Three-body calculation of the $\Delta\Delta$ dibaryon candidate $D_{03}(2370)$

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The $D_{03}$ dibaryon is generated dynamically as a resonance pole in a $\pi N$+$\Delta$ three-body model, where $\Delta'$ is a stable $\Delta$ baryon. Using separable interactions dominated by the $\Delta(1232)$ isobar for $\pi N$ and by the $D_{12}(1250)$ isobar for $N\Delta'$, with $D_{12}(1250)$ the $N\Delta'$ dibaryon deduced in and constrained by $^1D_2$ $pp$ scattering, the model reduces to an effective two-body problem for $N\Delta'$ which is solved. The mass and width of $D_{03}$ are found close to those of the $I(J^P) = 0(3^+)$ resonance peak observed by WASA@COSY in pion-production $pn$ collisions at 2.37 GeV.

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Introduction. Arguments in favor of $N\Delta$ and $\Delta\Delta$ dibaryon resonances and estimates of their mass values relative to the respective thresholds at 2.17 and 2.46 GeV date back to 1964 [1], as soon as SU(6) symmetry proved useful in classifying baryons and mesons below 2 GeV. Since nucleons and $\Delta$'s belong to the SU(6) representation, product states of $N \times \Delta$ offer numerous nonstrange dibaryon candidates. Focusing on $L = 0$ $s$-wave dibaryons $D_{1S}$, with isospin $I$ and spin $S$, and on the $^{10}_2$ and $^{27}_{10}$ SU(3) multiplets that contain the deuteron $D_{01}$ and $NN$ virtual state $D_{10}$, respectively, these symmetry-based arguments leave only two additional nonstrange dibaryon candidates: $D_{12}$ and $D_{03}$ with predicted masses 2.16 and 2.35 GeV, respectively [1].

Of these two $s$-wave dibaryon candidates, the $D_{12}$ shows up experimentally as an $NN(^1D_2) \leftrightarrow \pi d(^3P_2)$ coupled-channel resonance corresponding to a quasi-bound $N\Delta$ with mass 2.15 GeV, near the $N\Delta$ threshold, and width about 0.115 GeV [2, 3]. Early versions of quark-model based calculations [4, 5, 10–13] of the (real) $D_{12}$ dibaryon resonance, particularly when interpreted as a $\Delta\Delta$ dibaryon, failed, while constituent quarks used in most past works to generate $N\Delta$ and $\Delta\Delta$ interaction potentials for solving a two-body (mostly Schrödinger) wave equation. Given the prominent role of the $\Delta(1232)$ resonance in $\pi N$ dynamics, we extend our hadronic building blocks to include $\Delta$'s as self-consistently as possible. Earlier attempts by Ueda to consider dibaryon resonances within $\pi NN$ and $\pi\pi NN$ dynamics were limited to Heitler-London estimates [15], followed by $\pi NN$ Faddeev calculations with unrealistic nonrelativistic kinematics [16]. In contrast, we solve three-body Faddeev equations with relativistic kinematics, $\pi NN$ for $D_{12}$ and $\pi N\Delta$-like for $D_{03}$, the latter substituting for $\pi\pi NN$ four-body Faddeev-Yakubovsky equations. Each of these derived dibaryon resonances is the lowest, and perhaps the only $s$-wave dibaryon within its own class. Here we focus on $D_{03}$ in the first quantitative phenomenological few-body calculation to confront realistically the recently observed WASA@COSY 2.37 GeV resonance peak [14].

Three-body model of $D_{03}$. In Table I we list two-body thresholds relevant for cluster decompositions of the $I(J^P) = 0(3^+)$ $\pi\pi NN$ system. At 1 fm separation distance, the $\ell = 2$ centrifugal energy upward shifts of at least 200 MeV make the channels $(NN)_d-(\pi\pi)_s(500)$ and $N-(\pi\pi)_sN^*(1440)$ competitive against the $\Delta\Delta$ channel with $\ell = 0$ threshold at 2460 MeV. For $\pi \rightarrow (\pi NN)_{D_{12}(2150)}$, the $\Delta$-dominated $\pi N$ interaction should reduce the $\ell = 1$ upward shift from 300 MeV (no interaction) to about 150 MeV, so that the effective threshold here becomes somewhat lower than the $\ell = 0 \Delta\Delta$ threshold. This singles out $\Delta(1232)$ and $D_{12}(2150)$ as the most likely fermionic degrees of freedom of which pions may benefit in forming the $I(J^P) = 0(3^+) \pi\pi NN$...
system, with \(D_{03}\) likely to emerge as a quasibound state in \((\Delta\Delta)_{\text{upper}}-\pi(1236)\) \(\to\pi\pi\) \(N\Delta\) coupled-channel calculations. This is consistent with the observation that \(D_{03} \to D_{12} + \pi\) provides a dominant doorway mode into two-pion final states \[17\]. We also note that by assigning \(s\)-wave \(N\Delta\) structure to \(D_{12}\), both angular-momentum and isospin recoupling coefficients for transforming \(\pi D_{12}(2150)\) with \(p\)-wave pion in a \(I(J') = 0(3^+)\) state into \(s\)-wave \(\Delta\Delta\) are equal to 1, thus maximizing the coupling between these channels.

These arguments led us to reduce the \(I(J') = 0(3^+)\) \(\pi\pi NN\) system to a system of three hadrons \(\pi, N, \Delta'\) interacting via pairwise separable potentials. This approximates one of the \(pN\) pre-existing pairs in the four-body system by a stable \(\Delta\) of mass 1232 MeV and zero width, here denoted \(\Delta'\). In this model the \(\pi N\) interaction is limited to the \(I(J') = \frac{3}{2}(\frac{3}{2}^+)\) \(p\)-wave channel (denoted \(P_{33}\) dominated by the \(\Delta\) resonance. The \(N\Delta'\) interaction is limited to the \(I(J') = 1(2^+)\) \(s\)-wave channel dominated by the \(D_{12}\) dibaryon and, finally, the \(\pi\Delta'\) interaction is neglected because of the mass of the lightest \(N^*(\frac{3}{2}^+)\) isobar candidate \(N^*(1680)\) is too high for our purpose. This \(\pi N\Delta'\) three-body model leads to a \(\Delta\Delta'\) eigenvalue problem for the \(T\) matrix diagram shown in Fig. \[4\] where starting with \(\Delta\Delta'\), the \(\Delta\) resonance isobar decays into a \(\pi N\) pair followed by \(N\Delta' \to N\Delta'\) scattering via the \(D_{12}\) isobar (marked \(D\) in the figure) with a spectator pion, and finally by \(\pi N \to \Delta\) fusion back into the \(\Delta\Delta'\) channel.

![Graphical representation of the \(\Delta\Delta'\) dibaryon \(T\)-matrix pole equation. \(D\) denotes the \(D_{12}\) isobar.](image)

We note that, whereas the \(\Delta\)-isobar decay to the \(\pi N\) channel is fully accounted for, the \(D_{12}\) isobar is allowed to decay only to \(N\Delta'\). Additional width contribution will arise upon allowing the stable \(\Delta'\) in the \(\pi N\Delta'\) model to acquire normal \(\Delta \to N\pi\) decay width. However, this added \(\Delta'\) width is partly suppressed by quantum-statistics correlations between the decay \(N\pi\) pair and the pre-existing \(N\pi\) pair. Thus, for \(s\)-wave nucleons and \(p\)-wave pions, \(J' = 3^+\) implies space-spin symmetry for nucleons as well as for pions. With total \(I = 0\), Fermi-Dirac (Bose-Einstein) statistics for nucleons (pions) allows for isospins \(I_{NN} = I\pi = 0\), forbidding \(I_{NN} = I\pi = 1\), with weights 2/3 and 1/3, respectively, obtained by recoupling the two isospins \(I_{NN} = 3/2\) in the \(I = 0 \Delta\Delta'\) state.

We now specify the pairwise interactions in the \(\pi N\Delta'\) three-body model. The interaction between particles 1,2 is denoted \(V_{12}\), etc. Since the \(\pi N\) interaction \(V_{33}\) is dominated by the \(\Delta\) resonance, it is limited here to a \(P_{33}\) rank-one separable-potential of the form

\[V_{33}(p_3, p_3') = \lambda_{33} g_3(p_3) g_3(p_3'),\]

so that the corresponding \(T\) matrix is given by

\[t_3(\omega_3; p_3, p_3') = g_3(p_3) \tau_3(\omega_3) g_3(p_3'),\]

where \(\tau_3(\omega_3)\) is the propagator of the \((\Delta 1232)\) isobar in the two-body cm system, with \(\omega_3\) the two-body \(\pi N\) cm energy. In the three-body cm system, with \(W\) the total three-body cm energy and \(q_3\) the momentum of the spectator \(\Delta'\) with respect to the two-body \(\pi N\) isobar, it assumes the form

\[T_{3}^{-1}(W; q_3) = \lambda_{3}^{-1} \int_{0}^{\infty} \frac{[g_3(p_3)]^2 p_3^2 dp_3}{W - E_{\Delta}(q_3) - E_3(p_3, q_3) + i\epsilon},\]

where \(E_3(p_3, q_3) = [(E_N(p_3) + E_\pi(p_3))]^2 + q_3^2\), with \(E_h(p) = (m_h^2 + p^2)^{1/2}\) for hadron \(h\). For \(q_3 = 0\), when the two-body cm system degenerates to the two-body cm system, \(T\) and \(\tau_3\) are related by a simple mass shift, \(T_3(W; q_3 = 0) = \tau_3(W - m_\Delta')\), as expected.

For the \(\pi N\) \(p\)-wave form factor \(g_3\) we considered two forms, labeled here \(k=1,2\):

\[g_3(p_3) = p_3 \exp(-p_3^2/\beta^2) + C p_3^3 \exp(-p_3^2/\gamma^2),\]

falling off exponentially \[18\], and

\[g_3(p_3) = \frac{p_3}{(1 + p_3^2/\beta^2)^2} + C \frac{p_3^3}{(1 + p_3^2/\gamma^2)^3},\]

falling off as inverse power of \(p_3\), here as \(p_3^{-3}\). Both form factors, using the parameters listed in Table \[11\] reproduce perfectly the \(\pi N\) \(P_{33}\) phase shifts from Ref. \[19\]. The table also lists the distance \(r_0\) at which the Fourier transform \(g_3(r)\) flips sign, which roughly represents the spatial extension of the \(P_{33}\) \(p\)-wave form factor \[18\]. Together with a rank-two separable \(NN\) potential reproducing the \(^3S_1\) phase shift, these form factors lead in a relativistic \(\pi N\) Faddeev calculation to \(D_{12}\) dibaryon pole at \(E = M - i (\Gamma/2) = 2151(2) - i 60(3)\) MeV \[7\], in good agreement with accepted values \[2\].

### Table I: Two-body threshold energies \(E_{th}\) (in MeV) and lowest partial waves \(\ell\) for the \(I(J') = 0(3^+)\) \(\pi\pi NN\) system.

| \(E_{th}\) | \(\Delta - \Delta\) | \(d - \sigma(500)\) | \(N - N^*(1440)\) | \(\pi - D_{12}(2150)\) |
|---|---|---|---|---|
| 2380 | 2380 | 2380 | 2290 |
| \(\ell\) | 0 | 2 | 2 | 1 |

### Table II: \(\pi N\) \(P_{33}\) form factor \(g_3(p)\) parameters [Eqs. \[4\], \[5\], \[6\]], and the zero \(r_0\) of the Fourier transform \(g_3(r)\) [18].

| \(k\) | \(\lambda_3\) (fm\(^{-1}\)) | \(\beta\) (fm\(^{-1}\)) | \(\gamma\) (fm\(^{-1}\)) | \(C\) (fm\(^{-2}\)) | \(r_0\) (fm) |
|---|---|---|---|---|---|
| 1 | -0.07587 | 1.04 | 2.367 | 0.23 | 1.36 |
| 2 | -0.04177 | 1.46 | 4.102 | 0.11 | 0.91 |
The $N\Delta'$ interaction $V_1$, dominated by the $D_{12}(2150)$ isobar resonance, is limited here to the $I(J^P) = 1(2^+)$ channel. $D_{12}$ shows up as an inelastic resonance in the $^1D_2 NN$ partial-wave above the $\pi NN$ threshold $[2, 20]$. To generate the necessary inelastic $\pi NN$ cut we introduced a third s-wave subchannel $NN'$ where $N'$ is a fictitious nonstrange stable baryon with $I(J^P) = \frac{1}{2}(\frac{3}{2}^+)$ and $m_{N'} = m_N + m_\pi$. We have then fitted the $^1D_2$ partial-wave amplitude of Arndt et al. $[20]$ using a coupled-channel separable potential

$$V_{1}^{mn}(p_1, p'_1) = \lambda_1 g_1^m(p_1)g_1^n(p'_1), \quad (m, n = 1 - 3),$$

where the three subchannels labeled $m, n$ are $1=NN$ (d-wave), $2=NN'$ (s-wave), and $3=NN'$ (s-wave). The $t$-matrix of the system is obtained by solving a Lippmann-Schwinger equation with relativistic kinematics which in the case of the separable potential $[0]$ has the solution

$$t_{1}^{mn}(\omega_1; p_1, p'_1) = g_1^m(p_1)\tau_1(\omega_1)g_1^n(p'_1),$$

where $\tau_1(\omega_1)$ is the propagator of the $D_{12}$ isobar in the two-body cm system. In the three-body cm system it assumes the form

$$T_{1}^{1}(W; q_1) = \lambda_1^{-1} - \frac{3}{r=1} \int_0^\infty \frac{[g_1^m(p_1)]^2 p_1^2 dp_1}{W - \epsilon_{r}(p_1)} + i\epsilon,$$

where $\epsilon_{r}(p_1, q_1) = [(E_N(p_1) + E_r(p_1))^2 + q_1^2]^\frac{1}{2}$, with masses $m_r = (m_N, m_N, m_{\Delta'})$ for $r = (1, 2, 3)$.

The form factors of the separable potential $[0]$ were taken in the following form:

$$g_1^n(p_1) = \frac{p_1^\ell}{[1 + p_1^2/(\alpha_n)^2]^{1+\frac{\ell}{2}}} \left(1 + A_n \frac{p_1^2}{1 + p_1^2/(\alpha_n)^2}\right),$$

where $\ell = 2$ for $n = 1$, and $\ell = 0$ for $n = 2, 3$. The range parameters $\alpha_n$ were limited to values $\alpha_n \lesssim 3$ fm$^{-1}$ to ensure that the physics of these coupled channels does not require explicit shorter-range degrees of freedom, for example $\pi N \rightarrow \rho N$. Good fits to the $NN^1D_2$ scattering parameters satisfying this limitation required that not all $A_n$ be zero. A comparison between the phase shift $\delta$ and inelasticity $\eta$ from our best fit and those derived from pp scattering analyses $[20]$ is shown in Figs. 2 and 3. Here, $\delta$ and $\eta$ are given in terms of $S$ and $T$ matrices by $S = 1 + 2iT = \eta \exp(i\delta)$. A variance of 0.02 was used for Re $T$ and Im $T$ in these fits. We note that the decrease of the inelasticity $\eta$ from a value 1 is due to the $r = 2 NN'$ subchannel which generates the inelastic cut starting at the $\pi NN$ threshold, and that no explicit $D_{12}$ pole term was introduced in the $r = 3 NN'$ subchannel.

**Results and discussion.** Applying standard three-body techniques $[18]$ to the Faddeev equations of our $\pi N\Delta'$ three-body model, the following homogeneous integral equation for $T$-matrix poles is obtained:

$$T_3(W; q_3) = T_3(W; q_3) \int_0^\infty dq'_3 M(W; q_3, q'_3)T_3(W; q'_3),$$

where, suppressing the dependence on $W$,

$$M(q_3, q'_3) = 2 \int_0^\infty dq_1 K_{31}(q_3, q_1)T_1(q_1)K_{13}(q_1, q'_3),$$

$$K_{31}(W; q_3, q_1) = \frac{1}{2} q_3 q_1 \int_{-1}^1 d\cos\theta g_3(p_3)g_1^{N\Delta'}(p_1) \frac{p_3 \cdot q_1}{W - E_1(q_1) - E_2(q_1, q'_3) - E_3(q_3) + i\epsilon},$$

with $K_{13}(W; q_1, q_3) = K_{31}(W; q_3, q_1)$. The factor 2 on the rhs of Eq. (11) takes into account that the decay
The integral equation (11) coincides with that depicted in Fig. III for the \( \Delta' \) eigenvalue problem. In practice, replacing \( A_3(W; q_i) \) by \( Z_3(W; q_i) = A_3(W; q_i)/T_3(W; q_i) \), Eq. (11) was transformed to a standard eigenvalue equation

\[
Z_3(W; q_i) = \int_0^\infty dq_3' M(W; q_3, q_3') \left( \frac{q_3}{q_3'} \right)^{1/2} Z_3(W; q_3')
\]

which was solved numerically. In order to search for resonance poles of (10), the integral equation was extended into the complex plane using the standard procedure \( q_i \rightarrow q_i \exp(-i\phi) \) [21] which opens large sections of the unphysical Riemann sheet so that one can search for eigenvalues of the form \( W = M - i(\Gamma/2) \).

**TABLE III:** Lowest \( \chi^2/N \) values of \( NN \) \( ^1D_2 \) fits, fitted range parameters \( \alpha_{i,2,3} \) (in fm\(^{-1}\)) of the \( NN-NN'-NN' \) form factors [9] upon fixing \( A_{1,2,3} \) (in fm\(^2\)), and Faddeev-calculated \( D_{03} \) pole positions \( W_k \) (in MeV) for \( \pi N \) form factors labeled \( k \) in Table II

| \( A_{i,2,3} \) | \( \chi^2/N \) | \( \alpha_1 \) | \( \alpha_2 \) | \( \alpha_3 \) | \( W_1 \) | \( W_2 \) |
|-----|-----|-----|-----|-----|-----|-----|
| 1,1,1 | 0.78 | 1.47 | 2.27 | 3.24 | 2383 - i 41 | 2343 - i 24 |
| 0,1,1 | 1.10 | 2.00 | 2.11 | 2.96 | 2384 - i 44 | 2356 - i 30 |
| 0,1,0 | 1.15 | 2.04 | 2.16 | 2.44 | 2392 - i 52 | 2380 - i 45 |

In the actual solution of the eigenvalue equation (13) we replaced the Breit-Wigner mass value \( m_{\Delta'} = 1232 \) MeV in the propagator \( T_3 \), Eq. (9), by an effective \( \Delta \)-pole complex mass value [19] \( m_{\Delta} = 1211 - i \left( 2/3 \right) \times 49.5 \) MeV, where the 2/3 suppression factor accounts for the \( \Delta \) decay phase space effective in the \( I(JF' = 0(3^+) \) \( \pi N \Delta \) three-body system, as discussed earlier. \( D_{03} \) resonance pole positions calculated using the three lowest \( \chi^2 \) fits of the \( NN-NN'-NN' \) form factors [9] to the \( ^1D_2 \) \( NN \) scattering parameters are listed in Table III for the two \( \pi N \) form factors specified in Table II. Comparing the calculated mass values \( W \) in the last column to those in the preceding column, one observes sensitivity to the spatial extension of the \( \pi N \) form factor, quantified by the values of \( r_0 \) from Table III which we consider as providing reasonable bounds on this spatial extension. The smaller \( r_0 \), the lower the calculated mass values are. Admitting values of \( r_0 \) appreciably below 0.9 fm requires the introduction of explicit vector-meson and/or quark-gluon degrees of freedom. Within a given column for \( W \) in the table, the calculated mass values display sensitivity to the \( N\Delta' \) form factor primarily through the fitted values of \( \alpha_3 \) listed here. The larger \( \alpha_3 \), the lower the calculated mass values are. In general, it was found impossible to get values of \( \alpha_3 \lesssim 2.5 \) fm\(^{-1}\), whereas going beyond \( \alpha_3 \sim 3 \) fm\(^{-1}\) was considered undesirable, again requiring the introduction of explicit short-range degrees of freedom. As for width values \(-2 \) Im \( W \), the calculated widths display little sensitivity to these form factors and the widths are determined mostly by the phase space available for decay. Averaging over the pole positions of the best-fit solutions in the first row of Table III the mass and width values for \( D_{03} \) are

\[
D_{03} : \quad M = 2363 \pm 20, \quad \Gamma = 65 \pm 17 \text{ (MeV)}. \quad (14)
\]

By scaling to the mass-width phenomenology of the known \( \Delta \) resonance [19], the corresponding \( D_{03} \) Breit-Wigner mass-width values for comparison with experiment [12] should each be about 10 MeV higher.

In summary, considering the \( \Delta \) resonance as a stable baryon \( \Delta' \), we have presented a dynamical \( \pi N \Delta' \) separable-potential model for the \( D_{03} \) dibaryon that captures the essential physics of the underlying pairwise interactions, using fitted form factors for the \( \pi N \) and the \( s \)-wave \( N\Delta' \) interactions in the channels dominated by the \( \Delta(1232) \) baryon resonance and the \( D_{12}(2150) \) dibaryon resonance, respectively. The corresponding three-body Faddeev equations were derived and the resulting effective two-body \( \Delta' \) eigenvalue equation of Fig. III was solved, replacing the \( \Delta' \)-spectator real mass in the in-medium \( \pi N \) propagator by a physical \( \Delta \) effective complex mass value. A robust \( D_{03} \) dibaryon pole was found above the \( \pi D_{12} \) threshold but below the \( \Delta \Delta \) threshold, with mass value \( M = (2.36 \pm 0.02) \) GeV in good agreement with the location of the \( pn \to d\pi\pi \) resonance observed by WASA@COSY. Furthermore, the calculated width of \( \Gamma = 65 \pm 17 \) MeV also agrees well with the observed width of 70 MeV.

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