SUMMARY AND CONCLUSION

In this paper, we have studied the kinematics and derived the cross section of the scattering process $ep \rightarrow e\Sigma^+$. The cross section can be expressed in terms of various form factors which are introduced to parameterize the physics at the vertices $p\gamma\Sigma^+$ and $pZ^0\Sigma^+$. We have estimated the cross section at small momentum transfer invoking known experimental information. In particular, we obtain a model independent first order (in momentum transfer) result for the photon-exchange part of the cross section, which is completely determined by the weak radiative decay width $\Gamma(\Sigma^+ \rightarrow p\gamma)$. With this result and the kinematics of the parity violation experiment at MAMI, we found a factor of $4 \times 10^{-15}$ suppression relative to the elastic process $ep \rightarrow ep$. This suppression largely results from the weak interaction involved and the gauge invariance at the vertex $p\gamma\Sigma^+$. The $Z^0$-exchange part and the interference part both depend on the flavor changing weak neutral current. Using the experimental result for the branching ratio of $K^0_L \rightarrow \mu^+\mu^-$, we estimate that these two parts are much smaller than the photon-exchange contribution.

In principle, the weak hyperon production process $ep \rightarrow e\Sigma^+$ discussed here can be measured experimentally in facilities with low energy electron beams. For low momentum transfer, MAMI, MIT-Bates, and TJNAF are the best candidates. However, our estimate shows that the cross section is severely suppressed relative to the $ep$ elastic scattering. To be specific, consider the parity violation experiment proposed at MAMI. There, about $10^{14}$ elastic events can be accumulated (in 700 hours) \cite{4}, making a meaningful measurement unlikely at this time.

ACKNOWLEDGMENTS

We are indebted to Frank Maas for suggesting the scattering process addressed here to us and providing useful information on the parity violation experiment at MAMI. We would like to thank Lisa Randall for useful conversations.
REFERENCES

[1] F.E. Maas, private communication.

[2] L. Montanet et al., Phys. Rev. D 50, 1173 (1994).

[3] There have been many theoretical works discussing hyperon radiative decays. For a recent review and references, see B. Bassalleck in Proceedings of the International Conference On Hypernuclear and Strange Particle Physics, [Nucl. Phys. A585, 255c (1995)].

[4] F.E. Maas et al., “The New Parity Violation Experiment at MAMI”, Proceedings of the Erice Summer school on the spin structure of the nucleon, Erice, Sicili, August 10-15, 1996.

[5] C. Itzykson and J. Zuber, “Quantum Field Theory” (McGraw-Hill, New York, 1980).