Teor. Fiz.— Pis'ma Redakt. 2, 105 (1965) [translation: JETP Letters 2, 65 (1965)]; H. J. Lipkin and F. Scheck, Phys. Rev. Letters 16, 71 (1966); H. J. Lipkin, ibid. 16, 1015 (1966); J. J. Kokkedee and L. Van Hove, Nuovo Cimento 42, 711 (1966).

[4]The unit of electric charge adopted is the proton charge.

[5]By fundamental particles we shall mean simply noncomposite particles. We use the notation FP for fundamental particle. The name quark is reserved for the fractionally charged fundamental particles of the Gell-Mann model.

[6]F. Gürsey, T. D. Lee, and M. Nauenberg, Phys. Rev. 135, B467 (1964); T. D. Lee, Nuovo Cimento 35, 933 (1965).

[7]Gell-Mann, Ref. 1.

[8]We restrict ourselves to relations involving total cross sections for which there exist relatively good experimental data at high energies. Except for relation (2), these relations are not new and have been derived previously in the context of the quark model (see Ref. 3).

[9]A, B, and $C_{1,2}$ refer to the quark, singlet-triplet, and triplet-triplet models, respectively. In model C_1 the pseudoscalar spin-zero mesons are assigned the structure $|\bar{\alpha}\alpha\rangle$ while in model C_2 they correspond to $|\bar{\beta}\beta\rangle$.

[10]SU(2) invariance has been tested up to energies of 20 BeV to be a fairly good symmetry of strong interactions, broken only by the electromagnetic interaction [see W. Galbraith et al., Phys. Rev. 138, B913 (1965)].

[11]Lee, Ref. 6.

[12]The essential reason for this is that the construction of particles with integral baryon number requires the use of three quarks (or multiples of three, excluding possible quark-antiquark pairs). Such three-quark combinations do not contain an SU(3) triplet representation. We cannot therefore build the triplet t out of quarks.

[13]The superscripts a_i, b_i, $\cdots$ are used to differentiate between different particles, or between particles and antiparticles.

[14]I. Ya. Pomeranchuk, Zh. Eksperim. i Teor. Fiz. 30, 423 (1956) [translation: Soviet Phys.—JETP 3, 306 (1956)]; L. B. Okun' and I. Ya. Pomeranchuk, Zh. Eksperim. i Teor. Fiz. 30, 424 (1956) [translation: Soviet Phys.—JETP 3, 307 (1956)].

[15]Lipkin and Scheck (Ref. 3) have already pointed out that relations (3)-(5) are trivially satisfied if one neglects charge exchange in the very high-energy limit.

[16]I. Ya. Pomeranchuk, Zh. Eksperim. i Teor. Fiz. 34, 725 (1958) [translation: Soviet Phys.—JETP 34, 499 (1958)].

FORMATION OF RESONANT STATES IN $K^- + p \to K^- + p$ BETWEEN 0.8 AND 1.2 GeV/c

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In a study of K^-p elastic scattering at 22 momenta between 0.777 and 1.182 GeV/c, we observe the formation of $Y_1^*(1760)$ and $Y_1^*(1820)$. An analysis of the total elastic and differential elastic cross sections, based on a model in which the background is described by diffraction scattering, yields for the two resonances the following parameters:

$$M_2 = 1758 \pm 11 \text{ MeV}, \quad \Gamma_2 = 113 \pm 25 \text{ MeV},$$

$$x_2 = 0.46 \pm 0.05, \quad M_3 = 1811 \pm 4 \text{ MeV},$$

$$\Gamma_3 = 73 \pm 10 \text{ MeV}, \quad x_3 = 0.67 \pm 0.08.$$

Our data are part of a comprehensive study, in the Saclay 81-cm hydrogen bubble chamber, of K^-p interactions between 0.777 and 1.226 GeV/c.[1-3] Preliminary results of this work have been presented elsewhere.[3,4] Approximately 47 000 photographs were scanned and rescanned and ~13 000 two-prong events were found within a restricted fiducial volume. Of these ~10 000 were fitted to the reaction

$$K^- + p \to K^- + p \tag{1}$$

while the remainder included mainly $KN\pi$, and $\pi^+\pi^-$ with unobserved Λ^0 or Σ^0 final states. $\Sigma^\pm\pi^\mp$, $\Lambda\pi^0$, and $\bar{K}^0n$ final states in which the Σ, Λ, or $\bar{K}^0$ decayed close to the vertex, so as to simulate two-prong events, were also identified and rejected.

The measurements of cross sections were based on K^- path lengths obtained from a track count in all scanned frames. Corrections to the path length accounted for the K^- attenuation due to decay and interaction, observational losses, track curvature, and beam contam-

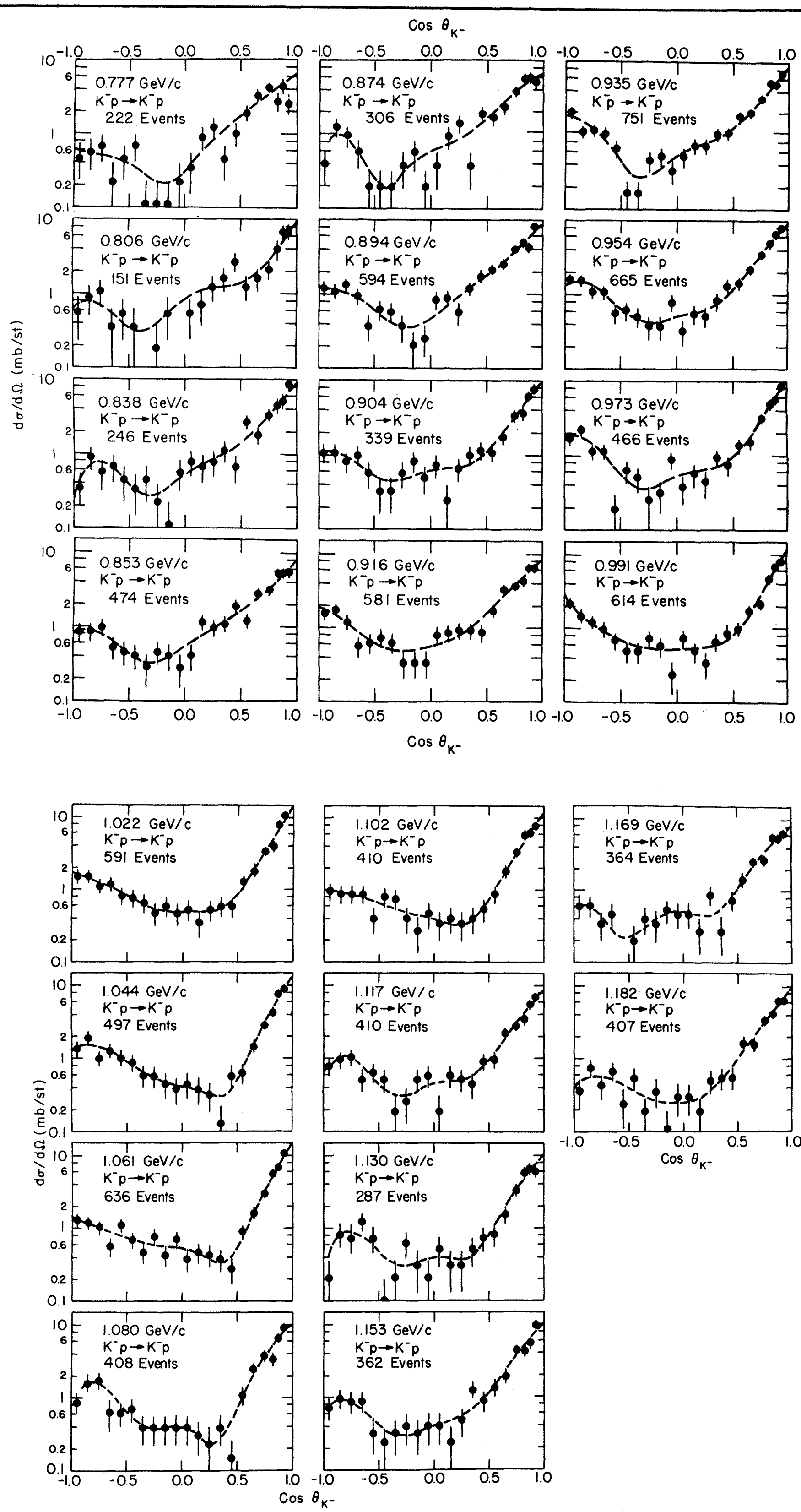

FIG. 1. Differential elastic cross sections fitted to the expansion $d\sigma/d\Omega = \lambda^2 \sum_n A_n P_n(\cos\theta)$. The curves represent the fit for $n = 6$.

ination. The latter was measured from δ-ray counts and found to vary from ~5% in the central momentum region to ~10% at the extremes.

The differential elastic scattering cross sections are shown in Fig. 1 for c.m. system scattering angles $-1.0 \leq \cos\theta_{K^-} \leq 0.95$. Small momentum- and angle-dependent corrections have been introduced in the range $0.80 \leq \cos\theta_{K^-} \leq 0.95$ to account for observational losses of events occurring in a plane normal to the front window of the bubble chamber.[5]

The differential cross sections have been expanded in the series

$$d\sigma/d\Omega = \lambda^2 \sum_n A_n P_n(\cos\theta), \tag{2}$$

where $A_0 = \sigma/4\pi\lambda^2$. Above ~900 MeV/$c$, a fifth-order fit is required by the data. We have, on the other hand, used sixth-order fits to represent the data over the entire K^- momentum region. The Legendre polynomial coefficients A_0 through A_6 are shown in Fig. 2 as a function of K^- momentum. They are consistent with those obtained in other experiments.[6,7]

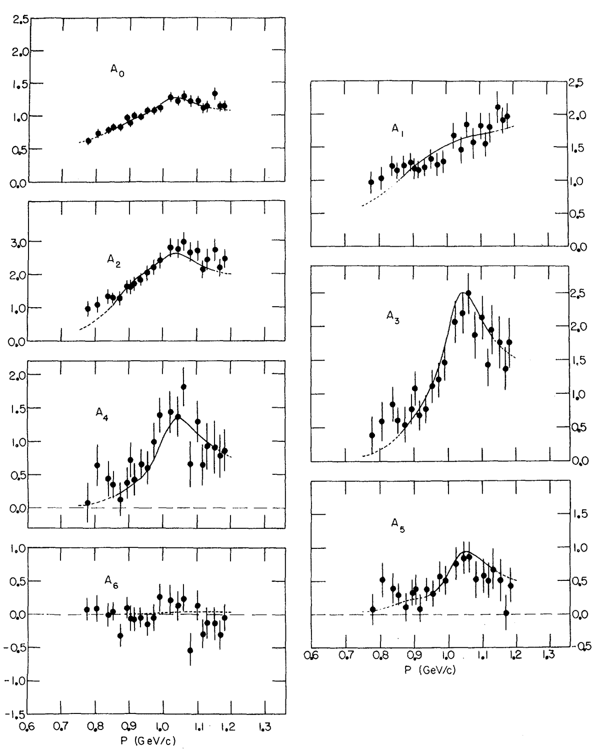

FIG. 2. Legendre polynomial coefficients from a fit to the differential cross sections for $n = 6$. The curves correspond to the nine-parameter BESTRAB fit for the parameters given in the text. The full lines correspond to the momentum region considered in the fit. The dashed line is an extrapolation.

Clear evidence for the formation of $Y_0^*(1820)$ is observed in the momentum dependence of A_2, A_3, A_4, and A_5. The evidence for the formation of $Y_1^*(1760)$, although less obvious, is indeed also very strong. In the first place, the very existence of a peak in A_5, for positive values of A_5 in $K^- + p \to K^- + p$ and for negative values of A_5 in $K^- + p \to K^0 + n$,[3,8] implies strong $I = 0, 1$ interference following the argument originally given by Barbaro-Galtieri, Hussain, and Tripp.[9]

The simplest assumption, however, that only the interference term between $D_{5/2}$, $I = 1$ [for $Y_1^*(1760)$] and $F_{5/2}$, $I = 0$ [for $Y_0^*(1820)$] resonating amplitudes contribute to A_5, leads already to disagreement with our data, for any choice of acceptable parameters for the two resonances. Furthermore, for the lower coefficients, the contribution of the two resonances, even brought to their unitarity limit, accounts for only a small part of the observed Legendre coefficients. Thus, large background contributions to the observed cross sections must be present and should be determined in order to extract from the data the resonant parameters.

A model, discussed elsewhere by Levi Setti and Predazzi,[10] accounts for the background contribution in terms of diffraction scattering which was parametrized by describing the forward peak in the elastic differential cross sections in the form

$$\frac{d\sigma}{dt} \equiv \frac{\pi}{k^2}\frac{d\sigma}{d\Omega} \approx A(k)e^{Bt}. \tag{3}$$

An inspection of Fig. 1 reveals that indeed our data exhibit an exponential-like behavior for small momentum transfers. In an attempt to interpret our data, then, as due to a superposition of resonant and diffraction scattering, a scattering amplitude was constructed which contained a resonant term only for the $l = 2$ and $l = 3$ partial waves [corresponding to $Y_1^*(1760)$ and $Y_0^*(1820)$], and in addition, a diffraction term in the spin-nonflip part of the amplitude:

$$g(k,\theta) = ik\frac{(c_1 + ic_2)}{(\pi k)^{1/2}}e^{bt/2} + \frac{1}{k}\sum_{l=l_R}[(l+1)a_{l+}^R + l a_{l-}^R]P_l(\cos\theta),$$

$$h(k,\theta) = \frac{1}{k}\sum_{l=l_R}[a_{l+}^R - a_{l-}^R]\sin\theta\frac{dP_l(\cos\theta)}{d\cos\theta}. \tag{4}$$

In the above expressions, l_R refers to the resonant angular-momentum states and $a_{l\pm}^R$ are the resonant partial-wave amplitudes, which for $K^- + p \to K^- + p$ are parametrized by

$$a^R = \tfrac{1}{2}x/(\epsilon - i), \qquad (5)$$

where $x = \Gamma_e(k)/\Gamma(k)$ is the elasticity of the resonance and $\epsilon = 2(E - E_R)/\Gamma(k)$. The usual l and k dependence of the widths was taken into account.[11] The first term in $g(k, \theta)$ represents the diffractive amplitude, where t is the momentum transfer squared and $c_1 + ic_2$ is a generalized amplitude for $t = 0$, corresponding to a partly absorptive interaction (for $c_2 = 0$, the interaction becomes completely absorptive).

Expressions for the Legendre polynomial coefficients of the angular distributions have been derived[10] from Eq. (4) and used in a fit called BESTRAB to the present data using the program MINFUN,[12] operating in its minimizing mode. Only the data for A_0 through A_5, in the broad enhancement between 0.87 and 1.13 GeV/c, have been included in the fit, since there are preliminary indications of other possible resonant effects beyond the above limits.[8] The nine-parameter fit (six parameters for the two resonances, three for the diffraction) gave a χ^2 of 72 for 90 data points and 80 degrees of freedom. The fitted resonance parameters are $M_2 = 1758 \pm 11$ MeV, $\Gamma_2 = 113 \pm 0.25$ MeV, and $x_2 = 0.46 \pm 0.05$ for $Y_1^*(1760)$, and $M_3 = 1811 \pm 4$ MeV, $\Gamma_3 = 73 \pm 10$ MeV, and $x_3 = 0.67 \pm 0.08$ for $Y_0^*(1820)$. The imaginary and real part of the forward diffractive amplitudes are $c_1 = 3.73 \pm 0.12$ mb$^{3/4}$ and $c_2 = 0.89 \pm 0.39$ mb$^{3/4}$, respectively. The slope of the forward peak due to diffraction is $b = 1.24 \pm 0.05$ mb. Finally, the curves which fit the A_n coefficients are shown in Fig. 2. The value for the real part of the forward diffractive amplitudes is also consistent with zero. In fact, an eight-parameter fit, obtained by setting $c_2 = 0$, could also satisfactorily reproduce the data.[10]

These resonant parameters agree with independent determinations[2,3,8,13-15] and provide convincing evidence also for the formation of $Y_1^*(1760)$. We would like to remark in this connection that in spite of the success in describing the background amplitudes, resonant parameters deduced from this channel are probably not as reliable as those from a study with other more favorable channels. Thus, our data have been also analyzed in conjunction with the $K^- + p \to \overline{K}^0 + n$ results.[8]

The limitations of the present preliminary interpretation of the $K^- + p \to K^- + p$ data consist primarily in (a) the neglect of the spin dependence of the background contributions; (b) the choice of a very crude parametrization of the diffraction amplitude; and (c) the neglect of the effects of resonances other than $Y_1^*(1760)$ and $Y_0^*(1820)$. Furthermore, no attempt has yet been made to parametrize the background waves in $K^- + p \to \overline{K}^0 + n$, which would enable us to separate the isospin amplitudes contributing to diffraction scattering.

On the other hand, the above analysis is an attempt at interpreting the nature of the $K^- + p \to K^- + p$ interaction at intermediate energies, and within accepted limitations, such attempt seems encouraging.

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[3]R. Armenteros, M. Ferro-Luzzi, D. W. G. Leith, A. Minten, R. D. Tripp, H. Filthuth, V. Hepp, E. Kluge, H. Schneider, R. Barloutaud, P. Granet, J. Meyer, J.-P. Porte, J. Doede, N. M. Gelfand, R. Hildebrand, R. Klem, R. Levi Setti, W. Männer, and C. Rose, reported by M. Ferro-Luzzi, in Proceedings of the Argonne National Laboratory Accelerator Group Users Meeting, May, 1965 (unpublished), p. 87.

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OBSERVATION OF THE THERMAL FLUCTUATIONS OF A GRAVITATIONAL-WAVE DETECTOR*

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We report operation of apparatus to measure the Fourier transform of the Riemann curvature tensor at sensitivity limited by the thermal fluctuations. The gravitational interaction drives a mechanical system which in turn is coupled to the electromagnetic field. Strains as small as a few parts in 10^{16} are observable for a compressional mode of a large cylinder.

The normal modes of an elastic body having appropriate symmetry may be excited by the dynamical curvature tensor.[1] A detector of gravitational radiation[1] may consist of a carefully suspended mass with instrumentation to observe the amplitude of the normal modes. We report observation of the thermal fluctuations of a $1\frac{1}{2}$-ton aluminum cylinder of length ~150 cm and diameter ~61 cm in the vicinity of its lowest compressional mode near 1657 cps (see Fig 1).

At these frequencies the center of mass[1] of the detector on the earth's surface behaves almost as though it were in free fall.[1] A normal coordinate[1] system is appropriate with the pole at the center of mass of the detector. In these coordinates, with the axis of the cylinder in the X^1 direction, the detector output measures the Fourier transform of the component R_{1010} of the Riemann tensor.

The mean squared relative displacement of the cylinder end faces, $\langle X^2\rangle$, associated with the thermal fluctuations is given for each degree of freedom approximately by

$$\tfrac{1}{2}m\omega^2\langle X^2\rangle = \tfrac{1}{2}kT. \qquad (1)$$

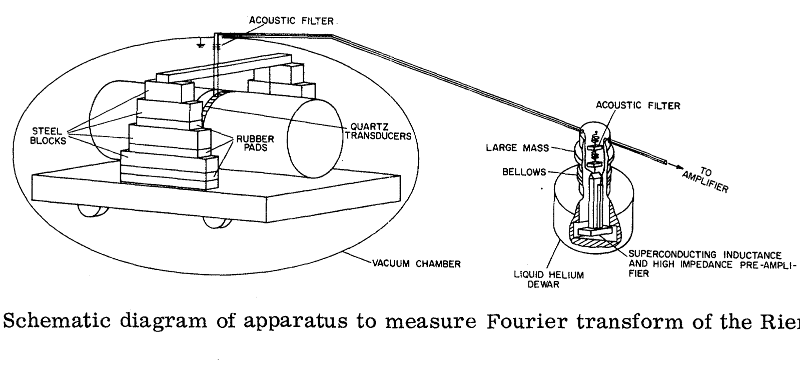

FIG. 1. Schematic diagram of apparatus to measure Fourier transform of the Riemann tensor.