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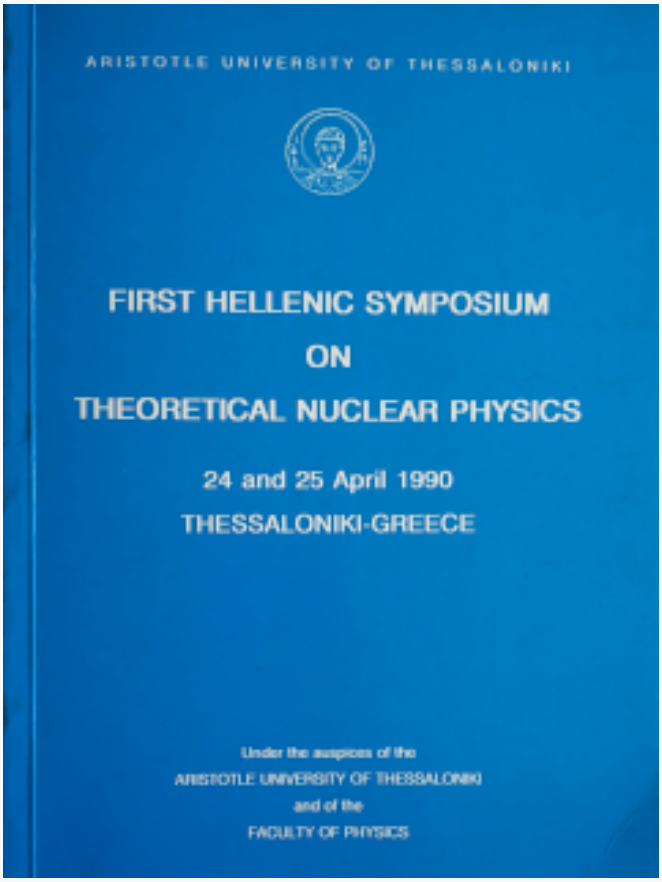

Microscopic Calculation of the Optical Model Potential from One Boson Exchange Potentials

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Microscopic Calculation of the Optical Model Potential from One Boson Exchange Potentials *

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ABSTRACT: A new method for calculating the optical model potential from One Boson Exchange Potentials (OBEPs) is developed. The G-matrix is calculated by solving the Bethe-Goldstone equation in momentum space. Using vector brackets these G-matrix elements can be transformed from the center of mass representation into the laboratory system. This allows the evaluation of the G-matrix interaction between nucleons in bound states and those in a plane-wave state. The lowest order contribution to the real part of the potential comes from the Hartree-Fock term, while the lowest order contribution to the imaginary part comes from the two-particle-one-hole (2p1h) diagram. Calculations for ^{16}O and ^{40}Ca have been carried out. Local approximations are obtained by describing the results for the central part in terms of a Woods-Saxon potential and those for the spin-orbit part in terms of the corresponding derivatives. The dependence of these potentials on energy and angular momentum is discussed. The parameters for these local approximations are in good agreement with empirical fits.

The calculation of properties of finite nuclei from nucleon-nucleon potentials is a longstanding problem in nuclear physics. In the early attempts (Kuo and Brown 1966, 1968), purely phenomenological potentials were used, like the Hamada-Johnston potential (Hamada and Johnston 1962). Nowadays more advanced one boson exchange potentials

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(OBEPs) exist, like the Bonn potential (Holinde *et al.* 1972a, 1972b) and the new version of it (Machleidt *et al.* 1987), derived by field theoretical techniques and viewed as a means of describing the low energy and small momentum transfer characteristics of QCD. It is therefore of great interest to try to determine properties of finite nuclei from modern OBEPs. The optical model potential for negative energies is particularly interesting, since, despite the richness of experimental data existing for the potential at positive energies (Hodgson 1984), little is experimentally known at negative energies. This lack of information has in fact motivated the recent extension of the optical potential into the negative energy region attempted through the dispersion relation method of Mahaux and Sartor (1987). The real part of a microscopically calculated optical potential can be directly compared to the results of this work.

Another source of motivation for the microscopic derivation of the optical potential for finite nuclei is the experimental data for the longitudinal quasielastic charge response (Barreau *et al.* 1983, Meziani *et al.* 1984, 1985), obtained by deep inelastic electron scattering. For their interpretation one needs an accurate calculation of the particle–hole interaction.

In this paper we use the self-energy approach of Mahaux *et al.* (1985). Our method is described in detail in Bonatsos and Müther (1989). We start by solving the Bethe–Goldstone equation for the Brueckner G-matrix in momentum space and in the center of mass coordinates. This equation reads

$$G(\omega) = V + V\frac{Q}{\omega - H_0}G(\omega), \tag{1}$$

where Q is the Pauli operator, V is the nonlocal nucleon–nucleon potential and ω is the starting energy. For V we use the relativistic OBEP of the Bonn group (Holinde *et al.* 1972a, 1972b). The equation is solved through use of the matrix inversion method (Haftel and Tabakin 1970). From this calculation we obtain matrix elements of the form

$$< klSJ_SKLT|G|k'l'SJ_SKLT >, \tag{2}$$

where k (l) is the relative momentum (angular momentum), K (L) is the corresponding center of mass quantity, S (T) stands for spin (isospin) and by definition $\vec{J}_S = \vec{l} + \vec{S}$.

The lowest order contribution to the real part of the optical potential comes from the Brueckner–Hartree–Fock term, which is given by (Kuo and Brown 1966)

$$V_{HF}(l_1, j_1; k_1, k_1') =$$

$$= \frac{1}{2(2j_1+1)} \sum_{n_2 l_2 j_2 JT} (2J+1)(2T+1) < k_1 l_1 j_1 k_2 l_2 j_2 JT|G|k_1' l_1 j_1 n_2 l_2 j_2 JT >, \quad (3)$$

where J denotes total angular momentum, k_1, l_1, j_1 are the quantum numbers of the incoming particle (unbound states we describe as plane waves), k_1', l_1, j_1 are these of the outgoing particle and n_2, l_2, j_2 are the quantum numbers of the hole, which is described as a harmonic oscillator eigenstate. In the case of ^{16}O the occupied levels are characterized by $n_2 = 0$, $l_2 = 0, 1$, the harmonic oscillator having $\hbar\omega_{osc} = 14.0$ MeV, which corresponds to an oscillator length in the laboratory system $\alpha = 1.72$ fm. In ^{40}Ca one has the levels with $n_2 = 0$, $l_2 = 0, 1, 2$ and $n_2 = 1$, $l_2 = 0$, the oscillator parameters being $\hbar\omega_{osc} = 10.4$ MeV, $\alpha = 2.00$ fm.

The lowest order contribution to the imaginary part of the potential comes from the two-particle–one-hole (2p1h) contribution to the self energy, given by

$$V_{2p1h} = \frac{1}{2(2j_1+1)} \sum_{n_2 l_2 j_2} \sum_{lL} \sum_{JJ_S ST} \int q^2 dq \int Q^2 dQ (2J+1)(2T+1)$$

$$< k_1 l_1 j_1 n_2 l_2 j_2 JT|G|qlSJ_SQLT >< qlSJ_SQLT|G|k_1' l_1 j_1 n_2 l_2 j_2 JT >$$

$$\pi\delta(\omega_1 + (\epsilon_{n_2 l_2} - \epsilon_F) - \frac{Q^2}{4m} - \frac{q^2}{m}), \quad (4)$$

where k_1, l_1, j_1 refer to the incoming particle, k_1', l_1 j_1 refer to the outgoing particle, q, Q, l, L refer to the intermediate states (in the center of mass system) and n_2, l_2, j_2 refer to the hole. The energy eigenvalues of the harmonic oscillator used for the description of the holes are

$$\epsilon_{n_2 l_2} = \hbar\omega_{osc}(2n_2 + l_2 + \frac{3}{2}), \quad (5)$$

with the parameter values given above. For ϵ_F we use the value corresponding to nuclear matter of Fermi momentum $k_F = 1.4$ fm^{-1}.

The matrix elements needed in V_{2p1h} we calculate from the matrix elements (2) in two steps. First, we transform the bra (or the ket) from the center of mass system into

the laboratory system. This transformation involves the use of vector brackets (Wong and Clement 1972), as described in detail in Bonatsos and Müther (1989). As a result we obtain matrix elements of the form

$$< klSJ_SKLT|G|k_1l_1j_1k_2l_2j_2JT > . \tag{6}$$

From them we can easily calculate matrix elements in a mixed representation (where some states are described by plane waves and some others by harmonic oscillator eigenstates) as shown by the following example

$$< klSJ_SKLT|G|k_1l_1j_1n_2l_2j_2JT >=$$

$$\int k_2^2 dk_2 \alpha^{3/2} P_{n_2l_2}(\alpha k_2) < klSJ_SKLT|G|k_1l_1j_1k_2l_2j_2JT >, \tag{7}$$

where by P_{nl} we denote harmonic oscillator radial wave functions.

In order to calculate the matrix elements needed in V_{HF}, we perform once more the transformation which led to the matrix elements (6), obtaining matrix elements of the form

$$< k_1l_1j_1k_2l_2j_2JT|G|k_1'l_1'j_1'k_2'l_2'j_2'JT > . \tag{8}$$

We then calculate matrix elements involving harmonic oscillator states in a way analogous to eq. (7).

Following common practice, we try to fit the microscopic results for the optical model potential by the Fourier–Bessel transform of a local potential. We consider a Woods–Saxon potential of the form

$$U(r) = \frac{-V_V}{1 + exp(\frac{r-r_V}{a_V})}, \tag{9}$$

and we calculate its Fourier–Bessel transform

$$< k|U|k' >= \int_0^\infty j_l(kr)U(r)j_l(k'r)r^2 dr, \tag{10}$$

where $j_l(kr)$ are the spherical Bessel functions of the first kind. The three free parameters of the Woods–Saxon potential we fit to the microscopic results by least square fitting.

In the same manner, we fit the microscopic results for the spin–orbit potential using a derivative Woods–Saxon potential (Hodgson 1984) of the form

$$U_{ls} = V_{ls}(\frac{\hbar}{m_\pi c})^2 \frac{1}{r}\frac{d}{dr}(\frac{1}{1+exp(\frac{r-r_{ls}}{a_{ls}})})\vec{l}\cdot\vec{\sigma}, \tag{11}$$

where

$$\vec{s} = \frac{\hbar}{2}\vec{\sigma}, \tag{12}$$

$$(\frac{\hbar}{m_\pi C})^2 = 2.000 fm, \tag{13}$$

and $m_\pi = 139.6$ MeV/c^2 (the pion mass).

Extensive applications of the method have been made to ^{16}O (Bonatsos and Müther 1989) and ^{40}Ca (Bonatsos and Müther 1990a). The main results are summarized here:

i) The Woods–Saxon local approximation to the microscopic optical model potential works well in the low momentum (k) region, while at large k the local approximation overshoots the microscopic results. This is the case for the real and the imaginary part. This disagreement can be avoided by using local potentials which are depressed at the origin and enhanced at the surface, as the ones suggested by Fiedeldey and Sofianos (1983).

ii) The Woods–Saxon potentials obtained from the fits of the real and imaginary parts show a systematic l-dependence, again similar to the one suggested by Fiedeldey and Sofianos (1983). As l increases the depth of the potentials decrease, while the radius and the steepness increase. In other words, the nucleus is more transparent for larger l.

iii) The radius and steepness parameters of the so obtained Woods–Saxon potential agree very well with the parameters determined empirically (Co' *et al.* 1988, Mahaux and Sartor 1988), while the depth parameter for the real part is found to be underestimated by roughly 25%–30%. This is attributed to the lack of binding, which is known to characterize the OBEP used in this calculation. Improved results can be obtained with the new Bonn potential (Machleidt *et al.* 1987).

iv) Also the parameters for the fit of the spin–orbit potentials are similar to the empirical ones, the depth of the potential being underestimated by the same percentage.

The contributions made by 3p2h, 1p2h and 2p3h diagrams to the optical model potential are currently under investigation (Bonatsos and Müther 1990b). In addition, the

improvements obtained through use of the new Bonn potential (Machleidt *et al.* 1987) are under study.

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