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

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(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-

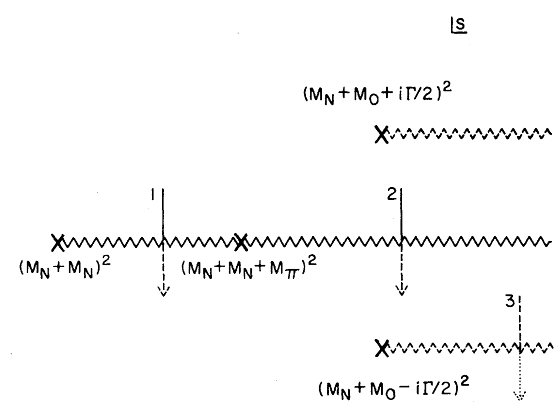

FIG. 1. Right-hand cut structure of the T matrix in the complex s plane. Each arrow leads to a different unphysical sheet.

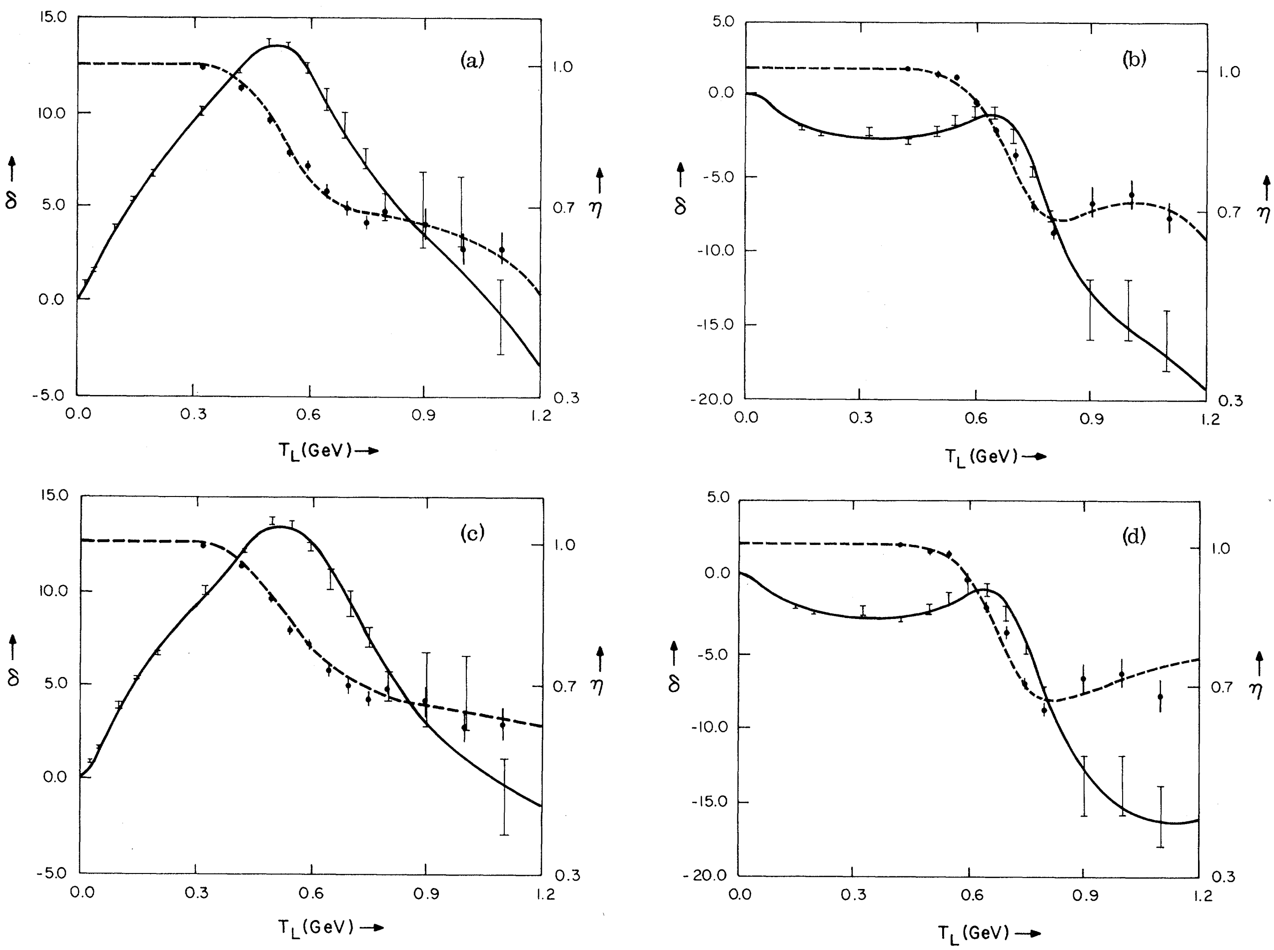

Fig. 2. M-matrix fit to (a) the $I=1$, $^1D_2$ and (b) the $I=1$, $^3F_3$ nucleon-nucleon partial waves of Arndt *et al.* (Ref. 4); K-matrix fit to (c) the $I=1$, $^1D_2$ and (d) the $I=1$, $^3F_3$ nucleon-nucleon partial waves of Arndt *et al.* (Ref. 4). Continuous line is the fit to the phase shift 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-

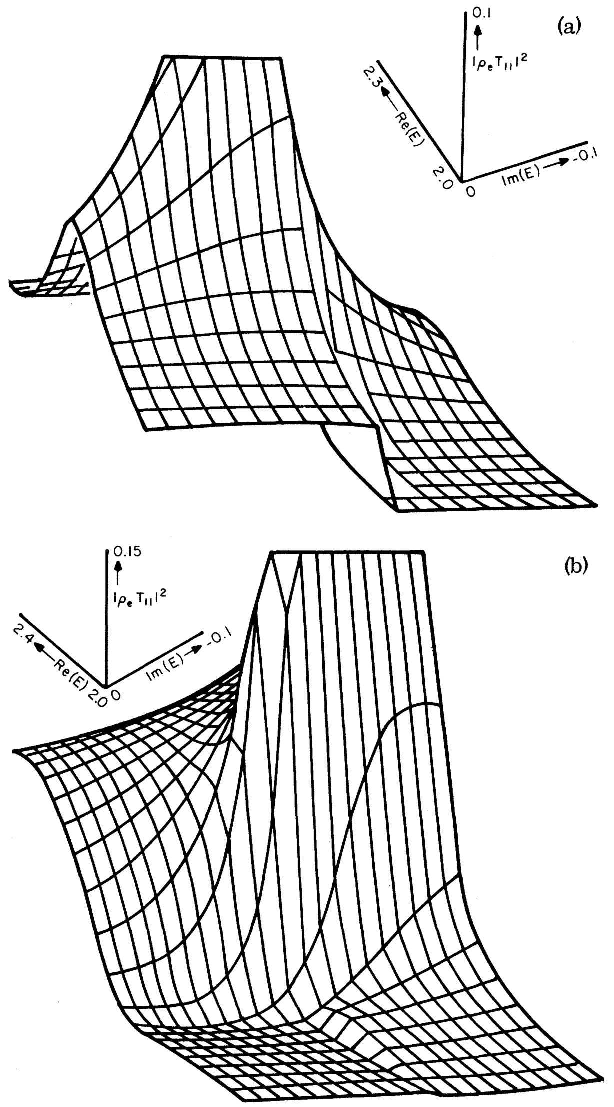

FIG. 3. Three-dimensional plots of $|\rho_e T_{11}|^2$ for (a) the $I = 1$, ${}^1D_2$ and (b) the $I = 1$, ${}^3F_3$ partial waves. Since the $N\Delta$ cut is shown running to the left, $|\rho_e T_{11}|^2$, calculated below the real axis, corresponds to the upper part of the sheet corresponding to arrow 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

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

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.

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