primarily on the former singularities, whereas the residue of the corresponding pole, and hence σ_t, depends more on the latter.

The main reason for the sensitivity of σ_t in (b) comes from the fact that $\sigma_t^{(\mathrm{in})}$ does not differ much from $\sigma_t^{(\mathrm{out})}$ over a very wide range of values of $\sigma_t^{(\mathrm{in})}$. In other words, $\sigma_t^{(\mathrm{out})}$ approaches $\sigma_t^{(\mathrm{in})}$ very slowly as $\sigma_t^{(\mathrm{in})}$ is varied. In view of all the approximations used, it was therefore thought more meaningful to calculate, not the self-consistent value of σ_t, but rather the value of $\sigma_t^{(\mathrm{out})}$ one obtains if one assumes the experimental value for $\sigma_t^{(\mathrm{in})}$. Using the elastic approximation, such a calculation gives $\sigma_t^{(\mathrm{out})} = 21$ mb, while the corresponding calculation in the black-disk approximation gives $\sigma_t^{(\mathrm{out})} = 17$ mb. These values do not differ much from the experimental value of 14 mb.

It is a great pleasure to be able to thank Dr. B. M. Udgaonkar for several very helpful discussions, and Professor J. R. Oppenheimer for his hospitality at the Institute for Advanced Study.

---

[1]L. A. P. Balázs, Phys. Rev. 128, 1939 (1962); 129, 872 (1963).

[2]L. B. Auerbach, T. Elioff, W. B. Johnson, J. Lach, C. E. Weigand, and T. Ypsilantis, Phys. Rev. Letters 9, 173 (1962); D. D. Carmony and R. T. Van de Walle, Phys. Rev. Letters 8, 73 (1962); J. Button, G. R. Kalbfleisch, G. R. Lynch, B. C. Maglić, A. H. Rosenfeld, and M. L. Stevenson, Phys. Rev. 126, 1858 (1962), to which the reader is also referred for additional references on experimental $\pi-\pi$ scattering.

[3]S. C. Frautschi, M. Gell-Mann, and F. Zachariasen, Phys. Rev. 126, 2204 (1962); G. F. Chew and S. C. Frautschi, Phys. Rev. Letters 7, 394 (1961); 8, 41 (1962); G. F. Chew, S. C. Frautschi, and S. Mandelstam, Phys. Rev. 126, 1202 (1962); R. Blankenbecler and M. L. Goldberger, Phys. Rev. 126, 766 (1962).

[4]G. F. Chew and S. C. Frautschi, Phys. Rev. 123, 1478 (1961).

[5]G. F. Chew and S. Mandelstam, Phys. Rev. 119, 467 (1960).

[6]There is some evidence that this resonance could be identified with the f^0, recently seen by W. Selove, V. Hagopian, H. Brody, A. Baker, and E. Leboy, Phys. Rev. Letters 9, 272 (1962). The mass and half-width of the f^0 are 1250 MeV and 50 ± 25 MeV, respectively.

[7]F. Zachariasen and C. Zemach, Phys. Rev. 128, 849 (1962).

[8]High-energy inelastic effects would be automatically taken into account by the mechanism described in the Appendix of I.

[9]B. M. Udgaonkar and M. Gell-Mann, Phys. Rev. Letters 8, 346 (1962). As in the first paper of reference 4 we have taken $s_0 = 2$.

[10]This value can be deduced from the $\pi-N$ and $N-N$ total cross sections. See M. Gell-Mann, Phys. Rev. Letters 8, 263 (1962); V. N. Gribov and I. Ya. Pomeranchuk, Phys. Rev. Letters 8, 343 (1962).

NUCLEON AS A REGGE POLE AND THE $J=\frac{1}{2}$, $T=\frac{1}{2}$ πN PHASE SHIFTS

N. N. Khuri* and B. M. Udgaonkar†
Institute for Advanced Study, Princeton, New Jersey
(Received 4 February 1963)

The idea that the nucleon may be a composite particle lying on a Regge trajectory has attracted considerable attention recently.[1] Several high-energy experiments have been suggested to test this idea.[2] Also, recently Balázs[3] has investigated the consequences of the assumption that the nucleon is a πN bound state on the low-energy behavior of the $J=\frac{1}{2}$, $T=\frac{1}{2}$ πN phase shifts, and he was able to make a rough prediction of the $P_{1/2}$ phase shift up to ~225 MeV. The purpose of this note is to propose a new "effective-range type formula" for this phase shift, based on the idea that the nucleon is a Regge pole. It will be seen that this formula gives a rather good description of this phase shift up to ~225 MeV.

This investigation was motivated by the fact that, as shown recently by one of us,[4] the scattering amplitude in potential scattering at low to intermediate energies may well be approximated by the contributions from a few leading Regge poles in the same channel. This approximation holds only if one uses the expression for the "full" contribution of a Regge pole which was given in reference 4. This latter expression has the further advantage of having the correct cuts in the z plane, and its partial-wave projection has the correct threshold behavior, two features which are missing in the usual expression for a Regge pole contribution.

The formal analytic continuations into the complex J plane of the various partial-wave amplitudes for πN scattering have been carried out by several authors.[5] We shall follow the notation of Singh. In this problem one, in general, has, for each value of the total isotopic spin, four functions of complex J and W, the total c.m. energy,

$a_{\pm}^{(e,o)}(J, W)$. Here the superscripts e and o refer to even or odd J-parity, respectively, with $a_{\pm}^{(e)}(J, W)$ coinciding with the physical amplitudes for $J=\frac{1}{2}, \frac{5}{2}, \frac{9}{2}, \cdots$; and $a_{\pm}^{(o)}(J, W)$ for $J=\frac{3}{2}, \frac{7}{2}, \frac{11}{2}, \cdots$. The ($\pm$) subscripts differentiate the two orbital amplitudes according to $l=J\pm\frac{1}{2}$. In general, these four amplitudes have different J-plane singularities except for the correlation implied by the MacDowell symmetry,[6] $a_{+}^{(e,o)}(J, W)=-a_{-}^{(e,o)}(J, -W)$. This means that, in particular, if $a_{+}^{(e)}(J, W)$ has a pole at $J=\alpha(W)$, then $a_{-}^{(e)}(J, W)$ will have a pole at $J=\alpha(-W)$.

Assuming the validity of the Mandelstam representation for πN scattering, one can easily show that the above four amplitudes are analytic functions of J regular in the region $\mathrm{Re}J>J_{\min}$, where $J_{\min}$ is related to the number of subtractions in the Mandelstam representation.[7] If one now assumes that these functions can be continued to the left of the line $\mathrm{Re}J=J_{\min}$, then the nucleon, being a $P_{1/2}$ state of the πN system, will demonstrate itself as a pole in $a_{+}^{(e)}(J, W)$ at $J=\alpha(W)$ with $\alpha(m)=\frac{1}{2}$, where m is the nucleon mass. By the symmetry discussed above, there will also be a pole in $a_{-}^{(e)}(J, W)$ at $J=\alpha(-W)$.

If the nucleon trajectory is the only trajectory with quantum numbers $B=1$, $S=0$, $T=\frac{1}{2}$, and even J-parity, then, in the spirit of the approximation proposed in reference 4, the contribution of this Regge trajectory should at low energies give a good approximation to the amplitudes of even J-parity. The argument is basically a familiar polology-type argument. In other words, the function $a_{+}^{(e)}(J, W)$ at $J=\frac{1}{2}$, i.e., the $P_{1/2}$ amplitude, is determined by the contribution of the nearest singularity in the J plane which for values of W just above threshold is just the nucleon pole $J=\alpha(W)$. For low energies, $\mathrm{Re}\,\alpha(W)$ would not have moved too far from its value at $W=m$. It is our aim to check this approximation with experiment.

In the usual form the contribution of the nucleon pole to the two $T=\frac{1}{2}$ Pauli amplitudes for πN scattering is[5]

$$f_2(W,z) = -\frac{\pi}{2}\frac{\beta(W)}{\cos\pi\alpha(W)}\left[P'_{\alpha(W)+\frac{1}{2}}(-z)+P'_{\alpha(W)+\frac{1}{2}}(z)\right] + \frac{\pi}{2}\frac{\beta(-W)}{\cos\pi\alpha(-W)}\left[P'_{\alpha(-W)-\frac{1}{2}}(-z)-P'_{\alpha(-W)-\frac{1}{2}}(z)\right],$$

$$f_1(W,z)=-f_2(-W,z). \qquad (1)$$

Here z is the cosine of the c.m. scattering angle and $\beta(W)$ is the residue of $a_{+}^{(e)}(J, W)$ at $J=\alpha(W)$. The terms in the brackets in (1) have cuts in the z plane starting at $z=1$ and $z=-1$, respectively. However, we know that the correct thresholds in the z plane are given by

$$z=\cosh\xi_1=1+2/k^2 \quad \text{(right-hand cut)};$$

$$z=-\cosh\xi_2=\left[(W^2-m^2-2)/2k^2-1\right] \quad \text{(crossed nucleon pole)}. \qquad (2)$$

We have taken the pion mass to be unity, and k is the c.m. momentum.

To obtain the full contribution of the nucleon pole, we use the method of reference 4 to modify the derivation of (1). We get

$$f_2(W,z) = \left[R_1(W,z;\ \alpha(W)+\tfrac{1}{2})+R_2(W,z;\ \alpha(W)+\tfrac{1}{2})\right] + \left[R_1(-W,z;\ \alpha(-W)-\tfrac{1}{2})-R_2(-W,z;\ \alpha(-W)-\tfrac{1}{2})\right],$$

$$f_1(W,z)=-f_2(-W,z). \qquad (3)$$

The functions R_1 and R_2 are given by

$$R_1(W,z;\ \lambda) = \beta(W)\left\{\frac{-\frac{1}{2}\pi P_\lambda'(-z)}{\cos(\lambda-\frac{1}{2})\pi}+\frac{1}{4\sqrt{2}}\int_{-\infty}^{\xi_1}\frac{\exp[(\lambda+\frac{1}{2})x]}{(\cosh x-z)^{3/2}}dx\right\},$$

$$R_2(W,z;\ \lambda) = \beta(W)\left\{\frac{-\frac{1}{2}\pi P_\lambda'(z)}{\cos(\lambda-\frac{1}{2})\pi}+\frac{1}{4\sqrt{2}}\int_{-\infty}^{\xi_2}\frac{\exp[(\lambda+\frac{1}{2})x]}{(\cosh x+z)^{3/2}}dx\right\}, \qquad (4)$$

where $\lambda=\alpha(W)+\frac{1}{2}$. R_1 has a right-hand cut in the z plane starting at $z=\cosh\xi_1=1+2/k^2$ and no left-hand cut. R_2 has only a left-hand cut starting at $z=-\cosh\xi_2$.

We can now project out the contribution of the nucleon pole to the $P_{1/2}$ partial wave amplitude and obtain

$$a_{+}(\tfrac{1}{2}, W)=\frac{1}{2}\frac{\beta(W)}{\alpha(W)-\frac{1}{2}}\left[\exp\{[\alpha(W)+\tfrac{1}{2}]\xi_1\}\exp(-\xi_1) + \exp\{[\alpha(W)+\tfrac{1}{2}]\xi_2\}\exp(-\xi_2)\right] \qquad (5)$$

The terms with $\alpha(-W)$ do not contribute at all to the $P_{1/2}$ amplitude. However, the $S_{1/2}$ amplitude,

$a_-(\frac{1}{2}, W)$, will only have contributions from the $\alpha(-W)$ terms which can be obtained from (5) by using the MacDowell symmetry.

If the nucleon Regge pole is the only $T=\frac{1}{2}$ trajectory of even J-parity that goes above or near $\mathrm{Re}J=0$ for low energies, then according to the arguments of reference 4 the right-hand side of (5) is a good approximation to the $a_+(\frac{1}{2}, W)$ amplitude in the region where $\xi_1>1$ and $\xi_2>1$. Given $\beta(W)$ and $\alpha(W)$, we can easily compute the expression in (5) and compare with experimental results on the δ_{11} phase shift to check our assumption that there are no other important singularities contributing to the $P_{1/2}$ amplitude. Unfortunately, we do not know β and α, and we have to proceed through an approximation and an extrapolation procedure.

Let us list the properties of $\alpha(W)$, $\beta(W)$, and $a_+(\frac{1}{2}, W)$ which we know and which will help us in estimating (5). First, we know that $\alpha(m)=\frac{1}{2}$, and in the region $m \leq W \lesssim m+2$ it is reasonable to assume that $\mathrm{Re}\alpha(W)$ can be approximated by a straight line. We write

$$\mathrm{Re}\alpha(W) \approx \tfrac{1}{2}+\epsilon(W-m). \qquad (6)$$

The imaginary part of $\alpha(W)$ is zero near threshold, and we shall assume that it is negligible in the low-energy region we are considering. Furthermore, we know that $\beta(W)$ has the following threshold behavior

$$\beta(W) \sim (k^2/s_0)^{\alpha_0+\frac{1}{2}}; \quad \alpha_0=\alpha(W)\big|_{k^2=0}. \qquad (7)$$

We also know that as $W \to m$, $a_+(\frac{1}{2}, W)$ approaches the usual perturbation theory pole term, namely,

$$a_+(\tfrac{1}{2}, W) \to 3f^2[1/(W-m)], \qquad (8)$$

where $f^2=0.08$.

The most obvious way to proceed is to factor out the threshold behavior of $\beta(W)$ and then extrapolate to the pole in (8). However, in that case we will have to determine the constant s_0. To avoid this we follow a path which is more suited to the expression (5) that we have.

The main point is to notice that the exponentials in (5) have the following threshold behavior:

$$\exp[(\alpha+\tfrac{1}{2})\xi_1] \sim (\tfrac{1}{4}k^2)^{-[\alpha_0+\frac{1}{2}]},$$

$$\exp[(\alpha+\tfrac{1}{2})\xi_2] \sim [k^2/(2m-1)]^{-[\alpha_0+\frac{1}{2}]}. \qquad (9)$$

Thus the product $\beta(W)\exp[(\alpha+\frac{1}{2})\xi_1]$ would be slowly varying and essentially real near threshold.

Our approximation consists of writing[8]

$$\beta(W)\exp[(\alpha+\tfrac{1}{2})\xi_1]=C, \qquad (10)$$

where C is a real constant. This is certainly correct near threshold and might not be too far off even for $W \approx m+2$.

We can now substitute (10) and (6) in (5) and obtain

$$a_+(\tfrac{1}{2}, W)=(C/2\epsilon)\{\exp(-\xi_1)+\exp[(\alpha+\tfrac{1}{2})(\xi_2-\xi_1)] \times\exp(-\xi_2)\}/(W-m). \qquad (11)$$

The constant C/ϵ can now be determined by extrapolation and comparison with (8). We get

$$a_+(\tfrac{1}{2}, W)=-\tfrac{3}{2}f^2[\exp(-\xi_1)+\exp\{[1+\epsilon(W-m)](\xi_2-\xi_1)\} \times\exp(-\xi_2)]/(W-m). \qquad (12)$$

This last expression has no arbitrary constants. For the slope ϵ of the trajectory, we take the usual value determined by the position of the N^{***}, in our units $\epsilon=0.4$. Furthermore, it is easy to check both in (5) and (11) that the expressions for $a_+(\frac{1}{2}, W)$ have the threshold behavior of a P-wave amplitude, $a_+(\frac{1}{2}, W)\sim k^2$ as $k^2 \to 0$.

One can now easily compute (12). In Fig. 1 we plot our results for the phase shift δ_{11}, where $a_+(\frac{1}{2}, W)=\sin\delta_{11}\exp(i\delta_{11})/k$. As can be seen from this figure, our result is in fairly good agreement with the experimental data.[9] When compared to previous effective-range type calculations,[10] which all tended to give a phase shift considerably larger in magnitude than that observed, our calculation gives a result which stays small all the way up to ≈200 MeV. For $k^2 \ll 1$, our result gives $a_+ \simeq -0.96f^2k^2$ while the pure perturbation-theory pole gives $a_+ \simeq -3f^2k^2$.

It may be noted that our phase shift does not change sign in the energy range considered. Thus the Fermi (i) and Yang solutions of Deahl et al.[11] at 225 MeV would not be consistent with our results, while their Fermi (ii) solution would.

In closing, we make a few remarks about the other two $T=\frac{1}{2}$ phase shifts, namely, the $S_{1/2}$ and the $P_{3/2}$. It is clear from (5) and the MacDowell symmetry that to compute the contribution of the nucleon pole to the $S_{1/2}$ amplitude we would have to know $\alpha(-W)$ for $m \lesssim W \lesssim m+2$. Extrapolation would be dangerous in this case as it involves going from $W \simeq m$ to $W \simeq -m$. In addition, even

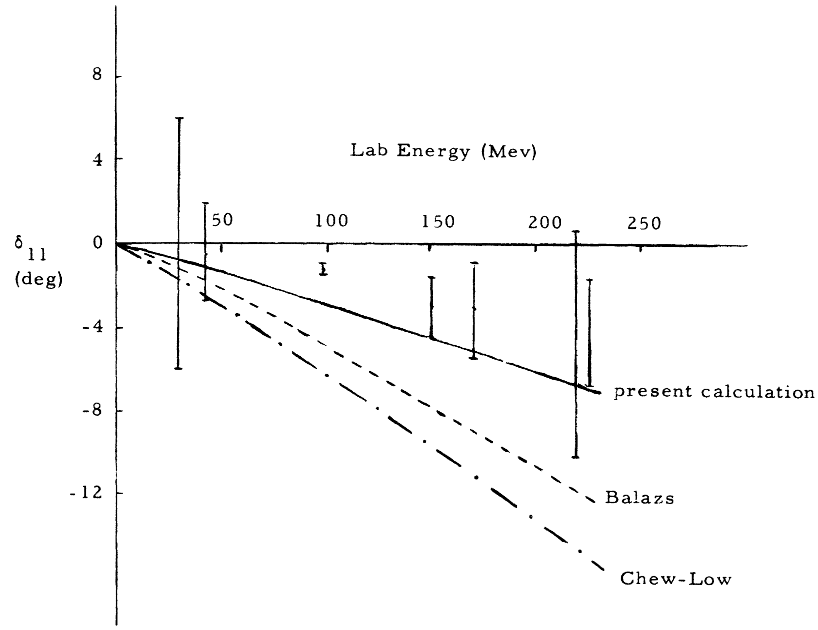

FIG. 1. Plot of δ_{11} phase shift as a function of pion lab energy. The experimental values are taken from S. W. Barnes, M. Winick, K. Miyake, and K. Kinsey, Phys. Rev. 117, 238 (1960) (30 to 98 MeV); from solution A of H.-Y. Chiu and E. L. Lomon, Ann. Phys. 6, 50 (1959) (150 to 220 MeV); and from the Fermi (ii) solution of J. Deahl, M. Derrick, J. Fetkovitch, T. Fields, and G. B. Yodh, Phys. Rev. 124, 1987 (1961).

if we know how to extrapolate that far, it is not clear whether if $\alpha(-W)$ becomes negative this singularity will continue to be the dominant one.

The nucleon pole makes a negligible contribution to the $P_{3/2}$ phase shift δ_{31}, since the latter is connected with an amplitude of odd J-parity. In the framework of the present approach, the main contribution to δ_{31} should come from the Regge trajectory of odd J-parity on which the N^{**} resonance ($D_{3/2}$) lies.

The authors wish to thank Professor J. R. Oppenheimer for the hospitality extended to them at the institute for Advanced Study.

---

*Work supported in part by the National Science Foundation. On leave of absence from the Physics Department of the American University of Beirut, Beirut, Lebanon.

†The study was supported in part by the U. S. Air Force Office of Scientific Research, Research Grant No. AF-AFOSR-42-63. On leave of absence from Tata Institute of Fundamental Research, Bombay, India.

[1]R. Blankenbecler and M. L. Goldberger, Phys. Rev. 126, 766 (1962); G. F. Chew and S. C. Frautschi, Phys. Rev. Letters 7, 394 (1961); M. Gell-Mann and M. L. Goldberger, Phys. Rev. Letters 9, 275 (1962).

[2]G. F. Chew and S. C. Frautschi, reference 1; S. C. Frautschi, M. Gell-Mann, and F. Zachariasen, Phys. Rev. 126, 2204 (1962); V. N. Gribov, Proceedings of the International Conference on High-Energy Nuclear Physics, Geneva, 1962 (CERN Scientific Information Service, Geneva, Switzerland, 1962), pp. 547-551.

[3]L. Balázs, Phys. Rev. 128, 1935 (1962).

[4]N. N. Khuri, Phys. Rev. (to be published).

[5]S. C. Frautschi, M. Gell-Mann, and F. Zachariasen, reference 2; Virendra Singh, Phys. Rev. 129, 1889 (1963); V. N. Gribov and I. Ya. Pomeranchuk (to be published).

[6]S. W. MacDowell, Phys. Rev. 116, 774 (1959).

[7]$J_{\rm min} < \frac{3}{2}$. See Virendra Singh and B. M. Udgaonkar, Phys. Rev. 123, 1487 (1961). See also Virendra Singh, thesis, University of California Radiation Laboratory Report No. UCRL-10254, May 1962 (unpublished).

[8]More generally one can set $\beta\{\exp[(\alpha+\frac{1}{2})\xi_1]+a\exp[(\alpha+\frac{1}{2})\xi_2]\}=c$, where both a and c are real constants. However, in doing the extrapolation one finds out that a must be zero to get a real residue at $W=m$.

[9]The experimental points used by us for comparison with our calculation are the same as those used by Balázs in reference 3. The phase shifts evaluated by Woolcock and reproduced in Fig. 7 of Hamilton et al. [J. Hamilton, P. Menotti, G. C. Oades, and L. L. J.

Vick, Phys. Rev. 128, 1881 (1962).] seem to be in somewhat better agreement with our calculations. See also W. S. Woolcock, Proceedings of the Aix-En-Provence Conference on Elementary Particles, 1961 (Centre d'Etude Nucleaire de Saclay, Seine-et-Oise, France, 1961), Vol. 1, pp. 459, 461.

[10]E.g., reference 3; see also G. F. Chew and F. E. Low, Phys. Rev. 101, 1570 (1956).

[11]J. Deahl, M. Derrick, J. Fetkovich, T. Fields, and G. B. Yodh, Phys. Rev. 124, 1987 (1961).

SPIN OF THE Y_1^{*}†

L. Bertanza,* V. Brisson,‡ P. L. Connolly, E. L. Hart, I. S. Mittra,|| G. C. Moneti,⊥ R. R. Rau, N. P. Samois, I. O. Skillicorn,** and S. S. Yamamoto
Brookhaven National Laboratory, Upton, New York

and

M. Goldberg, J. Leitner, S. Lichtman, and J. Westgard
Syracuse University, Syracuse, New York
(Received 10 January 1963)

A resonance in the $(\Lambda\pi)$ system, Y_1^*, with a mass of 1385 MeV has been reported by many groups.[1] Although most of the properties of this resonance are well known,[2] its spin has not been determined with a high degree of confidence, and its parity is unknown. The strongest evidence for the Y_1^* spin has been obtained by Ely et al.,[3] in the reaction $K^- + p \to \Lambda^0 + \pi^+ + \pi^-$ at 1.1 BeV/c, which indicates a spin $\geq\frac{3}{2}$. However, many other experiments[4] performed under varied conditions failed to give clear confirmation of this result. Furthermore, it has been pointed out by Adair[5] that the spin-$\frac{3}{2}$ result could be stimulated by a spin-$\frac{1}{2}$ Y_1^* interfering with a D-wave background intensity of as little as 5%. In this Letter we present further evidence on the Y_1^* spin, the observed decay correlations indicating a spin $\geq\frac{3}{2}$.

Our sample of Y_1^*'s were produced by the interaction of 2.24-BeV/c K^- mesons in the 20-in. BNL hydrogen bubble chamber via the reaction

$$K^- + p \to \Lambda^0 + \pi^+ + \pi^-. \qquad (1)$$

Details of the exposure and beam are discussed elsewhere.[6] All events consisting of a neutral V decay associated with a two-prong vertex were analyzed using the BNL TRED-KICK system[7]; 326 events were found to fit Reaction (1), 76 of which could be interpreted as Y_1^{*+} production events. The background events which occur with reasonable frequency and which can simulate Reaction (1) are

$$K^- + p \to \Sigma^0 + \pi^+ + \pi^-, \qquad (2)$$

$$\to \Lambda^0 + \pi^+ + \pi^- + \pi^0. \qquad (3)$$

An examination of the following neutral missing mass distribution,

$$M_0 = \{(E_{K^-} + M_p - E_{\pi^-} - E_{\pi^+})^2 - (\vec{p}_{K^-} - \vec{p}_{\pi^-} - \vec{p}_{\pi^+})^2\}^{1/2},$$

shows no peaking at the Σ^0 or π^0 masses, respectively, indicating that contamination from (2) and (3) is small.[8] Detailed studies indicate that the total contamination from all sources is $\leq 20\%$.

The Dalitz plot for Reaction (1) along with the invariant mass distribution for the $(\Lambda\pi^+)$ state are shown in Figs. 1(a) and 1(b). The peak in the mass distribution has been fitted by a Breit-Wigner resonance formula of the form $[(m - m_0)^2 + (\frac{1}{2}\Gamma)^2]^{-1}$, giving a Y_1^* mass m_0 of 1380 ± 3 MeV, and a half-width $\frac{1}{2}\Gamma$, of 25 ± 5 MeV, in excellent agreement with the currently accepted values.[2] The strong production of the Y_1^* in the positive charge state, along with little or no production of either ρ^0 or Y_1^{*-}, are evident features of Fig. 1(a). As a result, the interference between resonant channels is negligible. Furthermore, at this energy, dynamical interference effects[9] between the pions must be small because of the high velocity of the Y_1^* in the $K^- - p$ rest frame.

If the Y_1^*'s are polarized in the production process, parity conservation constrains the polarization to lie along the normal to the production plane, $\hat{n} \propto \hat{p}_{K^-} \times \hat{p}_{Y^*}$. The distribution of $\hat{n}\cdot\hat{p}_{\pi^+}$, where $\hat{p}_{\pi^+}$ is the direction of the π^+ from the Y_1^* decay as measured in the Y_1^* rest frame, may be used to investigate the spin if the Y_1^*'s are either polarized or aligned. Specifically, the predicted distribution is isotropic for $J = \frac{1}{2}$, and of the form $1 + A[\hat{n}\cdot\hat{p}_{\pi^+}]^2$ for $J = \frac{3}{2}$. The observed distribution in $[\hat{n}\cdot\hat{p}_{\pi^+}]$ for the 76 events in which the Y_1^* mass lies between $1340 < m_0 < 1420$ is shown in Fig. 2.