PHYSICAL REVIEW LETTERS

VOLUME 46 | 27 APRIL 1981 | NUMBER 17

Existence of Dibaryon Resonances in $I=1$, $^1D_2$ and $^3F_3$ Nucleon-Nucleon Scattering

R. Bhandari, R. A. Arndt, and L. D. Roper
Center for Analysis of Particle Scattering, Department of Physics, Virginia Polytechnic Institute and State University, Blacksburg, Virginia 24061

and

B. J. VerWest
Department of Physics, Texas A & M University, College Station, Texas 77843

(Received 16 October 1980)

Recent, precise analyses of p-p and n-p scattering data up to 800 MeV by Arndt *et al.* have provided the strongest evidence to date for the existence of dibaryon states in the $I=1$, $^1D_2$ and $^3F_3$ nucleon-nucleon channels. Model fits to their new phases reveal poles located near the "$N\Delta$" threshold ($2.15-0.05i$ GeV). Because of their strong coupling to this channel, these dibaryon resonances are highly inelastic.

PACS numbers: 13.75.Cs, 11.80.Gw, 25.10.+s, 25.40.Rb

Previous speculation on the existence of dibaryon resonances[1] has rested upon studies of incomplete scattering data. For example, in 1968 Arndt[1] investigated the possibility of a dibaryon resonance in the $^1D_2\,p$-p partial-wave amplitude and found that the fit most consistent with his set of phase shifts [existing only up to T_L (laboratory kinetic energy) $=400$ MeV] and a single, very imprecise datum point[2] at $T_L=660$ MeV revealed a pole close to the "$N\Delta$" threshold. Ever since then considerable experimental progress has been achieved, and the situation has changed remarkably.[3] The data available now allow partial-wave analyses over a broader energy range in addition to a more precise determination of the N-N partial-wave amplitudes. Recent, comprehensive analyses of the world data by Arndt *et al.*[4] show sharp energy variations for the $I=1$, $^1D_2$ and the $^3F_3$ phases in the ~2.08–2.25-GeV center-of-mass energy region. This structure correlates with the structure[3] in $\Delta\sigma_L$ (the difference between the p-p total cross sections for parallel and antiparallel longitudinal spin states) observed in approximately the same energy region. The latter has been interpreted by Hidaka *et al.*[5] as manifesta-

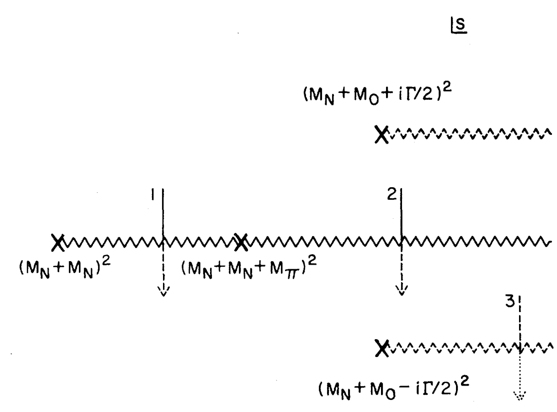

FIG. 1. Right-hand cut structure of the T matrix in the complex s plane. Each arrow leads to a different unphysical sheet.

© 1981 The American Physical Society

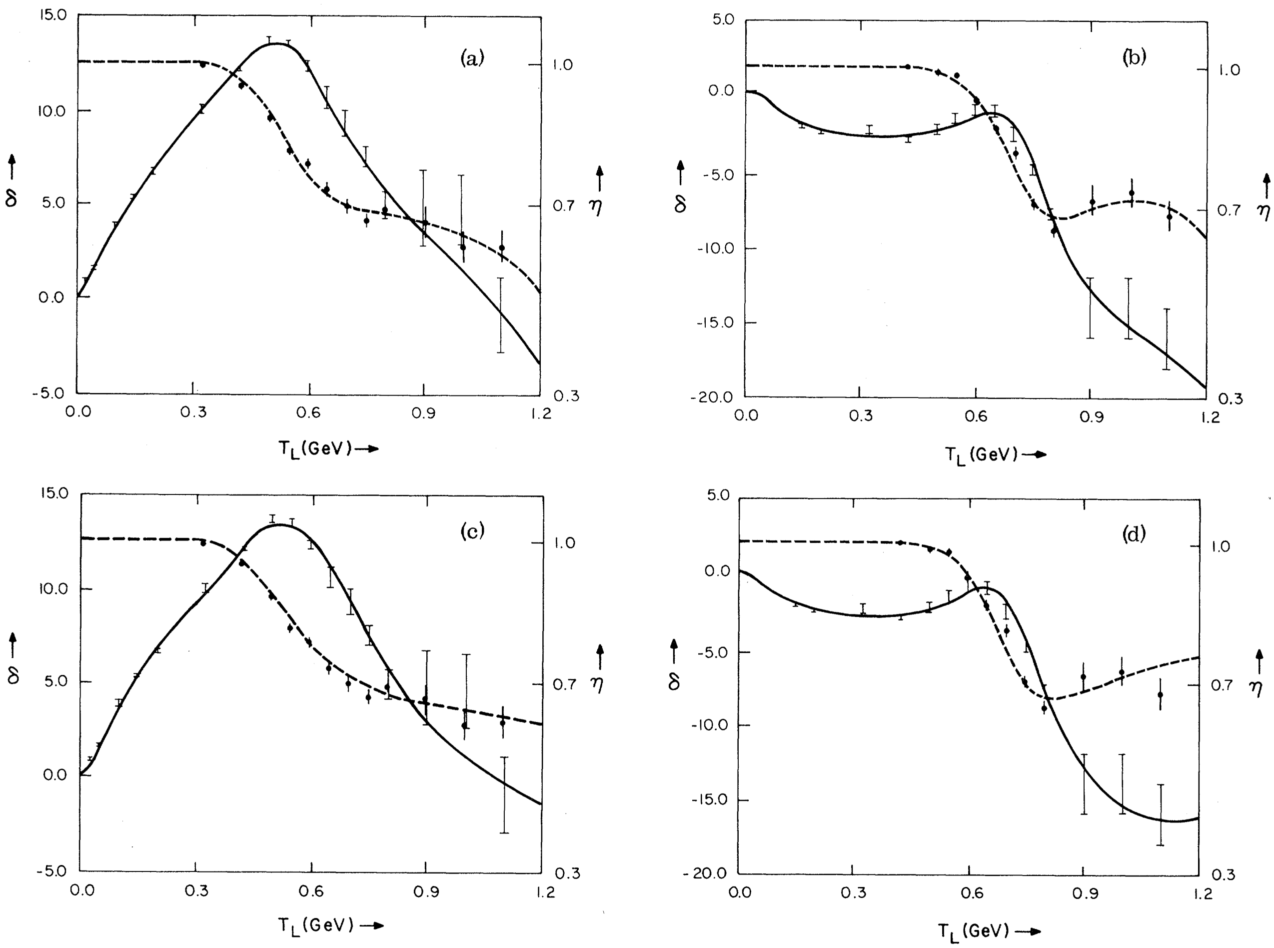

Fig. 2. M-matrix fit to (a) the $I=1$, ${}^1D_2$ and (b) the $I=1$, ${}^3F_3$ nucleon-nucleon partial waves of Arndt *et al.* (Ref. 4); K-matrix fit to (c) the $I=1$, ${}^1D_2$ and (d) the $I=1$, ${}^3F_3$ nucleon-nucleon partial waves of Arndt *et al.* (Ref. 4). Continuous line is the fit to the phase shift δ (in degrees) while the broken line is the fit to the elastic parameter η. T_L is the laboratory kinetic energy of the incident nucleon.

tion of the ${}^3F_3$ resonance. Suggestions for ${}^1D_2$ and ${}^3F_3$ resonances have also been made recently by Hoshizaki[6] from single-channel fits (a Breit-Wigner form plus a smoothly varying background) to his set of phase shifts. The purpose of this paper is to pursue the possibility of resonances in the $I=1$, ${}^1D_2$ and ${}^3F_3$ partial-wave amplitudes, by using as input the new, precise phase shifts of Arndt *et al.*[4] Because of the small size of errors on these phase shifts, we believe that a proper coupled-channel T-matrix fit[7] should be able to distinguish between a resonance and a nonresonance representation.

The S matrix for a system of coupled channels can be expressed as

$$S(s)=1+2i\{\mathrm{Re}[\rho(s)]\}^{1/2}T(s)\{\mathrm{Re}[\rho(s)]\}^{1/2}, \quad (1)$$

where T is the reduced scattering amplitude matrix, and ρ is a diagonal matrix of phase-space factors for the channels; s is the familiar Mandelstam variable, equal to the square of the center-of-mass energy E. The unitarity condition, $S^\dagger S=I$, immediately leads to

$$\mathrm{Im}[T^{-1}(s)]=-\mathrm{Re}[\rho(s)]. \quad (2)$$

In N-N scattering, inelasticity at intermediate energies is due to pion production and originates mainly in the $N\Delta$ channel.[8] Thus for our purposes it is adequate to use a 2×2 matrix representation in which the $N\Delta$ accounts for the inelasticity. Consequently, for the ρ matrix, we need two phase-space factors, ρ_e for the NN channel and ρ_i for the $N\Delta$ channel. For our calculations, we take these to be

$$\rho_e = [(s - s_e)/(s - c_e)]^{l_e + 1/2}, \tag{3a}$$

$$\rho_i = \frac{1}{(E - c_i)^{l_i + 1/2}} \int_{M_T = M_N + M_\pi}^{\infty} \frac{[E - (M_N + M)]^{l_i + 1/2} (M - M_T)^{3/2}}{(M + \alpha)^{l_i + 2}[(M - M_0)^2 + \Gamma^2/4]} \, dM; \tag{3b}$$

l_e and l_i are the orbital angular momentum in the elastic and inelastic channel, respectively. $E = \sqrt{s}$ and $s_e = (M_N + M_N)^2$; M_0 and $-\Gamma/2$ are the real and imaginary parts of the complex mass of Δ, $M_0 - i\Gamma/2$; c_e, c_i, and α are adjustable real constants. ρ_e and ρ_i provide the right-hand unitarity cuts for the T matrix. While the cut due to ρ_e is of a square-root nature, the right-hand cut at the three-body threshold, $E_i = M_N + M_N + M_\pi$, originating in ρ_i, is of a logarithmic nature.[9] One observes that $\mathrm{Re}(\rho_i)$ behaves as $(E - E_i)^{l_i + 3}$ near the three-body threshold. On the unphysical sheets attached to this cut are square-root branch points at complex-conjugate positions: $E_\pm = M_0 + M_N \pm i\Gamma/2$, and also at $E = M_N - \alpha$. For suitable values of α, the latter can be pushed out to the left away from the elastic threshold. The discontinuity across the right-hand cuts associated with branch points at $E = E_\pm$ behave as $(E - E_\pm)^{l_i + 1/2}$ near $E = E_\pm$. The form (3b) possesses the threshold and analytic properties required of a quasi-two-body channel, and in addition has the advantage that it can be evaluated analytically.[9] Figure 1 shows the right-hand cut structure of the T matrix with arrows indicating how different unphysical sheets can be reached. One notes that the complete T matrix can be obtained from Eq. (2) by writing

$$T^{-1} = A - i\rho, \tag{4}$$

where $A = K^{-1}$ or M, K and M being real symmetric matrices, free of threshold cuts.[10] Consequently, we can use the parametrization

$$K_{ij} \text{ or } M_{ij} = \sum_{m=1}^{N} a_m^{(ij)} T_L^{m-1}, \tag{5}$$

where T_L, the laboratory kinetic energy, is related to s by

$$T_L = [s - (M_N + M_N)^2]/2M_N. \tag{6}$$

The phase shifts of Arndt *et al.*[4] exist presently up to $T_L = 800$ MeV. They were supplemented in the fits with the data points of Hoshizaki[11] from 900 to 1100 MeV. Parameters $a_m^{(ij)}$ of Eq. (5) were varied and all good fits to the data required a maximum of up to twelve parameters. The parameters c_e, c_i, and α, on the other hand, were fixed, typical values being 2.6, 0, -0.8, respec-

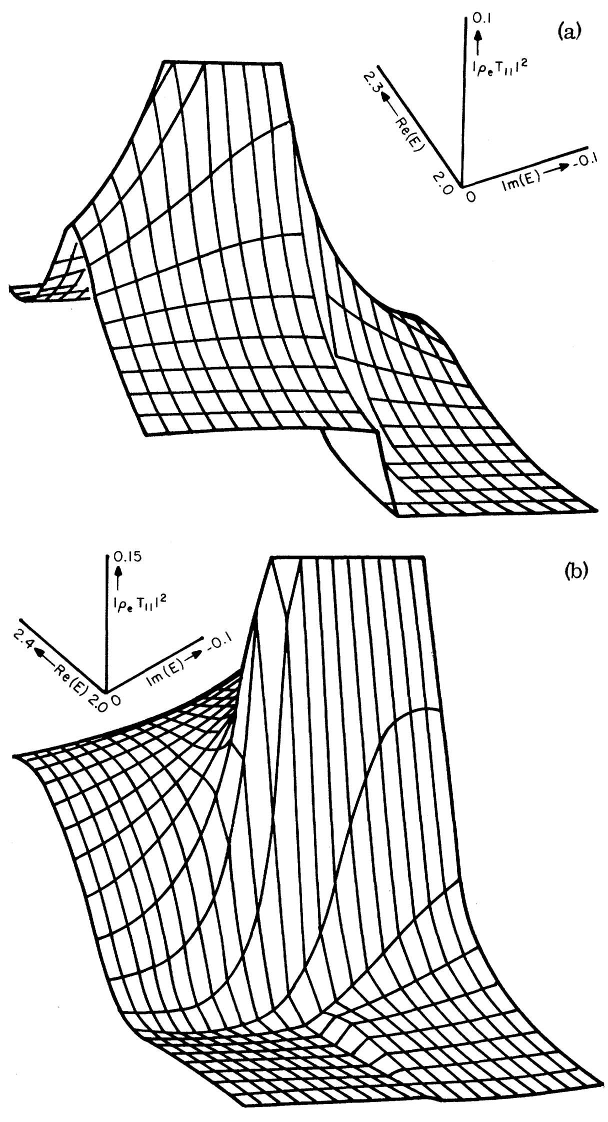

FIG. 3. Three-dimensional plots of $|\rho_e T_{11}|^2$ for (a) the $I = 1$, ${}^1D_2$ and (b) the $I = 1$, ${}^3F_3$ partial waves. Since the $N\Delta$ cut is shown running to the left, $|\rho_e T_{11}|^2$, calculated below the real axis, corresponds to the upper part of the sheet corresponding to arrow 2 and the lower part of the sheet corresponding to arrow 3 (see Fig. 1). The precise energy region (in GeV) and the maximum value of $|\rho_e T_{11}|^2$ are indicated at the top of each figure. The peaks (truncated) are due to poles.

TABLE I. Resonance parameters for the observed $I = 1$, $^1D_2$ and $^3F_3$ dibaryons. $|R|$ is the magnitude of the elastic residue of the poles. $E_p = E_R - i\Gamma_R/2$. The sheet on which a pole lies is indicated by the number of the corresponding arrow (see Fig. 1).

	Solution type	Pole position (E_p) E_R (GeV)	$\Gamma_R/2$ (GeV)	Arrow No.	$2\|R\|/\Gamma_R$
$^1D_2$	M matrix (best)	2.12–2.15	0.08–0.10	2	0.1–0.3
	M matrix (II best)	2.14–2.15	0.05–0.07	2,3	0.1–0.2
	K matrix (best)	2.04–2.05	0.10–0.12	2	0.15–0.20
	K matrix (II best)	2.13–2.14	0.04–0.05	2	0.1–0.15
$^3F_3$	M matrix	2.21–2.22	0.06–0.08	3	0.1–0.2
	K matrix	2.18–2.20	0.06–0.07	3	0.1–0.2

tively. l_i was 0 or 1, depending upon whether it was the $^1D_2$ or the $^3F_3$ partial wave under consideration. The mass of the Δ isobar was taken to be $1.21 - 0.05i$ GeV while values of 0.140 and 0.940 GeV were used for M_π and M_N, respectively. Figure 2 illustrates our best M- and K-matrix fits to the $^1D_2$ and $^3F_3$ phase shifts. The T matrix corresponding to such fits was analytically continued into unphysical sheets along the arrows of Fig. 1, and a search revealed poles near the $N\Delta$ branch point in each one of the above partial values. The influence of these poles on the real energy axis is illustrated in Figs. 3(a) and 3(b), which are three-dimensional plots of $|\rho_e T_{11}|^2$ calculated from the fits. The pole positions as determined from the fits are given in Table I. They lie close to the $N\Delta$ branch point which is at $2.15 - 0.05i$ GeV. Near such poles, the T matrix can be expressed as

$$T_{kj} = \gamma_k \gamma_j / (E_p - E) + B_{kj}, \quad (7)$$

where $E_p = E_R - i\Gamma_R/2$, and B is a background matrix. The residue, R, of the pole for the elastic Argand amplitude, $T_e = \rho_e T_{11}$, is then equal to $\rho_e \gamma_1^2$. If we regard the quantity $|R|/(\Gamma_R/2)$ as a measure of elasticity, we find that the $^1D_2$ and the $^3F_3$ resonances are indeed highly inelastic. Results pertaining to elasticity are also summarized in Table I.

In summary, we find that fits to the $I = 1$, $^1D_2$ and $^3F_3$ scattering phases of the analyses of Arndt *et al.*[4] reveal poles coupled strongly to the $N\Delta$ channel. The precise pole positions are uncertain and also depend to some extent upon the particular parametrization scheme which is being employed. The existence of poles, nevertheless, seems to be a compelling feature of the above partial-wave analyses which, it appears, will be very difficult to fit with a nonresonance hypothesis.

This work was supported by the U. S. Department of Energy.

[1]R. A. Arndt, Phys. Rev. 165, 1834 (1968); L. M. Libby and E. Predazzi, Lett. Nuovo Cimento 18, 881 (1969); J. H. Hall, T. A. Murray, and L. Riddiford, Nucl. Phys. B12, 573 (1969); H. Suzuki, Prog. Theor. Phys. 54, 143 (1975); D. D. Brayshaw, Phys. Rev. Lett. 37, 1329 (1976).

[2]L. S. Azhgirey *et al.*, Phys. Lett. 6, 196 (1963).

[3]For a review of experimental situation, see A. Yokosawa, Argonne National Laboratory Report No. ANL-HEP-PR-80-16 (to be published).

[4]R. A. Arndt, L. D. Roper, B. J. VerWest, Robert Clark, R. A. Bryan, and P. Signell, "Nucleon-Nucleon Partial-Wave Analyses from the Virginia Polytechnic Institute and State University Interactive Dial-In System" (to be published).

[5]K. Hidaka *et al.*, Phys. Lett. 70B, 479 (1977).

[6]N. Hoshizaki, Prog. Theor. Phys. 60, 1796 (1978), and 61, 129 (1979).

[7]A recent coupled-channels T-matrix fit to the $^1D_2$ phase shifts by B. J. Edwards and G. H. Thomas, Argonne National Laboratory Report No. ANL-HEP-PR-80-30 (to be published), uses the older set (year 1978) of the phase shifts of Arndt *et al.*

[8]S. Mandelstam, Proc. Roy. Soc. (London) A244, 491 (1958); D. V. Bugg, J. Phys. G5, 1349 (1977); I. Duck and E. Umland, to be published.

[9]R. Bhandari, Virginia Tech Report No. CAPS-80-1 (to be published).

[10]R. H. Dalitz and S. Tuan, Ann. Phys. (N.Y.) 10, 307 (1960); M. Ross and G. Shaw, Ann. Phys. (N.Y.) 13, 147 (1961).

[11]N. Hoshizaki, private communication.