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SURFACE DELTA INTERACTION MODEL FOR NUCLEAR STRUCTURE ANALYZED BY PROTON INELASTIC SCATTERING

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December 1966

SURFACE DELTA INTERACTION MODEL FOR NUCLEAR STRUCTURE ANALYZED BY PROTON INELASTIC SCATTERING*

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ABSTRACT

Differential cross-sections and polarizations for 17.8 MeV protons scattered by nickel isotopes have been calculated by solving the coupled equations. The nuclear states were described in terms of their microscopic composition in the two-quasiparticle approximation. Two such model calculations were used. The effect of the presence of pure "two-phonon" states was investigated by generating them from the collective operator for the 2_1^+ state. The effects of the spin-flip term in the direct interaction were included. Comparison of the results with experiments suggests that, except for the collective 2_1^+ state, the two-quasiparticle method is inadequate for treating nuclear states. The vibrational treatment also fails. A much more elaborate structure calculation which treats the mixing between the two-quasiparticle and the "two-phonon" configurations is indicated.

1. INTRODUCTION

As is well known, the S-wave phase shift for nucleon-nucleon scattering decreases with increasing relative energy until it eventually becomes negative around 115 MeV, reflecting the repulsion at short distance. Inside the nucleus, the relative energy of two interacting nucleons is on the average rather high, so that a reduction in their interaction energy may be expected. At the surface however, the relative energy is low enough that there will be a net attraction in the S-states. Of course this behaviour is not automatically reproduced by the phenomenological forces with which it is convenient to work in nuclear structure calculations. Moszkowski and Green[1] were thus led to introduce a phenomenological force with the above surface localization assured by the inclusion of a delta function, which allows the force to be effective only when the interacting particles are at the nuclear surface. This they called a surface delta interaction (SDI). The zero range and sharp localization are naturally simplifying assumptions which it is hoped will simulate an authentic phonomenon.[2]

Within the last year the SDI has been applied to single-closed-shell nuclei,[3] to double-closed-shell nuclei,[4] to the (s,d)-shell nuclei[5] and to heavy deformed nuclei.[6,7] In all these cases the simple SDI was able to reproduce the experimental energy spectra as well as conventional forces which have more parameters (Gaussian with exchange forces; pairing plus quadrupole or octupole force model).

Faessler and Plastino[6] tested the SDI-wave functions in deformed nuclei by calculating electromagnetic transition probabilities. They obtained good agreement with the experimental results. But it is well known that the information which one gets from inelastic scattering with finite momentum transfer

is more sensitive to details of the nuclear structure. Here we want to test the SDI-functions in single-closed-shell nuclei[3] by inelastic proton scattering. We compare the results also with the scattering cross section calculated from wave functions for which a residual interaction of conventional form was used.[8,9]

The importance of the inelastic scattering as a means of investigating nuclear structure was recognized already long ago. But only transitions between pure shell model states had been studied.[10,11] After the recognition that collective states, which have a strong coupling with the electromagnetic field, have also a large inelastic scattering cross section, considerable success has attended the application of macroscopic collective models to the description of inelastic scattering by nuclei.[12]

For vibrational nuclei the macroscopic treatment was satisfying as far as the enhanced collective transitions were concerned (the 2_1^+ and 3_1^-). However little success has attended the application of the model to the so-called two-phonon states. Whereas a single deformation parameter is associated with all the states built up from quadrupole vibrations, in practice one has been forced to use almost as many parameters as there are transitions.[13] This reflects the well known property that from the two-phonon triplet, perhaps only the 4+ is as collective as described by the vibrational model.[14] The transition from the second 2+ to the first 2+ has in most cases about only half the electromagnetic strength predicted by the harmonic vibrational model. The 0+ is normally the least collective judging by the E2-transition from the 0+ of the triplet to the one-phonon 2+ state.

Such behaviour in spherical even-mass nuclei is at least qualitatively reproduced by microscopic calculations.[15,16] It seems therefore natural to use these microscopic descriptions for inelastic scattering calculations in place of the macroscopic collective one. Treatments of this problem were recently given by several authors.[17] In these initial calculations the distorted wave Born approximation was used. But this method has all the disadvantages of a first order perturbation theory. Therefore microscopic treatments of nuclei were incorporated recently into the coupled channel method.[18,19]

In this paper we calculate the (p,p')-cross sections to a number of states in the nickel isotopes, by solving the coupled equations. For the description of the nuclei we take the wave functions of Plastino, Arvieu and Moszkowski,[3] which were calculated utilizing the SDI.[1] We make some comparisons with results of calculations which utilize wave functions produced by an interaction which is effective throughout the nucleus.[8,9]

In Sect. 2 we review briefly the nuclear model utilizing the SDI. In Sect. 3 the scattering theory is sketched. In Sect. 4 our calculations are described and compared with experimental measurements and in Sect. 5 we discuss the results.

2. THE NUCLEAR MODEL

We describe the nucleus in the conventional way with a shell model potential and a residual interaction:

$$H = \sum_{\alpha} \epsilon_{\alpha} c_{\alpha}^{\dagger} c_{\alpha} + \frac{1}{4} \sum_{\alpha\beta\gamma\delta} \langle \alpha\beta | V(1,2) | \gamma\delta \rangle \, c_{\alpha}^{\dagger} c_{\beta}^{\dagger} c_{\delta} c_{\gamma} \qquad (1)$$

For the residual interaction, the Surface Delta Interaction (SDI) is utilized:

$$V(1,2) = - \frac{4\pi G}{(U_0 R_0)^4} (1 + \eta + \eta P_{\sigma}) \, \delta(r_1 - R_0) \, \delta(r_2 - R_0) \, \delta(\theta_{12}) \qquad (2)$$

The capital G is the coupling constant. The parameter η gives the relative strength of the spin-exchange force P_{σ}. The three δ-functions ensure that particles 1 and 2 only interact if they are at the same place on the nuclear surface. If one compares the radial wave functions for an oscillator, and also for a Saxon-Woods potential,[20,21] at the nuclear surface, one finds that for each potential, all its wave functions, to a good approximation, are equal (say to U_0). (The phase is chosen always in such a way that the tail of the radial wave function is positive). In Eq. (2) this amplitude $U_0 = U(R_0)$ is absorbed into the coupling constant in such a way that the radial integral with the two radial δ-functions and the factor $(U_0 R_0)^{-4}$ gives unity. In that way results are essentially independent of whether an oscillator or Woods-Saxon central potential is used.

Here we are considering only single-closed-shell nuclei. The nucleons outside closed shells can therefore only interact in the isospin state T = 1. Because a zero-range force gives only a finite matrix element for spatially symmetric states, the eigenvalue of P_σ is always $P_\sigma = -1$ and so the value of η does not enter in these calculations.

The method of solving the problems whose Hamiltonian is given by Eq. (1) is described in detail in Ref. 3. Here we only outline briefly the procedure.

The quasiparticle Tamm-Dancof approximation (QTDA) is utilized. The numerical calculations are very much simplified by the fact that the monopole part of the particle-particle matrix element is state independent (pairing force). Therefore one can find the quasiparticle (Bogoliubov-Valatin) transformation by solving two simple BCS-equations.[15] The diagonalization of the residual interaction between quasiparticles within a two-quasiparticle basis leads then to the following secular equation:

$$\sum_{cd} P_{J\pi}(j_a j_b j_c j_d)\ \xi^i_{cd} = E^i_{J\pi}\ \xi^i_{ab} \tag{3}$$

The two-quasiparticle wave function has the form:

$$|JM;\pi i\rangle = \sum_{a\leq b} [1+\delta_{a,b}]^{-1/2}\ \xi^i_{ab}\ A^\dagger_{JM}(a,b)|0\rangle$$

$$A^\dagger_{JM}(a,b) = [\alpha^\dagger_a\ \alpha^\dagger_b]^M_J \tag{4}$$

where the square bracket denotes vector coupling. The characters

$$\alpha^+_{am} = U_a\ C^+_{am} - (-)^{j_a-m} V_a\ C_{a-m} \tag{5}$$

are quasiparticle creastion operators and $|0\rangle$ is the BCS-ground state (quasiparticle vacuum). The superscript i distinguishes between different solutions to the same total angular momentum J and the same parity π.

The matrix P and detailed expressions for the antisymmetrized particle-particle and the direct and the exchange particle-hole elements are given in Ref. 3. There also the single-particle energies and the values for the coupling parameter G are listed.

3. SCATTERING THEORY

The direct interaction. Since the nuclear states are here described in terms of their nucleonic structure, the interaction of the scattered particle with the nucleus is by way of a two-nucleon potential rather than the one-body deformed potential used in macroscopic treatments. For the form of the interaction between nucleons we use

$$V(\underset{\sim}{r},\underset{\sim}{r}_i) = (V_0 + V_1\, \underset{\sim}{\sigma} \cdot \underset{\sim}{\sigma}_i) \exp\left[-\left(\frac{|\underset{\sim}{r} - \underset{\sim}{r}_i|}{\rho}\right)^2\right] \qquad (6)$$

where $\underset{\sim}{r},\underset{\sim}{\sigma}$ are the position and spin coordinates of the scattered particle, and $\underset{\sim}{r}_i,\underset{\sim}{\sigma}_i$ of the i'th nucleon in the nucleus. The spin scalar and vector parts, V_0 and V_1, depend in general on the isospin so that, as noted previously,[22] neutrons and protons may be expected to scatter differently in exciting the same nuclear state (except when the excitation is characterized by T=0).

If we use some simple ideas about the free two-nucleon potential one finds that for protons scattering from states described by neutron correlations, the spin-dependent part V_1 is small compared to V_0.[22] Recent evidence

from the analysis of scattering on the zirconium isotopes also supports this conclusion.[23] Nevertheless the direct interaction that should be used, may be very different from the vacuum interaction, as discussed qualitatively elsewhere,[22] and may quite likely depend upon the region of the periodic table and the details of the nuclear description. This is so because in practice, both the structure and scattering calculations are carried out in truncated spaces of the basic states. This causes the interaction within the truncated spaces to be modified from the free interaction by terms referring to the eliminated part of the space. It may therefore be referred to as an effective interaction: effective within a prescribed subspace of basic states. Its structure may be very complicated, since in principle, it can be both non-local and complex valued. Our approach corresponds to parameterizing it, assuming for simplicity that it is a real local central potential of Gaussian shape. For the range we take 1.85 F, such as has often been used in shell model calculations. For most of our calculations $V_1=0$, but the effect of spin-dependence was examined for ^{60}Ni.

Form factors. We expand the interaction in a multipole series to isolate the parts giving rise to the transfer of specific units of spin, orbital, and total angular momentum:

$$V(\underline{r},\underline{A}) = \sum_{L,S,J} (-)^{L+S+J} V_S \mathfrak{F}_{LSJ}(r,\underline{A}) \cdot \mathcal{Y}_{LSJ}(\hat{r},\underline{\sigma}) \qquad (7)$$

The sum on S=0 or 1, refers to the scalar and vector parts of Eq. (6). The operators $\mathfrak{F}$ and $\mathcal{Y}$ were defined in Ref. 22. Here the nuclear coordinates, denoted by $\underline{A}$, are contained only in the factor $\mathfrak{F}$, which also depends on the coordinate $|\underline{r}|$ of the scattered particle. The reduced matrix element of $\mathfrak{F}$

with respect to a pair of nuclear states α_1, α_2 is called a form factor

$$\mathcal{F}^{\alpha_1\alpha_2}_{LSJ}(r) = \langle \Psi_{\alpha_1 J_1}(\underline{A}) \| \mathcal{J}_{LSJ}(r,\underline{A}) \| \Psi_{\alpha_2 J_2}(\underline{A}) \rangle \tag{8}$$

and is to speak the way the nucleus appears to the scattered particle, as a function of their separation, for that particular transition. The structure of this form factor is actually rather simple. This follows from the fact that the scattered particle acts as a one-body operator on the particles of the nucleus because the interaction between nucleons is two-body.[24] Thus if the form factor for exciting a valence nucleon from the single-particle state a to b is

$$F^{ab}_{LSJ}(r) = \langle \psi_a(\underline{r}') \| \mathcal{J}_{LSJ}(r,\underline{r}') \| \psi_b(\underline{r}') \rangle \tag{9}$$

then $\mathcal{F}$ is merely some linear combination of such elementary form factors, however complicated the structure of the states may be:

$$\mathcal{F}^{\alpha_1\alpha_2}_{LSJ}(r) = \sum_{a,b} S^{ab}_{LSJ}(\alpha_1\alpha_2)\, F^{ab}_{LSJ}(r) \tag{10}$$

The amplitude, S^{ab}, for the single-particle transition $a \rightarrow b$ can be calculated from the nuclear wave functions. As an example, its form for the transition from the ground state to an excited state can be seen in Eq. (36) of Ref. 22. Some examples of form factors for transitions in the nickel isotopes can also be seen there. These form factors are the physical ingredients of the <u>off-diagonal</u> matrix elements of the interaction that appears in the coupled equations describing the elastic and inelastic processes. The <u>diagonal</u> elements are not calculated in

the same way but instead are represented by a spherical optical potential. There are two reasons for this departure from a strict microscopic description. The first is that the elastic scattering, represented by the diagonal parts, is a highly coherent process (every nucleon must recoil) and is therefore hard to calculate accurately. The second is that there may be important contributions from channels that are not explicitly treated. Such contributions can be accounted for approximately in the optical potential. This makes a phenomenological treatment of the diagonal parts desirable even if the first difficulty could be surmounted.

Coupled channel method. Within the subspace of channels chosen, the coupled channel method solves the scattering problem to all orders in the interaction. It does however require appreciable computer time, and one wants to establish first whether the distorted-wave Born method, which solves the problem to first order, would be adequate. It is clear that for the collective states it should be fiarly good. For the non-collective states however, multiple scattering may compete with the direct transition from the ground, and was shown to be important for some of the levels in the nickel isotopes.[19] Therefore we must in general use the coupled channel method when we are interested in non-collective as well as collective transitions. All results reported here have been obtained by solving the coupled equations.

4. RESULTS

Two-quasiparticle states. Proton cross-sections for the coupled system comprising two 0^+, four 2^+ and one 4^+ state in the nickel isotopes of mass 58, 60, and 62 were calculated using the two-quasiparticle description of these isotopes. Two nuclear descriptions were used, one based on the surface-delta-interaction calculation of Faessler and Plastino, and one on the volume finite-range interactions calculation (referred to for brevity as VFI) of Arvieu, Salusti, and Veneroni. The results of the coupled channel calculations using these microscopic descriptions can be compared in Figs. 1-3, where the experimental data[25] is also shown.

The parameters of the calculation consist of those describing the optical model potential and the direct interaction between nucleons. The optical model parameters are essentially fixed by the elastic scattering and therefore are not available for adjustment of the excited state cross-sections. (See Table I.) The range of the direct interactions we take as having the fixed value of 1.85 F typically used in shell model calculations. The scalar form factor of the collective 2^+ state is dominant over its vector form factor so that V_0 can be determined by adjusting it to obtain agreement in magnitude with this level. The angular distributions based on both structure calculations are seen to be in quite good agreement with experiment. The strength V_1 of the spin-dependent part of the interaction is relevant only for those higher lying states whose vector form factor (giving rise to spin-flip transitions) is comparable to the scalar form factor. Since the wave functions of these states are not established, we cannot use them to determine V_1. As remarked earlier, there is some reason to believe that $V_1 \ll V_0$. We therefore have performed most of our calculations with

$V_1=0$. In comparing the results for the higher lying levels with experiment it should therefore be kept in mind that no adjustment of parameters was made.

In Fig. 4 we do show the effect of taking various values for V_1 including a repulsive value. The cross-section to the collective 2_1^+ state is insensitive to any of these variations, as are several of the others. The 2_4 level, which has a vector form factor of magnitude comparable to its scalar, is quite sensitive to the value of V_1, especially the polarization. However there is no data available on this level, which lies presumably in the region above 3 MeV. The excited 0^+ level, which cannot be excited directly by the spin-flip interaction, nonetheless is fairly sensitive to V_1, especially the polarization due to its coupling to other excited states. There is no polarization data, and the computed differential cross-section is not in good agreement with the data.

Two-phonon states. The available experimental evidence concerning so-called two-phonon states indicates considerable departures from the pure harmonic vibrational model. The evidence consists of the sometimes large splittings between the members of the two-phonon triplet, deviations from the predicted intensity rules for E2 radiation,[14] and the finite static quadrupole moment of the two-phonon 2^+ state in those nuclei for which the measurements have been made. Nevertheless some of the two-phonon character is expected to persist, though it may be shared by other neighboring states. The type of nuclear structure calculation described in Sect. 2 does not produce two-phonon states because they involve four-quasiparticle configurations which were not included. However we can use the coherent operator that produces the 2_1^+ state to generate a "two-phonon" triplet of wave functions. Presumably these are not eigenfunctions of the problem, and would mix with neighboring two-quasiparticle states, but we can

consider them as an extreme limit. The form factors by which each member of the triplet is connected to the collective 2_1^+ state can be obtained easily upon using the quasiboson commutation rules for the two-quasiparticle operators introduced in Eq. (4). The coupled equations for proton scattering on a model nucleus having excited states consisting of two 2^+ states calculated in the two-quasiparticle model, including the collective one, plus a triplet of states built on the latter, were solved and the results are shown in Fig. 5.

Surface direct interaction. For the same reason that the residual interaction in the structure calculation is thought to be concentrated at the nuclear surface, the direct interaction of the scattered particle with the bound nucleons of the nucleus may also be most effective in the surface. We therefore recalculated the cross sections for ^{60}Ni after introducing a radius dependent factor that suppresses the strength of the interaction in the interior. The factor used is $1 - \alpha f(r)$ where $f(r)$ is the Woods-Saxon form factor and α determines the degree of suppression. The shape of this suppression factor is such that the interaction has its full value outside the nucleus, is reduced by the factor $1 - 1/2\alpha$ at the surface, and by $1 - \alpha$ at $r = 0$. For the shape parameters of $f(r)$ we used those of the real part of the optical potential. The results of the calculation for $\alpha = 0.5$ are shown in Fig. 6. The polarizations are little effected while differential cross sections for several levels were somewhat altered, mainly in magnitude.

5. DISCUSSION

We take the point of view that the reaction calculation has been handled well enough to allow conclusions to be drawn about the nuclear description by comparing calculated cross sections with the data. It is true that the direct interaction is subject to considerable uncertainty as discussed earlier. Assuming a real central potential with a fixed range about that of the vacuum interaction, the one well depth V_0 was determined by the collective 2_1^+ state. Variation of the remaining parameter V_1 was not successful in achieving an overall agreement for all the remaining levels. The effective interaction may be more complicated than the one used so that our criticism of the nuclear description of these nuclei is somewhat blunted. Nevertheless we reach the tentative conclusion that aside from the fair success with the collective 2^+ state, the two-quasipárticle method has not so far, and possibly cannot, provide a very satisfactory description of the states that are weakly coupled to the ground. Some of the angular distributions are not bad, but often the absolute cross-sections are in error by a factor of two to three.

We can draw attention to the following points. For ^{58}Ni the quadrupole strength going to the levels above the collective one is stronger than the two-quasiparticle description predicts but when one of the levels is described as a "two-phonon" state, the strength comes out about right. However in ^{60}Ni the observed strength is greater than predicted in either case. On the other hand the 0^+ excited state which is known only in ^{60}Ni and may therefore be assumed small in ^{58}Ni, is weaker than predicted when described either as a two-quasiparticle or "two-phonon" state. The 4^+ level is weaker than predicted by the two-quasiparticle description, but stronger than the "two-phonon" description.

There is a suggestion in this that the overall picture would be improved by a four quasiparticle description of the nuclei since presumably some two-phonon like character would emerge, possibly admixed among several states. This apparently would yield better results than if it were concentrated into the triplet. Some progress on this difficult problem has been reported recently.[26]

Polarization data is becoming available. Examination of the figures suggests that it could be useful in deciding between alternate descriptions of a given state. However the available descriptions of the nickel isotopes are not yet good enough to take advantage of this potential.

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Table I. Optical model parameters (lengths in Fermi, energy in MeV) (after F. G. Perey, Phys. Rev. 131, 745 (1963).

	V	W	W_0	r_v	r_w	a_v	a_w	V_{SO}	r_{SO}	a_{SO}	r_c
^{58}Ni	47.59	0	10.0	1.251	1.245	0.677	0.414	9.77	1.251	0.677	1.25
^{60}Ni, ^{62}Ni	47.59	0	12.5	1.251	1.245	0.677	0.414	9.77	1.251	0.677	1.25

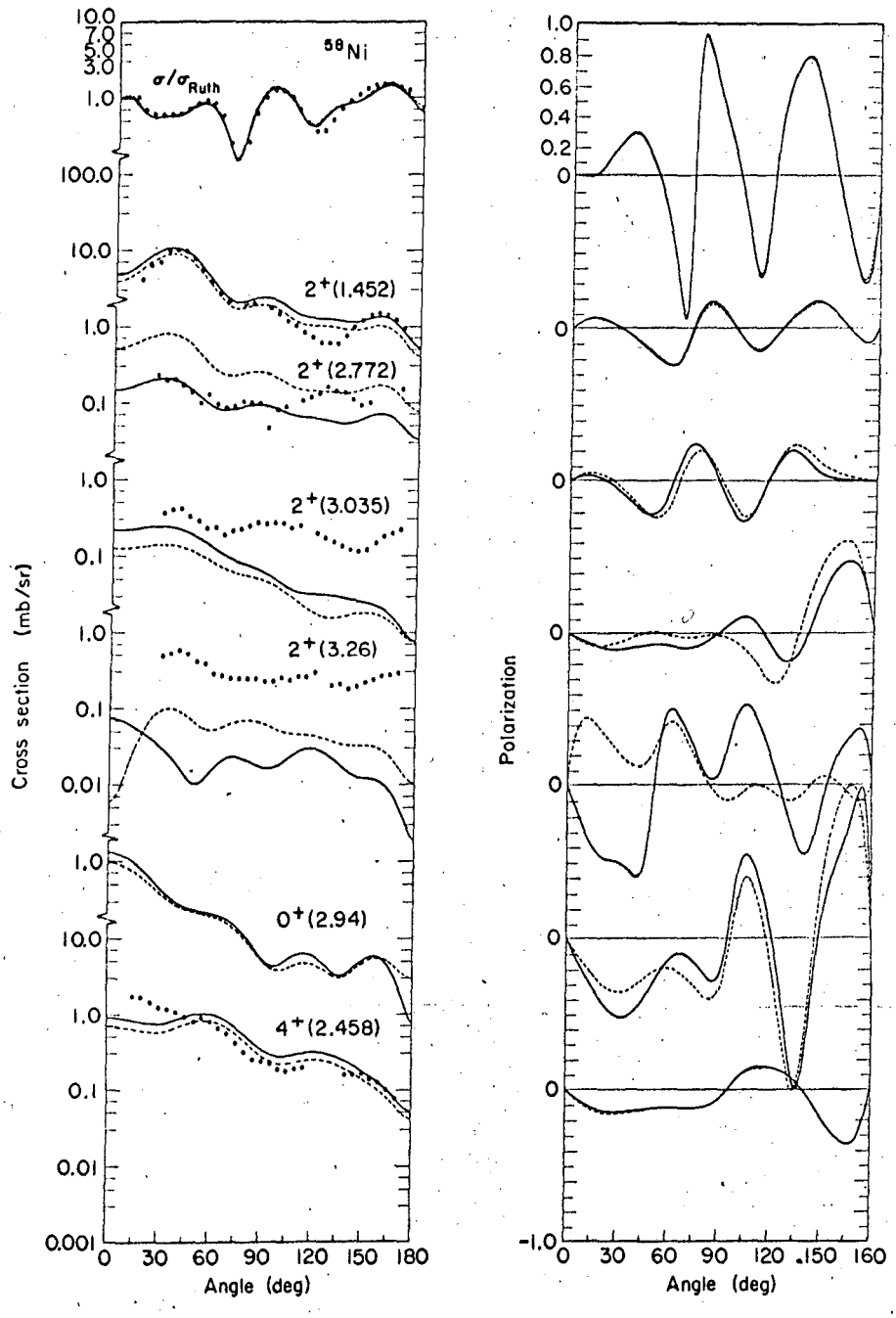

Fig. 1. Differential cross sections and polarizations for 17.8 MeV protons on ^{58}Ni. Curves represent coupled equation calculations using the microscopic descriptions of the nucleus mentioned in Sect. 4; solid line corresponds to the SDI and dashed to VFI structure calculation. Optical model parameters are listed in Table I. The direct interaction, Eq. (6), had the parameters V_0 = -65 MeV, V_1 = 0, ρ = 1.85F. Data is from Jarvis, Harvey, Hendrie, and Mahoney, Ref. 25.

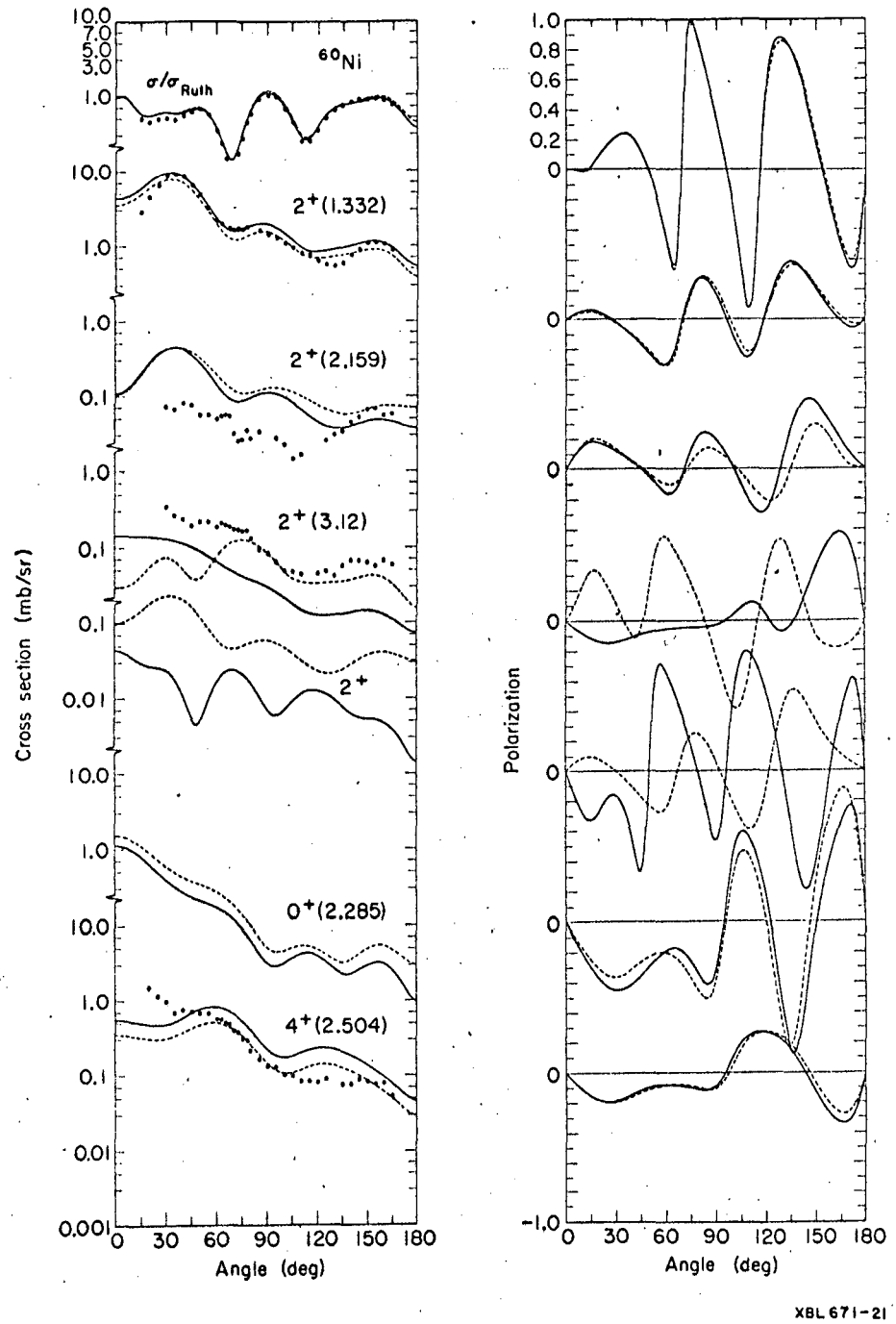

Fig. 2. Differential cross sections and polarizations for 17.8 MeV protons on ^{60}Ni. See caption to Fig. 1. Direct interaction, Eq. (6), had parameters V_0 = -55 MeV, V_1 = 0, ρ = 1.85 F.

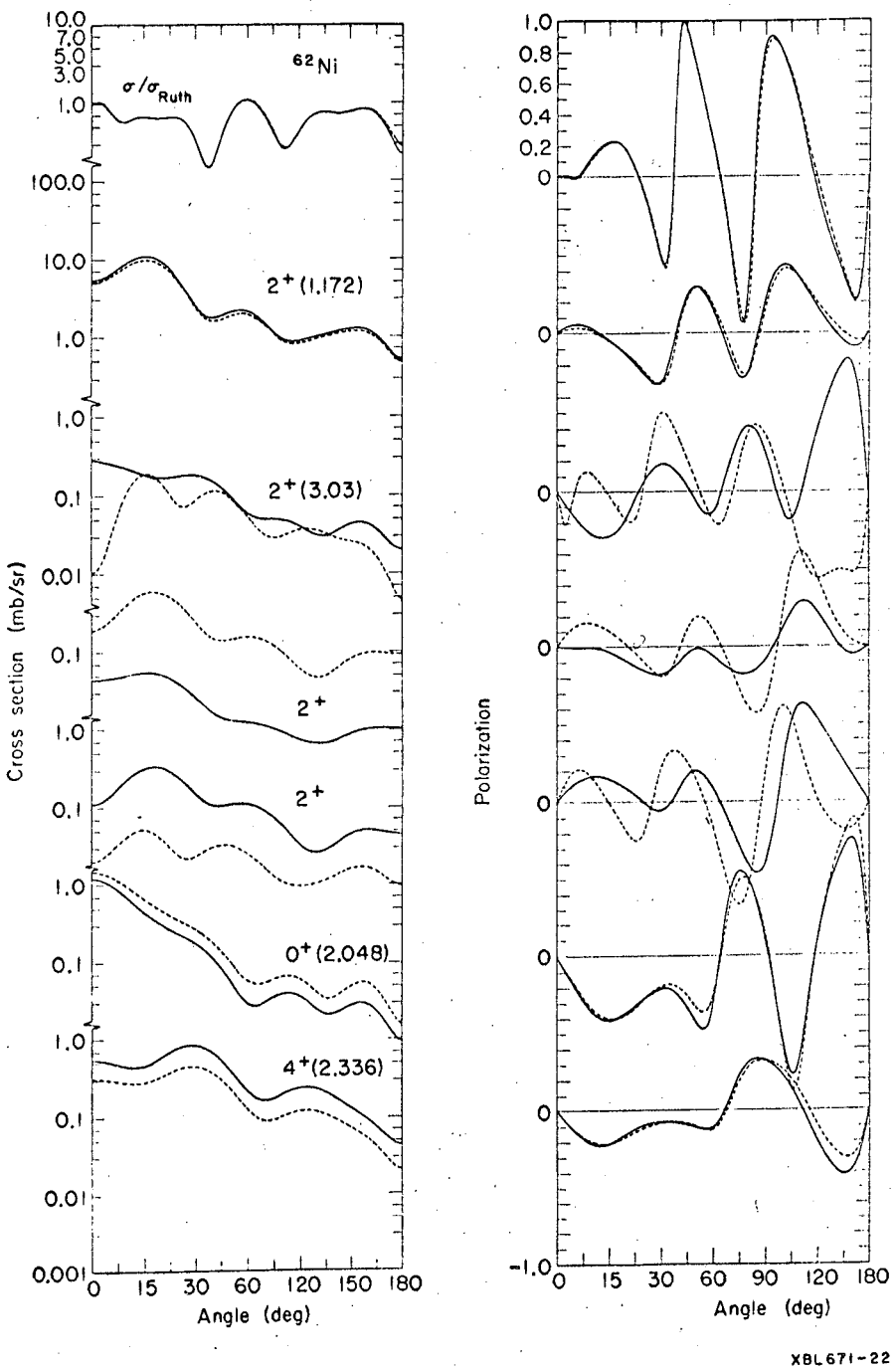

Fig. 3. Differential cross sections and polarizations for 17.8 MeV protons on ^{62}Ni. See caption to Fig. 1. Direct interaction, Eq. (6), had parameters V_0 = -55 MeV, V_1 = 0, ρ = 1.85 F.

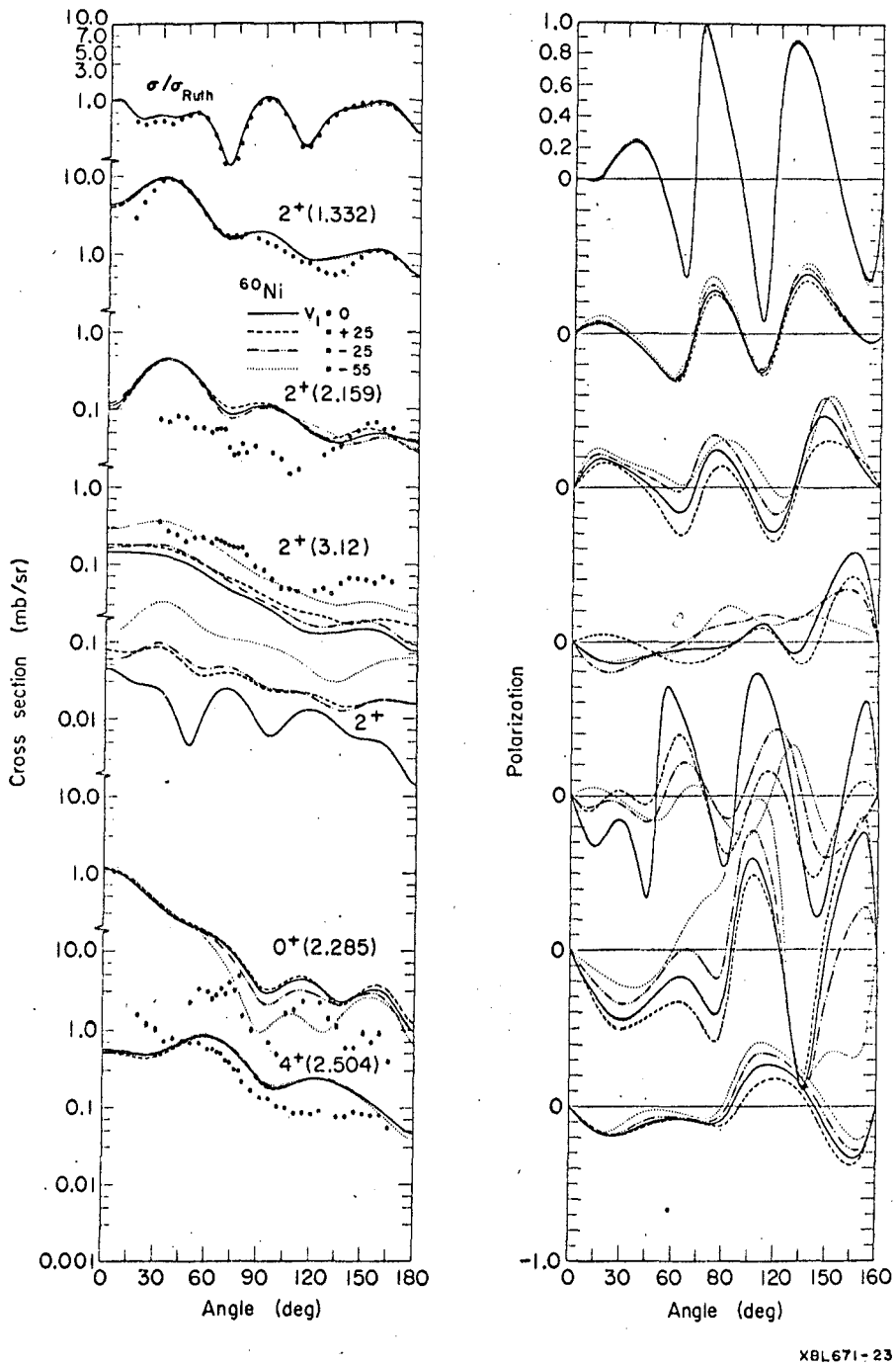

Fig. 4. Differential cross sections and polarizations for 17.8 MeV protons on ^{60}Ni. Curves illustrate effect of different values for the spin dependent part, V_1, of the direct interaction, Eq. (6), computed by solving the coupled equations using the microscopic wave functions of the SDI structure calculation described in Sect. 2. Other direct interaction parameters are V_0 = -55 MeV, ρ = 1.85 F.

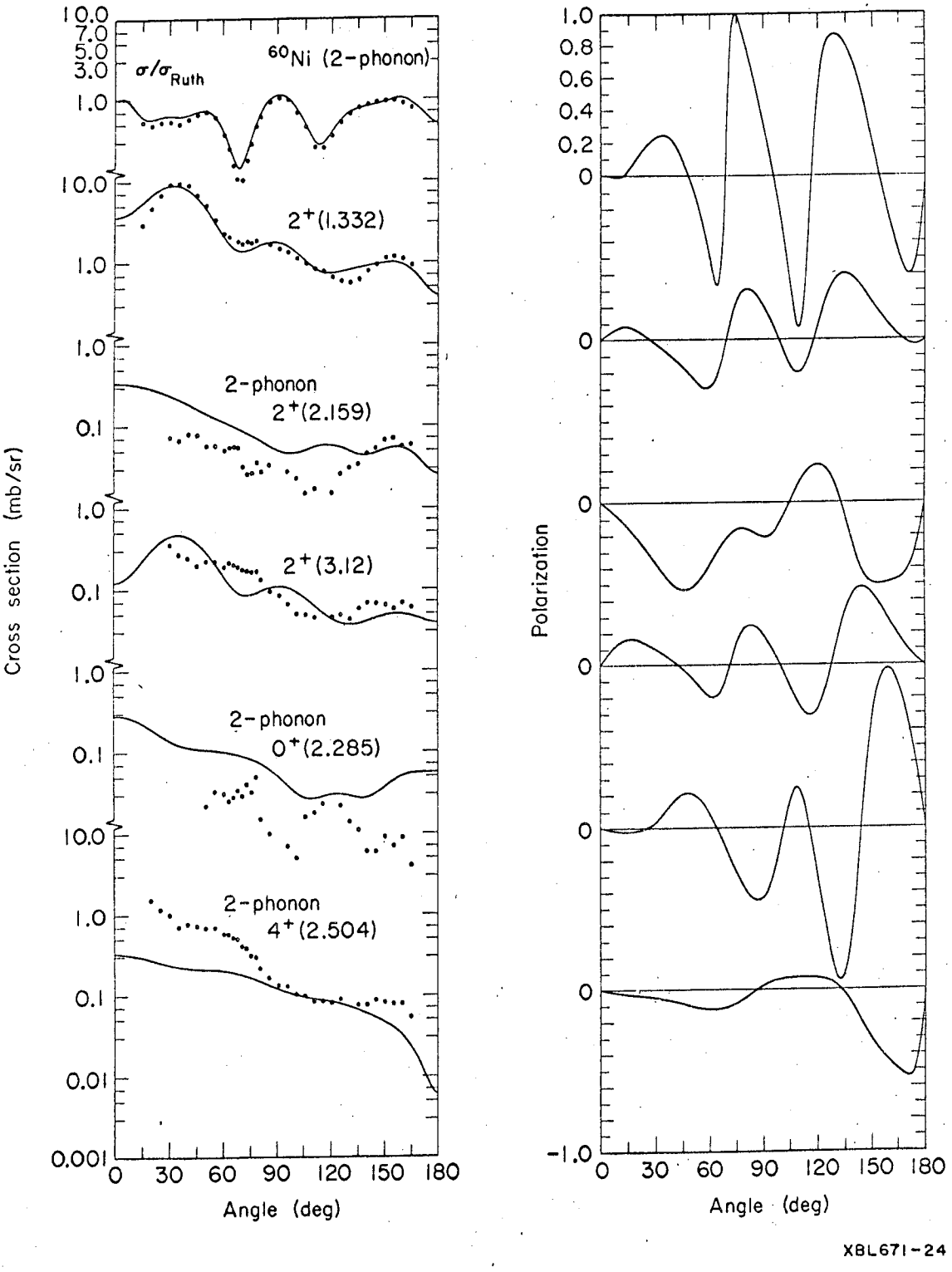

Fig. 5. Protons on ^{60}Ni at 17.8 MeV. Calculation includes two-phonon states described in terms of their microscopic structure as described in Sect. 4, as well as two-quasiparticle states generated by the SDI model of which the lowest 2_1^+ is the collective one. Direct interaction as in Fig. 2.

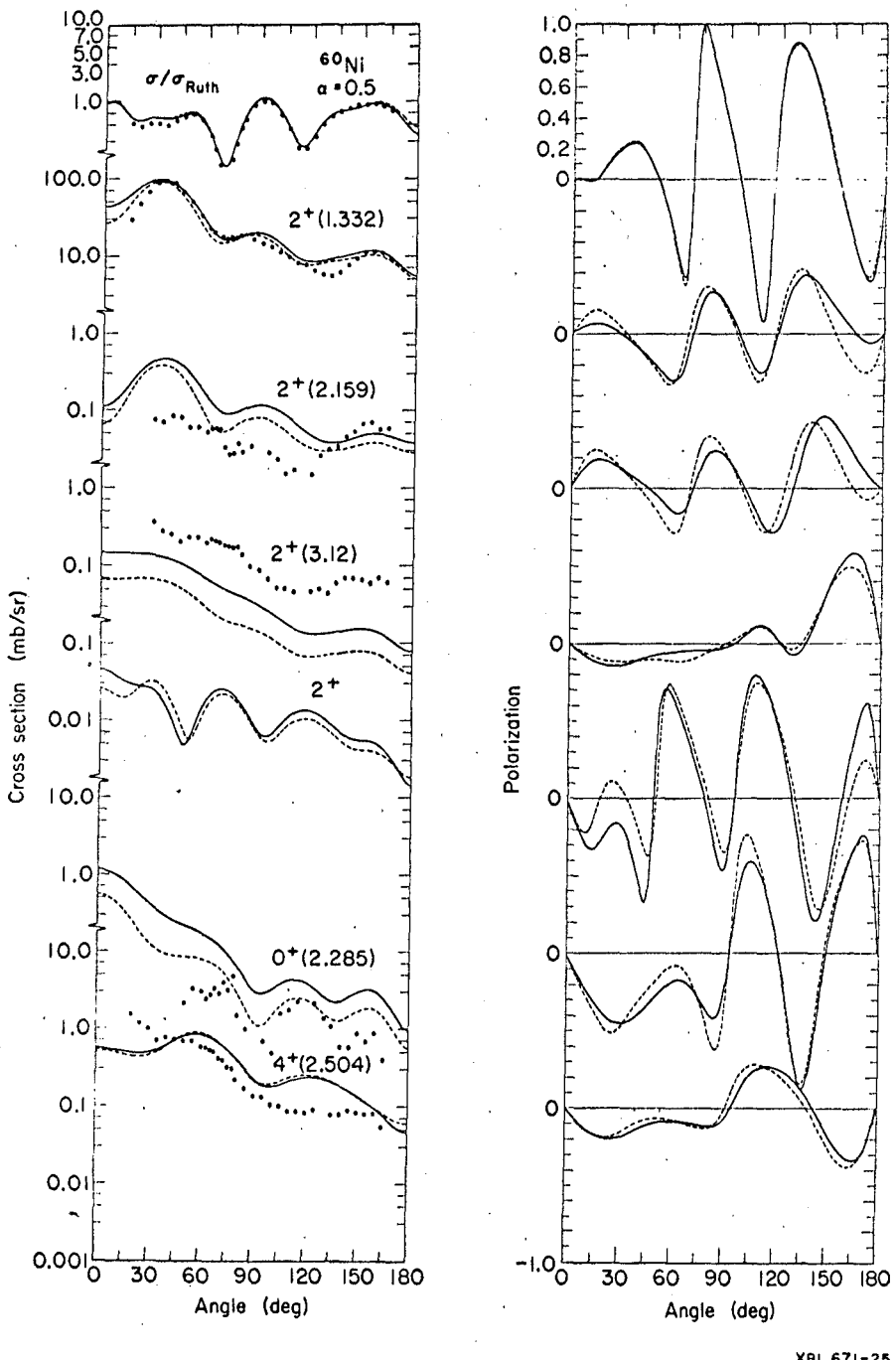

Fig. 6. Protons on ^{60}Ni at 17.8 MeV. Using the microscopic description of the SDI model described in Sect. 2, the coupled equations for the scattering were solved. including a radial suppression factor that weakens the strength of the direct interaction in the nuclear interior as discussed in Sect. 4. Results for this case are shown by dashed line for which the direct interaction Eq. (6), had parameters V_0 = -65 MeV, V_1 = 0 MeV, ρ = 1.85 F. Solid line corresponds to no suppression and parameters as in captions to Fig. 2.

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