Σ^0 Lifetime and Associated Photoproduction Poles*

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The Born-approximation calculations for the differential cross sections for the photoproduction of $K^+-\Lambda$ and $K^+-\Sigma^0$ pairs from hydrogen are generalized to include the effect of the $\Sigma-\Lambda$ mass difference and the possibility of odd $\Sigma-\Lambda$ parity. The residues of the poles in the unphysical regions of the center-of-mass production angle are computed. It is shown that if the laboratory photon energy is greater than about 2 Bev, one could obtain information about the Σ^0 lifetime and the anomalous moment of the produced hyperon by extrapolating the experimentally measured differential cross section in the backward direction of the K-production angle.

I. INTRODUCTION

THE lifetime of the Σ^0 particle is so short that it cannot be measured directly by present techniques. The rapid $\Sigma^0 \rightarrow \Lambda+\gamma$ decay also prevents the direct measurement of the Σ^0 static magnetic moment. This is unfortunate, since a knowledge of the Σ^0 lifetime and static moment would serve as a check to the various models of the interactions among Σ and Λ particles and photons. It is likely that our first knowledge of the Σ^0 decay amplitude λ and moment μ_Σ will result from the method of "polology," since the residues of the poles in the amplitudes for several well-known processes are proportional to λ and μ_Σ.[1] In this paper we are concerned with the poles in the two photoproduction processes, $\gamma+p\rightarrow K^++\Lambda$ and $\gamma+p\rightarrow K^++\Sigma^0$. The analytically continued amplitude for either of these processes at constant center-of-mass energy contains poles in the "backward" unphysical region of the K^+ production angle, the residues of which involve the Σ^0 decay amplitude and the static magnetic moment of the produced hyperon.[2]

The residues of the poles in the cross sections for associated photoproduction may be determined from the lowest order of the Born approximation. The Born approximation expressions given in the literature by Kawaguchi and Moravcsik[3] and by others[4,5] are incomplete for two reasons: (i) The possibility that the intrinsic $\Sigma-\Lambda$ parity is odd is not considered. This parity is crucial in those processes involving the $\Sigma-\Lambda$ vertex. (ii) The $\Sigma-\Lambda$ mass difference is neglected. Although this mass difference is small, its neglect leads to the vanishing of certain poles in the angular distributions, so that one must include the mass difference in order to obtain a complete understanding of the nature of the poles. The complete Born approximation expressions are given in Sec. III of this paper, while the possible experimental determination of the residues of poles (and hence of λ and μ_Σ) is discussed in Sec. IV.

The residues of the poles in the region $\cos\theta_K<-1$ depend not only on the hyperon moments, but also on the magnitudes of the $K\Lambda N$ and $K\Sigma N$ coupling constants G_Λ and G_Σ. In fact, the residues are proportional to G^2. These coupling constants, as well as the particle parities, may be determined from other experiments, so that extrapolation to backward poles in associated photoproduction may be considered as a method of measuring the hyperon moments and the Σ^0 lifetime.

II. HYPERON FORM FACTORS

Feldman and Fulton have derived the general form of the electromagnetic form factor for the $\Sigma-\Lambda$ transition, under both assumptions concerning the $\Sigma-\Lambda$ parity.[6] The "magnetic form-factor term" in their expressions is proportional to $f_2(k^2)\sigma_{\mu\nu}k_\nu$ if the $\Sigma-\Lambda$ parity is even, or to $f_2(k^2)\gamma_5\sigma_{\mu\nu}k_\nu$ if the $\Sigma-\Lambda$ parity is odd, where k is the four-momentum transfer. Only this f_2 term contributes on the mass shell, i.e., for real photon emission.[6] If we assume that the process $\Sigma^0\rightarrow\Lambda+\gamma$ is the only process contributing to Σ^0 decay, the Σ^0 lifetime depends on the form factor on the mass shell, and is given (in both $\Sigma-\Lambda$ parity cases) by the expression

$$\frac{1}{\tau}=\frac{|f_2(0)|^2}{\pi}\frac{k_0(M_\Sigma-M_\Lambda)^2}{M_\Sigma^2}, \tag{1}$$

where $k_0\approx M_\Sigma-M_\Lambda$ is the energy of the photon in the rest system of the Σ, and M_i is the mass of particle i. The constants $\hbar$ and c are set equal to 1 throughout this paper.

We define the "$\Sigma-\Lambda$ transition magnetic-moment" (denoted by μ_T) in terms of the value of f_2 on the mass shell, i.e.,

$$f_2(0)=\tfrac{1}{2}e\mu_T(M_\Sigma+M_\Lambda)/(2M_\Lambda),$$

where $e^2/(4\pi)=1/137$. The factor $(M_\Sigma+M_\Lambda)/(2M_\Lambda)$ has been included for convenience, so that all hyperon magnetic moments are defined in the same units, Λ magnetons.

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[1] See, for example, W. S. C. Williams, Nuovo cimento **19**, 1278 (1961).

[2] John G. Taylor, Phys. Rev. **116**, 768 (1959). This reference gives a brief discussion of the problem of extrapolating experimental data to such poles.

[3] M. Kawaguchi and M. J. Moravcsik, Phys. Rev. **107**, 563 (1957).

[4] A. Fujii and R. E. Marshak, Phys. Rev. **107**, 570 (1957).

[5] Richard H. Capps, Phys. Rev. **114**, 920 (1959).

[6] G. Feldman and T. Fulton, Nuclear Phys. **8**, 106 (1958); See Eqs. (13) and (14).

The form factor for the Λ hyperon is also given by the expressions of Feldman and Fulton,[6] if M_Σ is set equal to M_Λ. In this case the value of $2f_2/e$ on the mass shell is equal to μ_Λ, the static magnetic moment of the Λ in Λ magnetons. The static moment of the Σ^0 is defined in a similar fashion.

Since the poles in the processes $\gamma+p\rightarrow K^+ +\Lambda$ and $\gamma+p\rightarrow K^+ +\Sigma^0$ occur when the intermediate particle in the lowest order Born approximation is on the mass shell, the dependence of the residues of these poles on the $\Sigma-\Sigma$, $\Lambda-\Lambda$, and $\Sigma-\Lambda$ form factors may be expressed in terms of the magnetic moments μ_Σ, μ_Λ, and μ_T.

III. BORN APPROXIMATION FOR ASSOCIATED PHOTOPRODUCTION

Several different Feynmann diagrams contribute to the Born approximation for associated photoproduction.[7] The nature of these Born approximation amplitudes has been discussed by several authors.[3,8] We list below only the formulas for the differential cross section in the center-of-mass system, since this is the quantity most easily measured. The expressions are given only for the process $\gamma+p\rightarrow K^+ +\Lambda$; the corresponding expressions for $K^+ +\Sigma^0$ production may be obtained by obvious substitutions. The differential cross section for $K^+ +\Lambda$ production may be written in the form,

$$d\sigma/d\Omega=A_1+A_2\mu_N+A_3\mu_\Lambda+A_4\mu_N^2+A_5\mu_\Lambda^2+A^6\mu_N\mu_\Lambda +B_1\mu_T+B_2\mu_T\mu_N+B_3\mu_T\mu_\Lambda+B_4\mu_T^2, \quad (2)$$

where μ_N is the proton static, anomalous magnetic moment in nuclear magnetons, and μ_Λ and μ_T are the Λ and $\Sigma-\Lambda$ transition moments, defined as in Sec. II. The coefficients A_i and B_i are independent of μ_N, μ_Λ, and μ_T, and are functions of energy and production angle. The A_i are given by Kawaguchi and Moravcsik.[3] The formulas for the B_i are given below for the case in which the $K\Lambda N$ parity and $\Sigma-\Lambda$ parity are even. The corresponding equations for the other possible $K\Lambda N$ and $\Sigma-\Lambda$ parity cases may be obtained by applying the following rules:

Rule 1. If the $K\Lambda N$ parity is odd, replace M_N by $-M_N$ in all expressions.

Rule 2. If the $\Sigma-\Lambda$ parity is odd, replace M_Σ by $-M_\Sigma$ in all expressions.[9]

$$\begin{aligned}B_1\mu_T=&-\left(\frac{G_\Lambda G_\Sigma}{4\pi}\frac{\mu_T}{M_\Lambda}\right)\left(\frac{C}{kD_\Sigma D_K}\right)[2q^3kM_N\cos^3\theta\\&+q^2(W^2M_\Lambda-M_N^2M_\Sigma\\&+4kE_\Lambda M_N-2kE_KM_N)\cos^2\theta\\&+2qk(WE_\Lambda M_\Sigma-WE_KM_\Sigma\\&+E_\Lambda^2M_N-2E_\Lambda E_KM_N)\cos\theta\\&+q^2W^2(M_\Sigma-M_\Lambda)\\&-2kE_\Lambda E_K(WM_\Sigma+E_\Lambda M_N)],\end{aligned} \quad (3a)$$

$$\begin{aligned}B_2\mu_T\mu_N=&-\left(\frac{G_\Lambda G_\Sigma}{4\pi}\frac{\mu_T\mu_N}{M_\Lambda M_N}\right)\left(\frac{Cq}{D_\Sigma}\right)[q^2M_N^2\cos^2\theta\\&+qM_N(WM_\Sigma+WM_\Lambda+2E_\Lambda M_N)\cos\theta\\&+(WM_\Sigma+E_\Lambda M_N)(WM_\Lambda+E_\Lambda M_N)],\end{aligned} \quad (3b)$$

$$\begin{aligned}B_3\mu_T\mu_\Lambda=&\left(\frac{G_\Lambda G_\Sigma}{4\pi}\frac{\mu_T\mu_\Lambda}{M_\Lambda^2}\right)\left(\frac{C}{D_\Sigma D_\Lambda}\right)\{-2q^3kW\cos^3\theta\\&+q^2[-W^2E_\Lambda+3E_\Lambda M_N^2\\&+WM_N(M_\Sigma+M_\Lambda)]\cos^2\theta\\&+q[W^2E_\Lambda^2+3E_\Lambda^2M_N^2+2WE_\Lambda M_N(M_\Sigma+M_\Lambda)\\&+W^2M_\Lambda(M_\Sigma-M_\Lambda)]\cos\theta+WE_\Lambda[2E_NE_\Lambda^2\\&+E_\Lambda M_N(M_\Sigma+M_\Lambda)+WM_\Lambda(M_\Sigma-M_\Lambda)]\},\end{aligned} \quad (3c)$$

$$\begin{aligned}B_4\mu_T^2=&\frac{1}{2}\left(\frac{G_\Sigma^2}{4\pi}\frac{\mu_T^2}{M_\Lambda^2}\right)\left(\frac{C}{D_\Sigma^2}\right)\{-2q^3kW\cos^3\theta\\&+q^2(-W^2E_\Lambda+3E_\Lambda M_N^2+2WM_NM_\Sigma)\cos^2\theta\\&+q[W^2E_\Lambda^2+3E_\Lambda^2M_N^2+4WE_\Lambda M_NM_\Sigma\\&+W^2(M_\Sigma^2-M_\Lambda^2)]\cos\theta+WE_\Lambda[2E_NE_\Lambda^2\\&+2E_\Lambda M_NM_\Sigma+W(M_\Sigma^2-M_\Lambda^2)]\},\end{aligned} \quad (3d)$$

where the function C is given by the expression

$$C=\frac{1}{8}\frac{e^2}{4\pi}\frac{1}{qkW^3}, \quad (4)$$

and the denominator functions D_i are given by

$$D_K=\cos\theta-(E_K/q), \quad (5a)$$

$$D_\Lambda=\cos\theta+(E_\Lambda/q), \quad (5b)$$

$$D_\Sigma=\cos\theta+(E_\Lambda/q)+(M_\Sigma^2-M_\Lambda^2)/(2qk). \quad (5c)$$

All quantities refer to the center-of-mass system. The angle θ is the angle between the photon direction and K-particle direction, and k and q are the magnitudes of the momenta of the initial particles and final particles, respectively. The symbols M_i and E_i denote the mass and energy of the particle i, and $W=E_\Lambda+E_K=k+E_N$ is the total energy. The coupling constants $G_\Lambda/(4\pi)^{\frac{1}{2}}$ and $G_\Sigma/(4\pi)^{\frac{1}{2}}$ refer to the $K\Lambda N$ and $K\Sigma N$ interactions, defined as in Kawaguchi and Moravcsik,[3] while $e^2/4\pi=1/137$. The particle energies and momenta are all functions of W and the masses, of course. ($2WE_\Lambda=W^2+M_\Lambda^2-M_K^2$, etc.)

The author has pointed out previously that the Born-approximation expressions for associated photoproduction are very simple at the threshold energies.[5] If one takes into account the $\Sigma-\Lambda$ mass difference and the possible $\Sigma-\Lambda$ parity difference, the results of reference 5 must be modified. The properly modified threshold value of the differential cross section in the case of even

[7] See reference 4, Fig. 1.

[8] Fayyazuddin, Phys. Rev. **123**, 1882 (1961).

[9] In order for this rule to be consistent with the definition of the G_Λ and G_Σ given in reference 3, the sign of G_Σ must also be changed if the $K\Lambda N$ parity and $\Sigma-\Lambda$ parity both are odd.

$K\Lambda N$ and $\Sigma-\Lambda$ parity is

$$\frac{d\sigma}{d\Omega}=\tfrac{1}{2}q\frac{e^2G_\Lambda{}^2}{(4\pi)^2}\frac{M_\Lambda}{W_0{}^2(W_0{}^2-M_N{}^2)}\times\left[1+\left(\frac{\mu_N}{2M_N}\right)(W_0+M_N)-\left(\frac{\mu_\Lambda}{2M_\Lambda}\right)(W_0+M_N)-\left(\frac{\mu_T}{2M_\Lambda}\right)\left(\frac{G_\Sigma}{G_\Lambda}\right)\left(\frac{W_0M_\Sigma+M_\Lambda M_N}{M_\Lambda+\delta}\right)\right]^2, \quad (6)$$

where $W_0=M_K+M_\Lambda$, $\delta=W_0(M_\Sigma{}^2-M_\Lambda{}^2)/(W_0{}^2-M_N{}^2)$, and the anomalous moments are defined as in Sec. II of this paper, (i.e., measured in Λ magnetons). Again the corresponding equations for the other parity cases may be obtained from the rules preceding Eq. (3a).

IV. THE POLES

Throughout this section all listed formulas apply to the case of even $K\Lambda N$ and $\Sigma-\Lambda$ parities. The corresponding expressions for the other three parity cases may be obtained from the rules preceding Eq. (3a).

The differential cross section for the process $\gamma+p\to K^++\Lambda$ at fixed energy contains poles at three unphysical values of the center-of-mass production angle, i.e., at $\cos\theta=E_K/q$, $\cos\theta=-E_\Lambda/q$, and $\cos\theta=-(E_\Lambda/q)-(M_\Sigma{}^2-M_\Lambda{}^2)/(2qk)$. The poles correspond to K-particle, Λ-particle, and Σ-particle intermediate states, respectively. The second-order pole in the forward direction has been discussed by Moravcsik and by Sakurai.[10] In the backward direction there is a second-order pole only if the $\Sigma-\Lambda$ mass difference is included. This pole occurs at $-(E_\Lambda/q)-(M_\Sigma{}^2-M_\Lambda{}^2)/2qk$; the limiting value of the differential cross section at the pole may be determined from Eq. (3d) and is given by the expression,

$$D_\Sigma{}^2\frac{d\sigma}{d\Omega}(\text{at }D_\Sigma=0)=\frac{1}{64}\frac{e^2G_\Sigma{}^2}{(4\pi)^2}\frac{1}{qk^3W^2}\left(\frac{\mu_T}{M_\Lambda}\right)^2\times(M_\Sigma{}^2-M_\Lambda{}^2)^2[(M_N+M_\Sigma)^2-M_K{}^2], \quad (7)$$

where D_Σ is defined in Eq. (5c).

The energy dependence of this expression is the same as that of the second-order pole in the forward direction.[10] Both of these poles are difficult to detect at extremely high energies, because of the E^{-6} nature of the energy dependence of $D^2d\sigma/d\Omega$. Since the coefficient of the $D_\Sigma{}^2$ pole is proportional to the square of the small $\Sigma-\Lambda$ mass difference, this pole would be difficult to detect at all energies unless the $\Sigma-\Lambda$ transition moment is much larger than the nucleon anomalous moment. Hence we will neglect the backward-direction second-order pole in the rest of the discussion.

If terms of second order in $(M_\Sigma{}^2-M_\Lambda{}^2)$ are neglected, there are only first-order poles in the backward direction. The residue of the pole at $D_\Sigma=0$, determined from Eqs. (3), is [to first order in $(M_\Sigma{}^2-M_\Lambda{}^2)$],

$$D_\Sigma\frac{d\sigma}{d\Omega}(\text{at }D_\Sigma=0)=\frac{1}{8}\frac{e^2G_\Lambda G_\Sigma}{(4\pi)^2}\frac{1}{kW^3}\left(\frac{\mu_T}{M_\Lambda}\right)\left\{-\frac{2W^2M_\Lambda{}^2(M_\Sigma-M_\Lambda)}{W^2-M_N{}^2+M_\Sigma{}^2-M_\Lambda{}^2}-\frac{\mu_N}{M_N}W^2M_\Lambda M_\Sigma+\frac{\mu_\Lambda}{M_\Lambda}W^2M_\Lambda(M_\Sigma-M_\Lambda)+\frac{(M_\Sigma{}^2-M_\Lambda{}^2)}{2k}\left[\frac{\mu_N}{M_N}WM_N(M_\Sigma+M_\Lambda)-\frac{\mu_\Lambda}{M_\Lambda}W(W^2+M_\Sigma{}^2+2M_NM_\Sigma-M_K{}^2)+\frac{2W^3M_\Lambda-2WM_N{}^2M_\Sigma-2W(M_\Sigma-M_\Lambda)(M_\Lambda{}^2-M_K{}^2)}{W^2-M_N{}^2+M_\Sigma{}^2-M_\Lambda{}^2}\right]\right\}. \quad (8)$$

There is also a simple pole in the angular distribution at the angle $\cos\theta=-(E_\Lambda/q)$, arising from the interference between the μ_Λ and the μ_N terms in the angular distribution. The residue of this pole may be obtained from reference 3 or by making the replacements $M_\Sigma\to M_\Lambda$ and $G_\Sigma\to G_\Lambda$ in Eq. (8).

The part of the cross section that is linear in μ_T contains a simple pole at $\cos\theta=E_K/q$, in addition to the backward-direction poles. The residue of this pole is

$$D_K\frac{d\sigma}{d\Omega}(\text{at }D_K=0,\ \mu_T\text{ term})=-\frac{1}{4}\frac{e^2G_\Lambda G_\Sigma}{(4\pi)^2}\frac{\mu_T}{M_\Lambda}\frac{M_K{}^2(M_\Sigma-M_\Lambda)}{kW(W^2-M_N{}^2+M_\Sigma{}^2-M_\Lambda{}^2)}. \quad (9)$$

If the $\Sigma-\Lambda$ parity is even, this pole is expected to be hidden by the second-order pole discussed by Moravcsik.[10] If the $\Sigma-\Lambda$ parity is odd, however, so that the quantity $(M_\Sigma-M_\Lambda)$ in Eq. (9) becomes $(-M_\Sigma-M_\Lambda)$, this simple pole may be important. Theoretically, if one follows the procedure of Moravcsik and multiplies the differential cross section by $D_K{}^2$, the effect of the simple pole vanishes at the pole. However, if the $\Sigma-\Lambda$ and $K\Lambda N$ parities are odd, and if $\mu_T\gtrsim1$, the effect of the simple pole on the differential cross section at angles close to the forward direction may be greater than the corresponding effect of the second-order pole, so that extreme care must be taken in making the extrapolation.

[10] M. J. Moravcsik, Phys. Rev. Letters **2**, 352 (1959); J. J Sakurai, Nuovo cimento **20**, 1212 (1961).

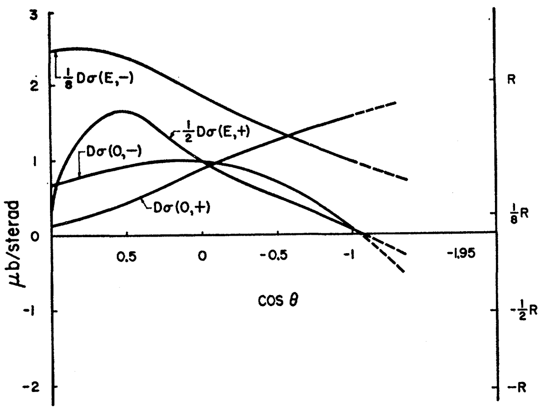

FIG. 1. Born approximation expressions for $D\sigma$, where $D=\cos\theta+E_\Lambda/q$. The first index in the expression $\sigma(a,b)$ denotes whether the parity is even or odd, while the second index denotes whether μ_Λ/M_Λ is equal to $\mu_N/|M_N|$ or $(-\mu_N/|M_N|)$. The moment μ_N is taken to be 1.9 and $G_\Lambda^2/(4\pi)$ is set equal to 2. The residue R is 2.05×10^{-30} cm². The extrapolated values of $D\sigma$ should go through R at $\cos\theta=-1.95$ in the cases $(E,-)$ and $(0,+)$, and should go through $(-R)$ in the other two cases.

At energies just above threshold, the backward-direction poles are so far from the physical region that it is not feasible to attempt to determine their residues by an extrapolation procedure. On the other hand, the terms quadratic in the magnetic moments dominate the Born-approximation cross section at very high energies (especially at extreme back angles), and the effect of the D_Σ and D_Λ poles may be large. Since the energy dependence of the leading terms in the residues of the backward poles (the $\mu_T\mu_N$, $\mu_\Lambda\mu_N$, and $\mu_T\mu_\Lambda$ terms) is of the type $(Wk)^{-1}$, the effect of these poles does not vanish as the energy becomes very high; i.e., the residue divided by the distance of the pole from the physical region approaches a constant at high energies. Hence it is not likely that these poles will be swamped by the non-Born contributions (dispersion integral corrections) at high energies. Of course, as is always the case in such extrapolations, one has to worry about other nearby singularities. For example, there are branch cuts, related to $\pi+\Lambda$ and $\pi+\Sigma$ intermediate states, that are not much farther from the physical region than the Λ- and Σ-particle poles; these cuts approach the physical region with the same energy dependence as the poles at high energies. The anomalous nucleon moment is quite large, however; if either μ_T or μ_Λ is of magnitude comparable to μ_N, it seems likely that pole effects will be large enough to be separable from the branch cut effects.

If the $\Sigma-\Lambda$ parity is even, the most important back-direction pole terms are the term proportional to $\mu_T\mu_N$ and the similar term involving $\mu_\Lambda\mu_N$, since these terms are the only ones that are not proportional to the small $\Sigma-\Lambda$ mass difference. If one neglects this mass difference, the D_Σ and D_Λ poles coincide, and the residue of this pole is simply

$$-\frac{1}{8}\frac{e^2G_\Lambda^2}{(4\pi)^2}\frac{M_\Lambda}{kWM_N}\mu_N\left(\mu_\Lambda+\frac{G_\Sigma}{G_\Lambda}\mu_T\right),$$

when the $\Sigma-\Lambda$ parity is even. On the other hand, if the $\Sigma-\Lambda$ parity is odd, the quantity $M_\Sigma-M_\Lambda$ in Eq. (8) becomes $(-M_\Sigma-M_\Lambda)$, and several terms become important at the $D_\Sigma=0$ pole.

In order to illustrate the main effects of the backward poles at high energies, we give below the high-energy limit of the Born-approximation terms in the differential $\gamma+p\rightarrow K^++\Lambda$ cross section, expressed in terms of the quantity $y=(1+\cos\theta)$:

$$\left(\cos\theta+\frac{E_\Lambda}{q}\right)\frac{d\sigma}{d\Omega}$$
$$=\frac{1}{16}\frac{e^2G_\Lambda^2}{(4\pi)^2}\frac{1}{kW}\Bigg\{q^2y(2-y)\left[\left(\frac{\mu_N}{M_N}\right)^2+\left(\frac{\mu_\Lambda}{M_\Lambda}+\frac{\mu_TG_\Sigma}{M_\Lambda G_\Lambda}\right)^2\right]$$
$$+\left[y+\frac{\mu_N}{M_N}(M_\Lambda+\tfrac{1}{2}M_Ny)-\frac{\mu_\Lambda}{M_\Lambda}(M_\Lambda+\tfrac{1}{2}M_Ny)\right.$$
$$\left.-\frac{\mu_TG_\Sigma}{M_\Lambda G_\Lambda}(M_\Sigma+\tfrac{1}{2}M_Ny)\right]^2$$
$$-2\frac{\mu_TG_\Sigma}{M_\Lambda G_\Lambda}(M_\Lambda-M_\Sigma)y\Bigg\}. \quad (10)$$

Only the first two terms in an expansion in $(M_i/q)^2$ have been retained, and the $\Sigma-\Lambda$ mass difference has been neglected, so that $M_\Sigma=\pm M_\Lambda$. In this approximation, the D_Σ and D_Λ poles occur at the same place, $y=-\frac{1}{2}M_\Lambda^2/q^2$. In making the extrapolation, values of y of order M_Λ^2/q^2 are important; for such values of y the leading terms of Eq. (10) are the q^2 term and the terms that are independent of y.

If μ_T or μ_Λ is comparable to μ_N, the residues of the backward poles are so large that the effect of these poles may be important at energies much lower than those at which the expansion in terms of M_i^2/q^2 is valid. In order to illustrate this point, we have calculated the differential cross section at the lab photon energy of 1.95 Bev in the Born approximation, assuming $\mu_T=0$ and $\mu_\Lambda/M_\Lambda=\pm\mu_N/|M_N|$. At this energy the D_Λ pole occurs at the point $\cos\theta=-1.95$. The results of this calculation are shown in Fig. 1, for the different assumptions concerning the sign of μ_Λ and the $\Sigma-\Lambda$ parity. (Very similar results are obtained from the assumptions $\mu_\Lambda=0$ and $\mu_T/M_\Lambda=\pm\mu_N/|M_N|$, provided $|G_\Sigma|=|G_\Lambda|$ and the $\Sigma-\Lambda$ parity is even.) It is seen that for these four cases a simple linear extrapolation between the angles $\cos\theta=-0.5$ and $\cos\theta=-1$ gives approximately the correct residue. (Of course, such a simple extrapolation procedure may not be valid with experimentally measured cross sections.[2])

We conclude that at any energy comparable to or greater than 2 Bev, the poles in the unphysical region $\cos\theta<-1$ may produce observable effects. Extrapolation to the backward pole (the two backward poles are very close together, so we consider them as one) in the angular distribution for the processes $\gamma+p\rightarrow K^{+}+\Lambda$ and $\gamma+p\rightarrow K^{+}+\Sigma^{0}$ will yield information concerning the magnetic moment of the produced hyperon and the Σ^0 lifetime, provided that the particle parities and the approximate values of the $K\Lambda N$ and $K\Sigma N$ coupling constants are known. The magnetic moment of the Λ may be measured directly,[11] so that this extrapolation method, when applied to $K^{+}+\Lambda$ production, is essentially a measurement of the Σ^0 lifetime. If the Σ^0 lifetime were known, the Σ^0 anomalous moment could be measured approximately by applying the extrapolation method to $K^{+}+\Sigma^{0}$ photoproduction.

V. REMARKS CONCERNING EXPECTED VALUES OF HYPERON MOMENTS

At present, no reliable theoretical prediction of μ_T, μ_Λ, and μ_Σ can be made, since our knowledge of Σ and Λ interactions is meagre. Various theoretical models do, however, lead to approximate predictions.[12] For example, the global symmetry model of pion interactions with nucleons and baryons, (in which the $\Sigma-\Lambda$ parity is even) leads to the predictions $\mu_\Lambda\approx\mu_{\Sigma^0}\approx 0$ and $\mu_T\approx\mu_N$ if the baryon mass differences and the K meson contributions to the moments are neglected.[5,12] The value $\mu_T=\mu_N$ for the transition moment, when substituted into Eq. (1), results in a Σ^0 lifetime of $\sim 0.8\times 10^{-19}$ sec.

It is generally assumed that the two-pion state is responsible for the largest part of the nucleon anomalous moment. The two-pion state that contributes to form factors is of isotopic spin 1, and hence cannot contribute to the Λ or Σ^0 magnetic moments.[5] Therefore, it would not be surprising if μ_Λ and μ_{Σ^0} were small.

The fact that the $(\pi+\Sigma)/(\pi+\Lambda)$ branching ratio of the Y_1^* is so small indicates that the $\pi+\Sigma$ and $\pi+\Lambda$ states of angular momentum, parity, and energy equal to that of the Y_1^* are not coupled together appreciably. A possible explanation for this lack of coupling results from the assumption that the $\pi\Sigma\Sigma$ coupling constant is zero,[13] since the exchange of pions cannot couple the $\pi+\Lambda$ and $\pi+\Sigma$ channels if the $\pi\Lambda\Sigma$ interaction is the only πY interaction. For a similar reason, the contribution of the two-pion state to the $\Sigma-\Lambda$ transition moment must be small if $f_{\pi\Sigma\Sigma}=0$. In this case one would expect that μ_T, as well as μ_Λ and μ_Σ, is smaller than μ_N.

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[11] R. A. Schluter *et al.* (private communication; and to be published).

[12] See, for example, Hiroshi Katsumori, Progr. Theoret. Phys. (Kyoto) **24**, 1371 (1960); Katsumi Tanaka, Phys. Rev. **122**, 705 (1961); J. Dreitlein and B. W. Lee, *ibid.* **124**, 1274 (1961).

[13] This assumption has recently been made by J. Franklin (to be published) in a model of the various low-energy pion-hyperon resonances.