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Distorted-Wave Born-Approximation Analysis of ^{36,38}Ar(d,p) to Neutron Resonances in ^{37,39}Ar

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^{36,38}Ar (d,p) angular distributions at 9- and 10-MeV incident energy, for 13 neutron-unbound states in ^{37,39}Ar, are described in terms of conventional distorted-wave Born-approximation theory using complex-energy eigenstates as form factors, and spectroscopic information is extracted which is consistent with previous studies of the bound neutron states of ^{37,39}Ar.

I. INTRODUCTION

In studies of the reaction ^{36}Ar(d,p)^{37}Ar at 9.16 MeV,[1] and ^{38}Ar(d,p)^{39}Ar at 10.06 MeV[2] it has been found that a number of neutron unbound states in ^{37}Ar and ^{39}Ar are populated. Sen *et al.*[1,2] were able to obtain angular distributions for these states in 5° steps over c.m. angles from 26 to 146°; specifically, data are available for states at 8.89 and 9.01 MeV in ^{37}Ar, and at 6.79, 7.00, 7.06, 7.14, 7.22, 7.34, 7.40, 7.50, 7.56, 7.63, and 7.73 MeV in ^{39}Ar.

Presented here are distorted-wave Born-approximation (DWBA) analyses of these 13 angular distributions, making use of complex-energy eigenstates to described the resonance states. It is shown that use of complex-energy eigenstates permits extraction of spectroscopic factors consistent with the usual bound-state single-particle spectroscopic factors.

II. COMPLEX-ENERGY EIGENSTATES

The complex-energy eigenstate, often called a Gamow state, is either of two equivalent facto. making up the residue of the Green's function of the system, at the pole corresponding to a given resonance.[3-5] Thus it is closely analogous to the bound-state function, which is again either factor of the residue of the Green's function of the system at the pole corresponding to a given bound state. It is straightforward to show that complex-energy eigenfunctions have normalization and orthogonality properties analogous to bound states,[6] and can form part of a basis for eigenstate expansion, in the sense that a quantum mechanical state Ψ can, under weak restrictions, be expanded as a sum over discrete bound and Gamow states, plus a contour integral over continuum states.[7] The choice of contour determines the number of Gamow states included in the discrete sum, and also the set of functions Ψ which may be so expanded. The situation is quite reminiscent of Regge-type representations of the scattering amplitude, in many ways.

Some confusion has resulted over the connection of Gamow states to the familiar Hilbert spaces of scattering theory. Berggren[8] has shown that norms can be introduced for Gamow states such that many analogous mathematical properties, concepts, and methods are indeed shared by bound states and complex-energy eigenstates. The application of sum rules to Gamow states remains a delicate question, since the contribution of the contour integral over continuum states has not been shown to be positive. However, if most of the single-particle strength resides in a few low-lying bound states and resonances, as is the case in the present analysis, there would seem to be no difficulty in practice.

Use of Gamow states presents a second difficulty stemming from their asymptotic behavior. A recent review by Huby[9] gives a succinct discussion of the available remedies for this and other divergences in scattering theory. In the work presented here, we use the so-called Abel method of regularization of the integrals involving Gamow states,[6-9] as discussed later.

Complex-energy eigenstates were obtained, for the analysis reported here, using a greatly modified version of the program NEP,[10] to solve the differential equation

$$[(-\hbar^2/2\mu)(d^2/dr^2 - l(l+1)/r^2) - \tilde{E} + U_{lj}(r)]g_{lj}(r) = 0\,, \tag{1}$$

where $U_{lj}(r)$ consists of a real Woods-Saxon potential plus a Thomas-type spin-orbit potential, and the c.m. energy $\tilde{E}$ is complex. The boundary conditions imposed on g_{lj} are that

$$g_{lj}(0) = 0 \tag{2a}$$

and

$$g_{lj}(r) \underset{r\to\infty}{\longrightarrow} [2\mu(-E_I)/\hbar^2\tilde{k}]^{1/2}[G_l(\tilde{k}, r) + iF_l(\tilde{k}, r)]\,. \tag{2b}$$

Here E_I is the imaginary part of $\tilde{E}$, $\tilde{k}$ is the wave number corresponding to $\tilde{E}$, and F_l (G_l) is the regular (irregular) Coulomb function for complex $\tilde{E}$.[11] Note that in our specific case the Coulomb parameter $\eta = 0$. Boundary condition (2b) is satisfied by an automatic search on $\tilde{E}$ for a given U_{lj}.[10] The complex-energy eigenstate so constructed uniquely corresponds to the single-particle resonance whose pole occurs in the single-particle Green's function at energy $\tilde{E}$.

To see more clearly how the complex-energy eigenstate is used in DWBA, consider the process $A(d,p)B^* \to A+n$, in which a low-lying neutron resonance is populated by a deuteron stripping reaction. The cross section for the process which puts a neutron into the continuum at an energy between E_n and $E_n + dE_n$, when the proton is detected within solid angle element $d\Omega_p$, is

$$\frac{d^2\sigma}{d\Omega_p dE_n} = \frac{2\mu_{pB}\mu_{dA}\mu_{nA}}{\pi^2\hbar^6 k_d{}^3 k_p k_n}\left(\frac{B}{A}\right)^2\left(\frac{2J_B+1}{2J_A+1}\right) \times \sum_{jmm_dm_p}\Big|\sum_{ls} C_{lsj}\bar{\beta}_{lsj}^{mm_dm_p}\Big|^2\,, \tag{3}$$

where $\bar{\beta}_{lsj}^{mm_dm_p}$ is defined as in Refs. 12–14, and contains the radial overlap integrals

$$f_{L_pJ_pL_dJ_d}^{lsj} = \int \chi_{L_pJ_p}(Ak_pr/B)[\chi_{lj}(k_nr)/r] \times \chi_{L_dJ_d}(k_dr)dr\,. \tag{4}$$

Here the form factor is $\chi_{lj}(k_nr)$, the radial state function for the continuum neutron, which for $E_n \approx E_R$ has the form

$$\chi_{lj}(kr) \approx -\frac{e^{2i\zeta_l^j}(\frac{1}{2}\Gamma_l^j)}{E_n - E_R + i\frac{1}{2}\Gamma_l^j}\,O_l(k_nr) + \chi_{lj}^{NR}(k_nr)\,. \tag{5}$$

Here ζ_l^j is the "background" phase shift and $O_l(k_nr)$ is a purely outgoing wave, while χ_{lj}^{NR} is

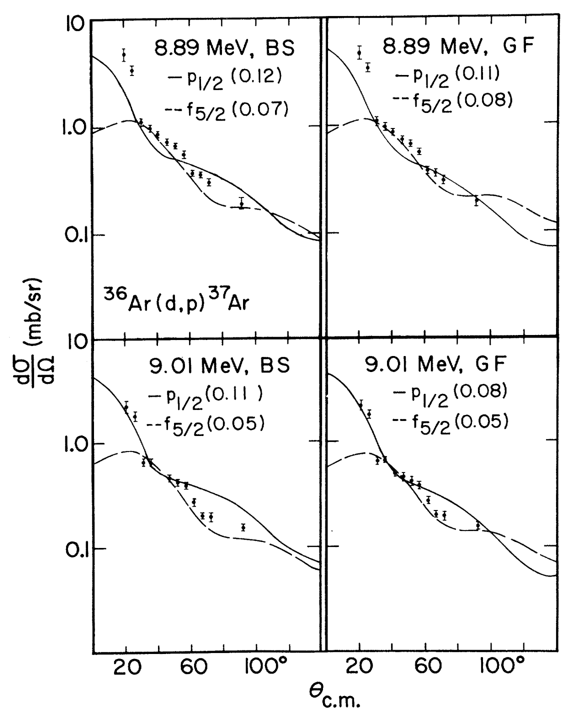

FIG. 1. The $^{36}Ar(d,p)^{37}Ar$ cross sections at 9.16-MeV center-of-target energy, from Ref. 1, for the neutron-unbound states at 8.89 and 9.01 MeV in excitation. On the left are results of DWBA calculations using the optical potentials of Ref. 1 and conventional form factors, for neutrons bound by 0.1 MeV. The solid curve is for $p_{1/2}$ transfer, the dashed curve for $f_{5/2}$. Spectroscopic factors are given in parentheses. On the right are results of DWBA calculations for the same states, using Eq. (8) and the Gamow functions for $p_{1/2}$ and $f_{5/2}$ resonances (see Table).

the nonresonant part of χ_{lj}.

The discrete proton groups corresponding to neutron-resonance population, observed in $A(d,p)B^* \rightarrow A+n$, sit on a background which consists in part of a proton continuum from three-body breakup, $A(d,np)A$. This process would be described using χ_{lj}^{NR} as the form factor and evaluating Eq. (3) directly (see, e.g., Ref. 12), but it is not of concern at present unless it interferes strongly with $A(d,p)B^* \rightarrow A+n$. Near resonance then, one simply omits χ_{lj}^{NR}.

Since the energy dependence of $O_l(k_n r)$, $\chi_{L_p J_p}(k_p r)$, and $\chi_{L_d J_d}(k_d r)$ is negligible within the small interval $\pm\Gamma_l^j$ around E_R, for an observable single-particle resonance, one can factor the strong energy dependence of the first term of Eq. (5) out of Eq. (3) and do the E_n integration analytically. One obtains

$$\frac{d\sigma}{d\Omega_p}=\frac{2\mu_{dA}\mu_{pB}\mu_{nA}}{\pi^2\hbar^6}\,\frac{1}{k_d{}^3k_pk_n}\left(\tfrac{1}{2}\pi\Gamma_l^j\right)\left(\frac{B}{A}\right)^2\left(\frac{2J_B+1}{2J_A+1}\right)\times\sum_{jmm_dm_p}\Big|\sum_{ls}C_{lsj}\tilde{\beta}_{lsj}^{mm_dm_p}\Big|^2, \qquad (6)$$

where $\tilde{\beta}_{lsj}^{mm_dm_p}$ contains the integrals

$$\int \chi_{L_pJ_p}(Ak_pr/B)[O_{lj}(k_nr)/r]\,\chi_{L_dJ_d}(k_dr)dr\,. \qquad (7)$$

Taking the factor $(\mu_{nA}\Gamma_l^j/\hbar^2k_n)$ inside the integral now gives a cross section identical in form to that given in the usual two-body DWBA,[15] and also permits one to continue the expression for $\tilde{\beta}$ to complex $\tilde{k}_n$. Then

$$\frac{d\sigma}{d\Omega_p}=\frac{\mu_{dA}\mu_{pB}}{\pi\hbar^4}\,\frac{1}{k_d{}^3k_p}\left(\frac{B}{A}\right)^2\left(\frac{2J_B+1}{2J_A+1}\right)\times\sum_{jmm_dm_p}\Big|\sum_{ls}C_{lsj}\beta_{lsj}^{mm_dm_p}\Big|^2, \qquad (8)$$

where β_{lsj} contains the overlap integrals

$$\int \chi_{L_pJ_p}(Ak_pr/B)[g_{lj}^R(\tilde{k}_n^R,r)/r]\,\chi_{L_dJ_d}(k_dr)dr \qquad (9)$$

and is thus identical to the β defined by Satchler[15] with g_{lj}^R/r as form factor. In performing the integral (9) with Gamow functions, a Gaussian weighting factor $\exp(-\alpha r^2)$ is inserted and the limiting value, as α goes to zero, used in Eq. (8).[6-9, 16]

The limit is obtained by a quadratic interpolation. If $\alpha_n = n\alpha_1 (n=1,2,3)$, the limiting value of the overlap integral, $I(0)$, is simply taken as $3[I(\alpha_1)-I(\alpha_2)]+I(\alpha_3)$. The advantage of this procedure is its simplicity and speed. Thus by repeating the calculation for smaller and smaller starting values of α, the reliability of the convergence is readily studied. The starting value for α_1 in the calculations reported here is 0.005 fm^{-2}, or more generally, $\alpha_1 = 10/R_c{}^2$, where R_c is the outer radial cutoff of the integral (9). For such small α one might expect the convergence to be very slow,[9] but in practice it is found that

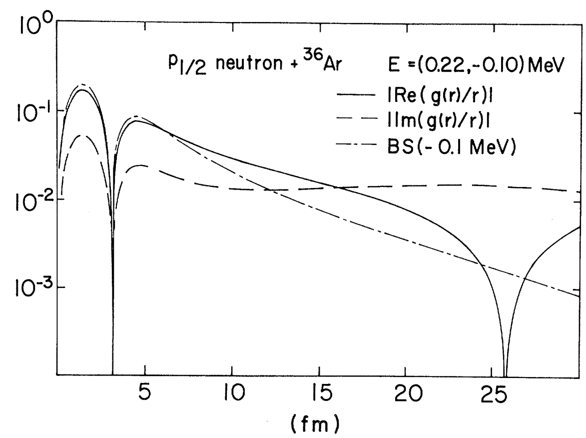

FIG. 2. A comparison of complex energy eigenstate and bound-state functions for the $p_{1/2}$ neutron plus ^{36}Ar system. The dot-dash curve is the modulus of the radial state function of a $2p_{1/2}$ neutron bound by 0.1 MeV, while the solid and dashed curves show the moduli of the real and imaginary parts, respectively, of the Gamow function corresponding to the $p_{1/2}$ neutron resonance at (0.22, −0.10) MeV. See also the lower half of Fig. 1. Note the close similarity of the bound-state function and Gamow state function within the nucleus, in shape and magnitude. Potential parameters are as given in the text and Table I.

TABLE I. Complex-energy eigenstates and results of analysis.

System	l_j	E_x	Real ($\tilde{E}$)	Imaginary ($\tilde{E}$)	V_0	S_{dp}
$n+{}^{36}$Ar	$p_{1/2}$	8.89	0.10	−0.029	49.96	0.11
	$f_{5/2}$		0.10	-2.3×10^{-6}	57.80	0.08
	$p_{1/2}$	9.01	0.22	−0.10	49.06	0.08
	$f_{5/2}$		0.22	-3.1×10^{-5}	57.57	0.05
$n+{}^{38}$Ar	$p_{1/2}$	6.79	0.20	-9.0×10^{-2}	47.55	0.18
	$f_{5/2}$		0.20	-2.3×10^{-5}	55.48	0.12
	$p_{1/2}$	7.00	0.41	−0.31	45.70	0.028
	$f_{5/2}$		0.41	-2.7×10^{-4}	55.07	0.023
	$p_{1/2}$	7.06	0.47	−0.42	44.96	0.017
	$f_{5/2}$		0.48	-4.7×10^{-4}	54.93	0.021
	$p_{1/2}$	7.14	0.54	−0.59	43.96	0.004
	$f_{5/2}$		0.55	-7.6×10^{-4}	54.78	0.016
	$f_{5/2}$	7.22	0.63	-1.2×10^{-3}	54.62	0.023
	$f_{5/2}$	7.34	0.75	-2.1×10^{-3}	54.38	0.025
	$f_{5/2}$	7.40	0.81	-2.8×10^{-3}	54.25	0.019
	$f_{5/2}$	7.50	0.91	-4.0×10^{-3}	54.06	0.017
	$f_{5/2}$	7.56	0.97	-5.0×10^{-3}	53.93	0.011
	$f_{5/2}$	7.63	1.038	-6.23×10^{-3}	53.79	0.025
	$f_{5/2}$	7.73	1.15	-8.77×10^{-3}	53.55	0.029

quite stable results are obtained for practical values of R_c (say, 40–100 fm) and practical numbers of partial waves (say, 30–50). See also Ref. 16.

The cross section resulting from such a limiting process is indeed found to be stable to within a few percent against wide variations in numerical methods, parameters, and procedures. Note that C_{lsj} has exactly the same significance as for bound states,[15] as should be clear from Eq. (8). Thus, spectroscopic factors are extracted, as usual, by $S_{dp} = d\sigma_{\text{exp}}/d\sigma_{\text{DWBA}}$.

The genuine divergence of Eq. (9) requires some effective regularization method. For example, in the case of the ^{36}Ar$(d,p)^{37}$Ar ($p_{1/2}$, 9.01-MeV) transition, if a fixed value of $R_c = 40$ fm is used and no convergence factor is inserted, the cross section at forward angles is in error by 50% and unrealistic oscillations are observed in the angular distribution. The pathology is of course not cured by increasing R_c or the number of partial waves, although large enough R_c and enough partial waves would guarantee convergence of Eq. (7).

At this point a legitimate question may be raised. The radial integral, Eq. (7), as it stands, is convergent.[9] Basically the convergence of the radial integral over three distorted waves is a consequence of the partial wave expansions, and interchange of summation and integration, which are crucial to the conventional derivations of the two- or three-body DWBA cross section.[12,15] Continuation of Eq. (7) to complex $\tilde{k}_n$ renders it a divergent integral, evaluation of which necessitates the use of some forceful regularization technique. Hence, why use Gamow states at all?

Restricting the discussion to single-particle potential resonances, as we do here, clarifies the question somewhat. Since the scattering problem can be solved "exactly," the apparatus of, say, shell-model reaction theory is unnecessary.[4] Hence the choice is between use of the full, resonating distorted wave [Eq. (5)] and the Gamow state $g_{lj}(r)$. But $\chi_{lj}(k_n r)$, evaluated at the resonance energy, is *not* the state function of the resonance in a quantum-mechanical sense. A remedy, which is time-consuming, is to compute $\chi_{lj}(k_n r)$ at a number of energies in the vicinity of the resonance and perform the E_n integration leading to Eq. (6) numerically. Even further, the slow convergence of Eq. (7) requires a time-consuming, careful numerical radial integration. Finally, one will not directly obtain the single-particle width of the resonance.

By going to the resonance pole itself one renders the radial integral strictly divergent but, assuming a speedy remedy for the divergence,

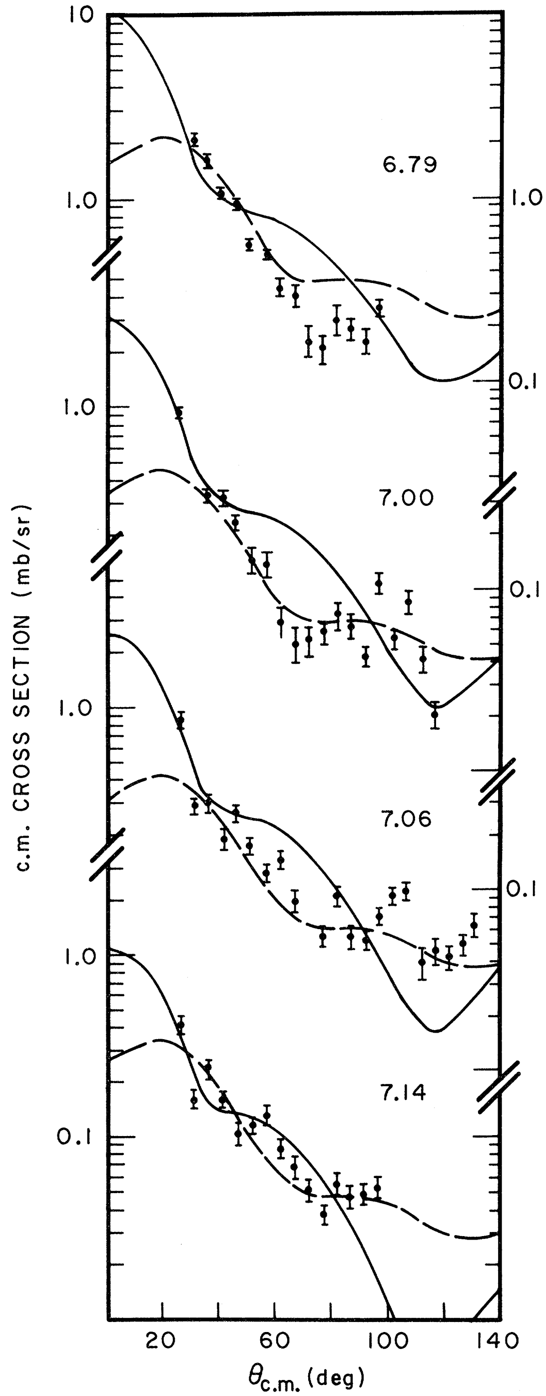

FIG. 3. The ^{38}Ar$(d,p)^{39}$Ar cross sections at 10.06-MeV center-of-target energy, from Ref. 2, for the four neutron-unbound states from 6.79 to 7.14 MeV in excitation. The solid curves are results of DWBA calculations using Gamow functions, as explained in the text, assuming $p_{1/2}$ resonances. The dashed curves are for $f_{5/2}$ resonances. Optical potentials are those of Ref. 2. Spectroscopic factors and resonance parameters are given in the Table.

the saving in time and labor is remarkable. The Gamow state is as fast and easy to compute as the conventional bound-state form factor, and so is the resulting DWBA cross section. Thus, conventional DWBA programs may be used, with straightforward modification. We would claim no more than this for the approach suggested here. It is of course also intuitively satisfying to describe the decaying resonant state by a nonstationary, decaying form factor, which the complex-energy eigenstate provides.

III. ANALYSIS AND SUMMARY

In Fig. 1 are shown DWBA fits to the ^{36}Ar(d,p) cross sections, at 9.16-MeV deuteron energy,

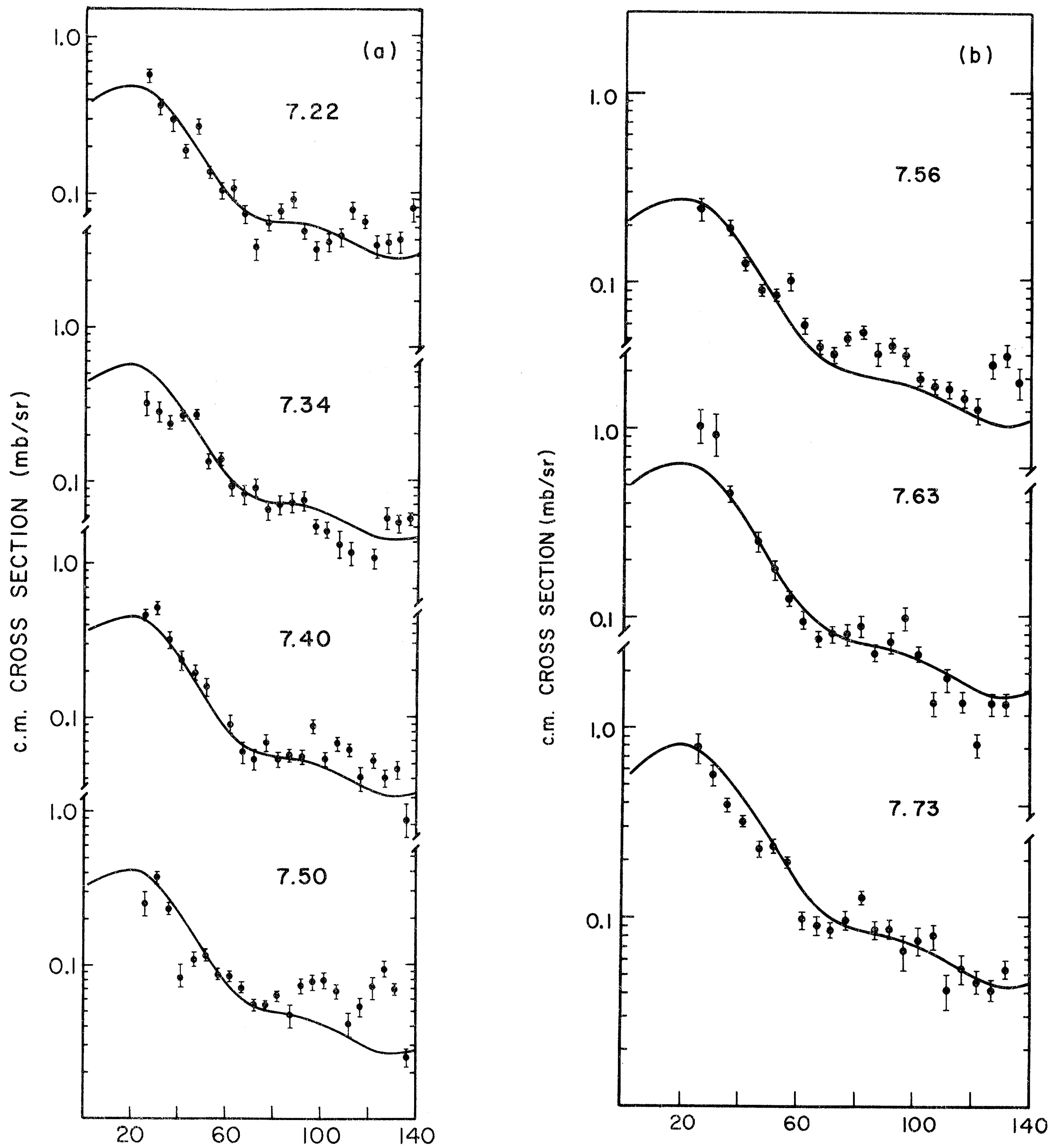

FIG. 4. (a) The ^{38}Ar$(d,p)^{39}$Ar cross sections at 10.06 MeV, from Ref. 2, for the four neutron unbound states from 7.22 to 7.50 MeV in excitation. Solid curves are results of DWBA calculations using complex-energy eigenstates, as explained in the text, for $f_{5/2}$ resonances. Optical potentials are those of Ref. 2. Spectroscopic factors and resonance parameters are given in the Table. (b) The ^{38}Ar$(d,p)^{39}$Ar cross sections at 10.06 MeV, from Ref. 2, for the three neutron unbound states from 7.56 to 7.73 MeV in excitation. Other information is as given for Fig. 3.

for population of the 8.89- and 9.01-MeV states in ^{37}Ar. On the left are conventional DWBA calculations assuming the states are bound by 0.1 MeV, while on the right are DWBA calculations using complex-energy eigenstates. In the Table, V_0 is the depth of the real Woods-Saxon potential used to compute the complex-energy eigenstate. Other parameters are $r_0 = 1.19$ fm, $a_0 = 0.65$ fm, $V_{so} = 5.8$ MeV, $r_{so} = 1.19$ fm, $a_{so} = 0.65$ fm. Optical parameters are those used in Ref. 1, and a D_0^2 value of 1.58×10^4 MeV^2 fm^3 is used in all calculations.[10, 15] The pole positions found and the resulting spectroscopic factors are summarized in the Table. Single-particle widths may be deduced from column 4, Im$\tilde{E}$. For these low-lying resonances, close agreement in magnitude between Gamow function (GF) and weakly bound-state (BS) cross sections is anticipated, and is found. Assuming the 8.89-MeV state is $p_{1/2}$ one has $S^{BS} = 0.12$, $S^{GF} = 0.11$, while assumption of $f_{5/2}$ gives $S^{BS} = 0.07$, $S^{GF} = 0.08$. Similarly if the 9.01-MeV state is $p_{1/2}$, $S^{BS} = 0.11$, $S^{GF} = 0.08$, or if it is $f_{5/2}$, $S^{BS} = S^{GF} = 0.05$. Without reliable data for c.m. angles less than 30° it is difficult to distinguish $f_{5/2}$ from $p_{1/2}$ at these excitation energies.[1]

A comparison of the complex-energy $p_{1/2}$ neutron eigenstate, used in the 9.01-MeV state calculations, with the corresponding bound state, is shown in Fig. 2. Thus these are the form factors relevant to the lower half of Fig. 1. Only the region from the origin to 30 fm is shown, although the state functions were originally computed to much larger distances. The striking similarity of bound and Gamow states in shape and magnitude is clearly seen. For a further comparison, including other types of unbound state functions, see Refs. 16 and 17.

Turning to ^{38}Ar(d, p) cross sections[2] for population of 11 resonances between 6.79 and 7.73 MeV in excitation, as shown in Figs. 3 and 4, it is again assumed the resonances are either $p_{1/2}$ or $f_{5/2}$. Other l_j values were tried, e.g. $d_{3/2}$, but only $p_{1/2}$ and $f_{5/2}$ give reasonable agreement or coincide with expectations from the analysis of lower-lying states.[2] Spectroscopic factors are given in the Table, and the optical parameters used are those of Sen *et al.*[2]

In the study of bound states in ^{39}Ar from 0.0 to 6.49 MeV in excitation, it was found that $\sum_i S_i^J$ is 0.81 for $p_{1/2}$ and ~0.24 for $f_{5/2}$.[2] If we assign every resonance observed to be $f_{5/2}$ the sum rises to ~0.57. Assigning the lowest one or three resonances as $p_{1/2}$ would essentially completely satisfy the $p_{1/2}$ sum rule.

In summary, treatment of stripping to fragmented single-particle resonances is seen to be reasonably straightforward even when complex-energy eigenstates are used to provide resonance form factors.

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