sult is as good as a similar test in fragmentation region, where a 15% agreement with factorization[10] was obtained. The theoretical picture gives the most detailed prediction near the KL. Our result indicates the validity of the approximation in Eq. (1). Future experiments of high missing mass at high energy, hopefully covering a wider range of t and s'/s, may allow us to study the problem of whether Pomeranchukons decouple and to find the t dependence of triple-Reggeon couplings systematically.

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[1]We give the following two review papers instead of putting down a long list of references: W. R. Frazer *et al.*, Rev. Mod. Phys. 44, 284 (1972); S. Gasiorowicz, Minnesota Report No. COO-1764-119, 1971 (unpublished).

[2]R. P. Feynman, Phys. Rev. Letters 23, 1415 (1969); J. Benecke *et al.*, Phys. Rev. 188, 2159 (1969).

[3]A. H. Mueller, Phys. Rev. D 2, 2963 (1970).

[4]H. D. L. Abarbanel *et al.*, Phys. Rev. Letters 26, 937 (1971).

[5]R. D. Peccei and A. Pignotti, Phys. Rev. Letters 26, 1076 (1971); M.-S. Chen, L.-L. Wang, and T. F. Wong, Phys. Rev. D 5, 1667 (1972); R. M. Edelstein *et al.*, Phys. Letters 35B, 408 (1971); J.-M. Wang and L.-L. Wang, Phys. Rev. Letters 26, 1287 (1971); J. H. Ting and H. J. Yesian, Phys. Letters 35B, 321 (1971); S. Ellis and A. I. Sanda, Phys. Rev. D 6, 1347 (1972). In the last reference, the authors constructed a simple two-term model which roughly explains the data. We are interested in a direct test of scaling and factorization.

[6]Y. A. Antipov *et al.*, Phys. Letters 40B, 147 (1972).

[7]J. V. Allaby *et al.*, data presented at Oxford Conference, 1972 (unpublished); CERN report, 1972 (unpublished).

[8]J. V. Allaby *et al.*, CERN Report No. 70-16, 1970 (unpublished).

[9]F. E. Paige and L.-L. Wang, BNL report, 1972 (unpublished).

[10]M.-S. Chen *et al.*, Phys. Rev. Letters 26, 1585 (1971).

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Method to Distinguish Between Multiparticle Production Mechanisms*

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Multiparticle production in particle-nucleus collisions at high energies is shown to provide a sensitive test for mechanisms of multiparticle production in particle-nucleon collisions at high energies.

I. INTRODUCTION

Various models for multiparticle production in high-energy collisions of elementary particles have been proposed since the discovery of cosmic rays. The accumulated data from cosmic rays and from high-energy accelerators have practically ruled out the models that had been proposed prior to the last decade, but some concepts of the original models have survived. Some aspects of models fashionable today can be traced back to old models like the Fermi statistical model[1] and the Heisenberg and bremsstrahlung models.[2] While the old models shared the advantage of being subject to simple experimental tests by which they could easily be ruled out, the present models for multiparticle production are generally too flexible. To our knowledge no crucial experimental tests have been found which could distinguish between the various models. It is our intention here to demonstrate that multiparticle production in particle-nucleus collisions at high energy may provide a sensitive test for mechanisms of multiparticle production in particle-nucleon collisions.

Models for multiparticle production in high-energy particle-nucleon collisions fall into two categories. The first category includes those models which assume that the final multiparticle state *is* *directly produced in a single step* from the colliding particles. This includes the particle-fragmentation model,[3] the multiperipheral model[4] and its multi-Regge generalizations,[5] and the bremsstrahlung analogy.[6] The second category consists of those models which assume *a two-step mechanism*, where in the first step one or two compound systems are produced which subsequently decay into the final multiparticle states with lifetimes long compared to the collision time. Such models include the diffraction excitation model[7] and the one- or two-fireball thermodynamic models.[8]

Below we will show that by measuring the dependence of the average multiplicity in particle-nucleus collisions on the incident energy and on the atomic number of the target nucleus, one can distinguish between the two classes of models. The approach here is sufficiently general that with minor adjustments the different models within each class may be accommodated. It is anticipated that with sufficiently accurate measurements it may even be possible to differentiate the models within a class. The derivation of our main formula involves some "nuclear optics" which we summarize first. Later we turn to the derivation of the average multiplicity in high-energy particle-nucleus collisions for the two classes of models.

II. ELEMENTARY NUCLEAR OPTICS

At high energy the absorption and refraction of a plane wave Ψ of incident particles of momentum k by a slab of matter of thickness dx and uniform density ρ (particles per unit volume) is described by the well-known optical formula[9]

$$d\Psi = \frac{i2\pi f(0)}{k}\rho\Psi\,dx\,, \tag{1}$$

where $f(0)$ is the forward scattering amplitude of the incident particle on a target particle, normalized such that

$$\frac{d\sigma}{d\Omega} = |f|^2\,.$$

Let us denote by β the ratio

$$\beta = \frac{\mathrm{Re}f(0)}{\mathrm{Im}f(0)}\,; \tag{2}$$

then using the optical theorem for the particle-particle collision,

$$\sigma = \frac{4\pi}{k}\,\mathrm{Im}f(0)\,,$$

Eq. (1) can be rewritten as

$$\frac{d\Psi}{dx} = -\tfrac{1}{2}\sigma(1 - i\beta)\rho\Psi\,. \tag{3}$$

Throughout we use σ to denote the total cross section of the beam particle with a target particle. Consider now a target which consists of $A = \int \rho\, d^3r$ particles and whose transverse dimensions are much larger than the incident wave length. It then follows from Eq. (3) that the cross sections for absorption and for elastic scattering and the total cross section are given in the impact-parameter representation by

$$\sigma_{\mathrm{abs}} = \int (1 - e^{-\sigma T(b)})d^2b\,, \tag{4}$$

$$\sigma_{\mathrm{el}} = \int \{1 + e^{-\sigma T(b)} - 2e^{-\sigma T(b)/2}\cos[\beta\tfrac{1}{2}\sigma T(b)]\}d^2b\,, \tag{5}$$

$$\sigma_{\mathrm{tot}} = 2\int \{1 - e^{-\sigma T(b)/2}\cos[\beta\tfrac{1}{2}\sigma T(b)]\}d^2b\,, \tag{6}$$

respectively, where the thickness function $T(b)$ is the total number of target particles encountered at impact parameter b:

$$T(\vec{b}) = \int_{-\infty}^{+\infty} \rho(\vec{b} + \hat{1}_z z)dz\,. \tag{7}$$

For high-energy collisions between elementary particles and nucleons, $\beta \ll 1$, and since most of the contribution to σ_{el} and σ_{tot} in particle-nucleus collisions comes from regions where $\sigma T(b) \gtrsim 1$, one can replace expressions (5) and (6) for medium and heavy nuclei by the approximate relations

$$\sigma_{\mathrm{el}} \cong \int (1 - e^{-\sigma T(b)/2})^2 d^2b\,, \tag{8}$$

$$\sigma_{\mathrm{tot}} \cong 2\int (1 - e^{-\sigma T(b)/2})d^2b\,. \tag{9}$$

Let us now examine multiparticle production in particle-nucleus collisions at high energies

III. INTRANUCLEUS CASCADING

Consider a situation where multiparticle final states in particle-nucleon collisions are produced in the lab frame in a time interval not longer than the collision time (typically of the order of 1 fm/c $= 0.3\times10^{-23}$ sec). If the production takes place inside the nuclear medium, and if the mean free path of the produced particles within the nuclear medium is short compared to the nuclear dimension, then an intranucleus cascading phenomena will develop: The first generation of produced particles will produce a second generation of particles via collisions while escaping from the nuclear interior; the second generation will produce a third generation, etc. The final multiparticle state will consist of all the particles from all generations that escape outside of the nuclear fragments. In principle, if all cross sections for all particle-nucleon collisions were known, one could carry out approximate numerical simulation of intranucleus cascading, using Monte-Carlo-type methods. However, since only a very small fraction of the total required information is available, and since Monte Carlo simulations are lengthy, tedious, and untransparent, we present here a simple analytic model based on simplifying assumptions, which we believe reproduces the significant features of the intranucleus cascading phenomena.

IV. INTRANUCLEUS CASCADING MODEL

The model is based on the following assumptions:

(i) The average multiplicity $\langle n(E)\rangle$ in particle-nucleon collisions increases logarithmically with E, the incident energy:

$$\langle n(E)\rangle = a\ln(E/E_0)\,, \tag{10a}$$

which has been shown[10] to hold approximately for the charge multiplicity in pp collisions between 10 and 1000 GeV/c. The average multiplicity is defined in the conventional manner

$$\langle n(E)\rangle = \sum_n n\sigma_n/\sigma\,, \tag{10b}$$

where σ_n is the cross section for the production of n particles.

(ii) The produced multiparticle final states in high-energy particle-nucleon collisions are confined *in the lab system* to a narrow forward cone along the beam axis.[11] Since nuclear dimensions are typically a few mean free paths for elementary particles, the intranucleus cascading will be confined to a narrow forward cone along the beam axis.

In addition to assumptions (i) and (ii), we also make the following simplifying assumptions:

(iii) At any stage of the cascading, on the average all the produced particles share equally the initial energy and have the same cross section, σ_0, for collisions with target nucleons. Thus all the produced particles are assumed to generate the same average multiplicity in secondary collisions.

Assumptions (i), (ii), and (iii) lead to the following differential equation for intranucleus particle multiplication along impact parameter b:

$$dN = N[\langle n(E/N)\rangle - 1]\sigma_0\rho\, dz\,, \tag{11}$$

where both N and ρ are functions of b and z, and where

$$\begin{aligned}\langle n(E/N)\rangle &= a\ln(E/NE_0)\\ &= \langle n(E)\rangle - a\ln N\,.\end{aligned} \tag{12}$$

The solution to Eq. (11) with the boundary condition $N(E,b,z=-\infty)=1$, i.e., the multiplication factor at $z=-\infty$ is equal to 1, is given by

$$N(E,b,z) = \exp\left\{\frac{\langle n(E)\rangle - 1}{a}[1-e^{-a\sigma_0 T(b,z)}]\right\}, \tag{13}$$

where $T(b,z)$ is the number of target nucleons encountered at impact parameter b between $-\infty$ and z,

$$T(b,z) = \int_{-\infty}^{z}\rho(b,z')dz'\,. \tag{14}$$

The total multiplication factor for impact parameter b is therefore given by

$$\begin{aligned}N(E,b) &\equiv N(E,b,z=\infty)\\ &= \exp\left\{\frac{\langle n(E)\rangle - 1}{a}[1-e^{-a\sigma_0 T(b)}]\right\}\end{aligned} \tag{15a}$$

or by

$$N(E,b) = \left(\frac{E}{E^*}\right)^{\{1-\exp[-a\sigma_0 T(b)]\}}, \tag{15b}$$

where

$$E^* = E_0 e^{1/a}\,.$$

The multiplicity for a particle-nucleus collision can now be calculated by analogy with (10b) by using

$$\langle N(E)\rangle = \frac{\int N(E,b)\sigma_{\text{tot}}(b)}{\int \sigma_{\text{tot}}(b)}\,, \tag{16}$$

where from Eq. (9) we read

$$\sigma_{\text{tot}}(b) = 2(1-e^{-\sigma T(b)/2})2\pi b\, db\,. \tag{17}$$

For heavy nuclei and small impact parameters, Eq. (15) has then an approximate simple form:

$$N(E,b) \cong E/E^*\,. \tag{18}$$

Equation (18) can be understood easily. $E^* = e^{1/a}E_0$ is the average threshold energy for producing a single particle per collision, as can be deduced from Eq. (10a). E/E^* is thus the maximum number of final particles that can be produced if the incident energy is E and if there are sufficiently many collisions to convert the energy into particles, i.e., if $\sigma_0 T(b) \gg 1$.

Equation (15) can be easily generalized to situations where the final multiparticle states in particle-nucleon collisions are actually produced via intermediate states composed of resonances and stable particles. If the average particle multiplicity of such intermediate states is $\alpha\langle n(E)\rangle$, then the nuclear multiplication factor is given by the following expression:

$$\begin{aligned}N(E,b) &= \frac{1}{\alpha}\exp\left\{\frac{\alpha\langle n(E)\rangle - 1}{a^*}[1-e^{-a^*\sigma_0 T(b)}]\right\}\\ &= \frac{1}{\alpha}\left(\frac{E}{e^{1/a^*}E_0}\right)^{\{1-\exp[-a^*\sigma_0 T(b)]\}},\end{aligned} \tag{19}$$

where

$$a^* = \alpha a ,$$

and where the over-all factor $1/\alpha$ arises from the decay of the resonances outside the nucleus due to time-dilatation effects at high energy. As expected, Eqs. (15) and (19) show that the intranucleus cascading phenomena strongly depend on the incident energy and on the atomic number of the target nucleus. In particular it transforms the dependence of the average particle multiplicity on the incident energy from a logarithmic dependence for particle-nucleon collisions into a linear dependence for particle-heavy nucleus collisions.

V. INTRANUCLEUS FORMATION AND TRANSPORT OF RESONANCES

We will now consider multiparticle production in nuclei under the assumption that multiparticle production in particle-nucleon collisions proceed mainly via intermediate states of one or two compound systems which subsequently decay into the final multiparticle states. To be more specific let us consider the diffractive-excitation models.[7] According to these models, one or both of the colliding particles are first excited into resonant states which have the same internal quantum numbers, excepting spin and parity, and approximately the same momenta as the original particles. These resonances subsequently decay via cascading (mainly via pion emission) with small Q values (typically 300 MeV) into the final multiparticle states. Since the average multiplicity increases very slowly with energy, and since the average Q value for the cascading decay of the resonances is also small, the important resonances are therefore of low-mass excitation and one may expect them to have lifetimes not much different from measured lifetimes of familiar low excitation states, say $\Gamma_{\text{average}} \sim 100$ MeV. Consequently at high energy a beam particle which is diffractively excited inside the nucleus behaves there like a stable particle due to time dilatation. On the other hand, resonances which are produced within the nucleus by diffractive excitation of target nucleons have only small momenta in the lab frame. They will start their cascading decay inside the nucleus; however, due to the small Q values in their cascading decays, the decay products will not have sufficient energy to multiply via secondary collisions while escaping out of the nucleus.

The average multiplicity in particle-nucleus collisions for the diffractive excitation mechanism can now be easily evaluated: Let us assume that the average multiplicity $\langle n(E)\rangle$ in the particle-nucleon collision is made of an average multiplicity $\langle n_B(E)\rangle$ associated with the decay of the excited states of the beam particle, and of average multiplicity $\langle n_n(E)\rangle$ associated with the decay of the excited states of the target nucleon,

$$\langle n(E)\rangle = \langle n_n(E)\rangle + \langle n_B(E)\rangle . \tag{20}$$

We also assume that the beam particle has the same cross sections in its ground and excited states, it will produce $\sigma T(b)$ nucleon resonances when it traverses the nucleus at impact parameter b, and after it emerges behind the nucleus it will decay on the average into $\langle n_B(E)\rangle$ particles. The beam particle is assumed to retain approximately its initial energy E throughout its collisions with target nucleons. The resonances of the target nucleons themselves thus have insufficient energy to create more resonances by subsequent collisions. The multiplication factor at impact parameter b therefore is given by

$$N(E,b) \cong \langle n_B(E)\rangle + \sigma T(b)\langle n_n(E)\rangle , \tag{21}$$

and the average multiplicity in particle-nucleus collision is given by

$$\langle N(E)\rangle = \frac{\int[\langle n_B(E)\rangle + \sigma T(b)\langle n_n(E)\rangle]\sigma_{\text{tot}}(b)}{\int \sigma_{\text{tot}}(b)} , \tag{22}$$

where again

$$\sigma_{\text{tot}}(b) = 2(1 - e^{-\sigma T(b)/2}) 2\pi b\, db . \tag{23}$$

For nucleon-nucleus collisions

$$\langle n_n(E)\rangle = \langle n_B(E)\rangle = \tfrac{1}{2}\langle n(E)\rangle \tag{24}$$

and Eq. (22) for heavy nuclei can be well approximated by

$$\frac{\langle N(E)\rangle}{\langle n(E)\rangle} = \left(\frac{1}{2} + \frac{A\sigma}{\sigma_{\text{tot}}}\right) . \tag{25}$$

Contrary to intranucleus cascading, which is associated with fragmentation or multiperipheral-type mechanisms, the diffractive excitation mechanism (and the two-fireball mechanism) leads to average multiplicity in particle-nucleus collisions which has the same energy dependence as the average multiplicity in particle-nucleon collisions. In addition, since σ_{tot} behaves as $\sigma A^{2/3}$, the diffractive excitation mechanism shows only a weak dependence on the atomic number of the target nucleus.

VI. RESULTS

Figure 1 presents numerical estimates based on the intranucleus cascading model for the average particle multiplicity in proton-nucleus collisions for different combinations of incident energy and target nucleus. The calculations are based on Eqs. (15)–(17) with the following input information:

(i) The high-energy proton-nucleon total cross section was taken to be 38.5 mb.

(ii) The cross section of the produced particles

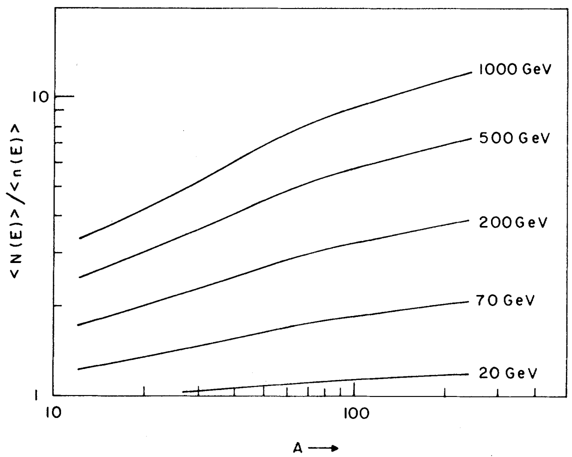

FIG. 1. Ratio of the average particle multiplicity in proton-nucleus collisions $\langle N(E)\rangle$, based on the intranucleus cascading model as described in the text, to the measured average multiplicity in pp collisions $\langle n(E)\rangle$ as a function of incident energy and atomic number of the target nucleus.

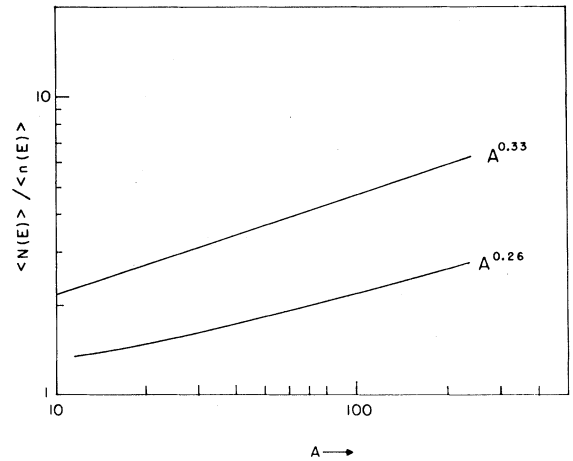

FIG. 2. Ratio of the average particle multiplicity in proton-nucleus collisions $\langle N(E)\rangle$, based on the diffractive excitation model as described in the text, to the measured average multiplicity in pp collisions $\langle n(E)\rangle$ as a function of atomic number of the target nucleus.

with target nucleons was taken to be 24 mb, the high-energy π-nucleon cross section.

(iii) The average multiplicity in particle-nucleon collision was taken as $\frac{3}{2}$ the observed[10,12] multiplicity of charged particles in pp collisions.

$$\langle n(E)\rangle \cong \tfrac{3}{2}\langle n_c(E)\rangle \cong 0.63\ln\left(\frac{E}{1.71}\right), \qquad (26)$$

where E is in GeV. The factor $\frac{3}{2}$ was introduced in order to account for unobserved neutral particles in the final states. It is based on the assumptions of charge independence and that most of the particles in the final multiparticle states are pions.

(iv) The nuclear density distributions were assumed to be Gaussian with the rms radius given by the nuclear charge distributions as deduced by Hofstadter[13] from electron-nucleus scattering.

Figure 1 clearly demonstrates the strong dependence of the average multiplicity on the atomic number of the target nucleus and on the incident energy. In particular note the large multiplication predicted for heavy nuclei and high energies, where we believe the analysis to be most trustworthy. Figure 2 presents numerical estimates based on the diffractive excitation mechanism for the average particle multiplicity in p-nucleus collisions as given by formula (22) with $\sigma = 38$ mb and with σ_{tot} calculated via the integration of (23). The calculated σ_{tot} was found to be in good agreement with the measurements of Belletini *et al.*[14] at incident proton momentum of 20 GeV/c. This agreement is desirable since σ is energy-independent from 20 to 1000 GeV, and by (23) σ_{tot} has the same property. Figure 2 clearly indicates the weak dependence of the average multiplicity on the atomic number of the target nucleus typical of the diffractive excitation model. In particular, for heavy nuclei the dependence is approximately described by an $A^{0.26}$ dependence law.

VII. CONCLUSIONS

The present paper clearly demonstrates that the dependence of the average particle multiplicity in particle-nucleus collisions on the incident energy and on the atomic number of the target nucleus provides a strong indication of the dominant mechanism for multiparticle production in high-energy collisions of elementary particles: For the case of one-step mechanisms, the dependence of the average multiplicity on the incident energy is changing drastically from a logarithmic dependence for proton and light nuclei into a linear dependence for heavy nuclei. However, in addition, the average multiplicity is approximately proportional to some power of A which increases with energy. For the two-step mechanisms, the average multiplicity has the same logarithmic energy dependence for all nuclei and it increases with atomic number as $A^{0.26}$. These qualitative conclusions are shown not to be sensitive to the specific details of the production mechanisms, since this formalism is easily adjusted to accommodate the various models in each class. A study of average multiplicity in high-energy particle-nucleus collisions as a function of incident energy and atomic number will provide a simple tool for distinguishing between mechanisms for multiparticle production in elementary-particle collisions.

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[1]E. Fermi, Phys. Rev. 81, 683 (1951).

[2]W. Heisenberg, Z. Physik 133, 65 (1952).

[3]J. Benecke, T. T. Chou, C. N. Yang, and E. Yen, Phys. Rev. 188, 2159 (1969); T. T. Chou and C. N. Yang, Phys. Rev. Letters 25, 1072 (1970).

[4]D. Amati, A. Stanghellini, and S. Fubini, Nuovo Cimento 26, 896 (1962).

[5]F. Zachariasen and G. Zweig, Phys. Rev. 160, 1322 (1967); 160, 1326 (1967).

[6]R. P. Feynman, Phys. Rev. Letters 23, 1415 (1969); in *High Energy Collisions*, edited by C. N. Yang *et al.* (Gordon and Breach, New York, 1969).

[7]R. K. Adair, Phys. Rev. D 5, 1105 (1972); M. Jacob and R. Slansky, Phys. Letters 37B, 408 (1971).

[8]R. Hagedorn, Suppl. Nuovo Cimento 3, 147 (1965); R. Hagedorn and J. Ranft, *ibid.* 6, 169 (1968).

[9]See for instance G. Molière, Z. Naturforsch. 2A, 133 (1947).

[10]See for instance L. Jones *et al.*, Phys. Rev. Letters 25, 1779 (1970); O. Czyzewski and K. Rybicki, INR, Cracow Report No. 215/70-RC/344/59, 1970 (unpublished).

[11]D. E. Lyon, Jr., C. Risk, and D. Tow, Phys. Rev. D 3, 104 (1971).

[12]G. R. Charlton and G. H. Thomas, Argonne National Laboratory Report No. ANL/HEP 7217 (unpublished).

[13]R. Hofstadter, Rev. Mod. Phys. 28, 214 (1956).

[14]G. Belletini, G. Cocconi, A. N. Diddens, E. Lillethun, G. Matthiae, J. P. Scanlon, and A. M. Wetherell, Nucl. Phys. 79, 609 (1966).

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Decay $\Sigma^{\pm} \rightarrow \Lambda e^{\pm} \nu$ *

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The ratio of the vector to axial-vector coupling constant for $\Sigma^{\pm} \rightarrow \Lambda e^{\pm}\nu$ decays using 186 events is determined to be -0.37 ± 0.20. The branching ratio for $\Sigma^{-} \rightarrow \Lambda e^{-}\nu$ is $(0.62 \pm 0.07) \times 10^{-4}$ and for $\Sigma^{+} \rightarrow \Lambda e^{+}\nu$ is $(0.21 \pm 0.05) \times 10^{-4}$. An upper limit on the magnitude of the ratio of the axial-magnetism to axial-vector coupling constants is 3.2.

Recently, three comparable experiments were performed[1-3] to study the strangeness-conserving decays $\Sigma^{\pm} \rightarrow \Lambda e^{\pm}\nu$. The conserved-vector-current hypothesis (CVC) relates the vector coupling in these decays to the amplitude for $\Sigma^0 \rightarrow \Lambda^0 \gamma$; it is thus expected that, to zero order in the momentum transfer, the vector coupling should not contribute to such decays.[4]

The sample of events available in each experiment was rather small. We therefore felt it worthwhile to repeat the analysis of the combined sample rather than simply averaging the results presented by each group. The combined sample contains 163 examples of Σ^- decays and 23 examples of Σ^+ decays. The present analysis has also been performed under slightly more general assumptions than before.

The decays $\Sigma^{\pm} \rightarrow \Lambda e^{\pm}\nu$ are particularly rich in information because the three-body final state can be completely kinematically analyzed by observing