[11]E. J. Johnson and D. M. Larsen, Phys. Rev. Letters 16, 655 (1966).

[12]See, for example, M. Born and K. Huang, Dynamical Theory of Crystal Lattices (Clarendon Press, Oxford, England, 1954), p. 368.

[13]P. A. Wolff, to be published.

PARITY MIXING IN NUCLEAR HARTREE-FOCK CALCULATIONS*

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(Received 12 December 1966)

It is well known that although the nuclear Hamiltonian may possess a certain symmetry, the Hartree-Fock (HF) Hamiltonian will, in general, not commute with the operator corresponding to that symmetry. Thus, restrictions on the variational single-particle wave functions should not be inferred from properties of the actual many-body wave function. Specifically, it is an assumption, a priori unjustifiable, to restrict the HF wave functions to be eigenstates of the parity operator. Such an assumption is especially dubious in the light of the fact that the tensor force can give no contribution (at least the direct term) if parity is not mixed in the trial wave functions.[1] Previous investigations[2-4] have shown that agreement with the observed spin-orbit splittings and magnetic moments could be obtained from the tensor force if parity mixing were allowed. However, in the HF calculations presently being performed with realistic forces, the lowest energy solution is always found to have single-particle wave functions of good parity. (The details of the HF procedure and the results of these calculations can be found in Baranger.[5])

Since the forces being used in these calculations (the Tabakin[6] and Yale-Shakin[7] forces) are rather complicated, a simplified force was constructed and a detailed investigation of parity mixing was carried out. The force consisted of the central and tensor parts of the Hamada-Johnston[8] force with the infinite hard cores replaced by finite cores of heights 150 and 0 MeV, respectively. The radius of the soft core was 0.67 F. This model force, though unrealistic, contains some of the features of the realistic potential of Bressel[9] and is well suited to a study of the relation between parity mixing and the strength of the tensor force. The nuclei considered were the closed-shell nuclei He^4, O^{16}, and Ca^{40}, and the deformed nuclei Be^8, C^{12}, Ne^{20}, and Si^{28}. The trial wave functions were taken to be linear combinations of harmonic oscillator wave functions:

$$|\lambda\rangle = \sum C\, |nljm\tau_z\rangle$$

with the summation on n, l, and j. Thus, though radial variations, deformations, and parity mixing were allowed, it was assumed that the HF solutions are axially symmetric and eigenstates of τ_z. The variational space included the 1s, 1p, 2s-1d, and 2p-1f shells. The orbits of all A nucleons were treated variationally.

It is worth going into some detail concerning the initial choice of the variational parameters (C). If the total Hamiltonian possesses a symmetry and the C's are suitably chosen for the first iteration, then the HF solutions may also have that symmetry. It was found in our other calculations that even if the C's for $+m$ states and $-m$ states were initially unrelated, the minimal solution was invariant under time reversal, i.e., if

$$|\lambda_+\rangle = \sum_{nlj} C_{nlj}^{m\tau_z} |nljm\tau_z\rangle,$$

and

$$|\lambda_-\rangle = \sum_{nlj} C_{nlj}^{-m\tau_z} |nlj-m\tau_z\rangle,$$

are occupied states, then

$$T|\lambda_+\rangle$$
$$= \sum_{nlj} (C_{nlj}^{m\tau_z})^* (-1)^{j+l-m} |nlj-m\tau_z\rangle = |\lambda_-\rangle,$$

or

$$(C_{nlj}^{m\tau_z})^* (-1)^{j+l-m} = C_{nlj}^{-m\tau_z}.$$

In most HF calculations, the C's are initially chosen to be real. Since the interaction is real, the C's will continue to be real throughout the iteration procedure. This, in fact, is equivalent to the further assumption that the HF solution is invariant under the combined operation of space inversion and rotation of 180° about the y axis, $PR_{\pi y}$. According to a suggestion of Villars,[10] the initial set of C's should be complex in the following way: All "wrong" parity components should enter with pure imaginary C's and all "right" parity components with pure real C's. For example:

$$|\lambda\rangle = C_1 |nljm\tau_z\rangle + iC_2 |n'l+1j'm\tau_z\rangle,$$

with C_1, C_2 real. This is equivalent to invariance under $R_{\pi y}$ as well as T. The bulk of the work reported here followed this latter suggestion. The results are shown in Fig. 1. There binding energy is shown versus α, where α is defined by

$$V = V_{\text{central}} + \alpha V_{\text{tensor}}$$

($\alpha = 1$ corresponds to the Hamada-Johnston strength.)

The dashed lines correspond to the initial choice of (real) C's such that $P|\lambda\rangle = \pm|\lambda\rangle$. (This property is, of course, retained by the solution.) It is perhaps surprising that these energies are independent of α, not only for the spherical nuclei, but for the deformed nuclei Be^8 and Ne^{20}, as well, over a wide range of tensor strengths. The solid lines correspond to the initial C's being chosen real for even parity and pure imaginary for odd parity. The jagged solid line indicates that these results are rather qualitative in that not many points were calculated in that region since that region is not physically interesting. For values of α less than the "transition strength" the parity mixing vanishes regardless of the initial choice of C's. This was verified for pure real and nontrivially complex, time-reversal-invariant and time-reversal-noninvariant C's, as well as the choice described above.

The energy dependence of the deformed nuclei C^{12} and Si^{28} as a function of α is more, as one would expect. For reasonable values of the tensor strength, the energy decreases slowly as the force becomes larger, while if α is allowed to take on very large values, the energy decreases rapidly. Here it is also true that below a certain critical value of

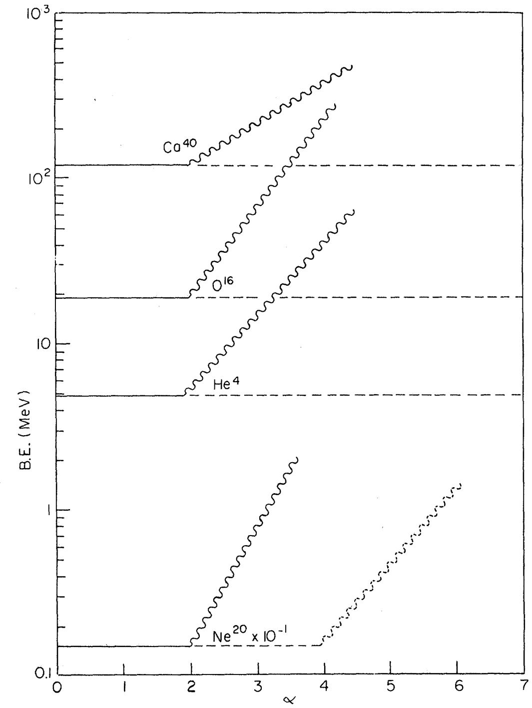

FIG. 1. HF binding energies as a function of the tensor-force strength. The dashed lines correspond to good parity solutions and the solid lines to mixed parity solutions. The two solutions coincide for $0 \leq \alpha \leq 1.9$. The Be^8 solutions have the same characteristics as the Ne^{20}. (The Ne^{20} curve has been scaled down for the sake of clarity.)

α (again approximately twice the realistic strength), there is absolutely no gain in energy achieved by mixing parity, and the lowest solution consists of single-particle wave functions which are eigenstates of parity. The critical value of α is nearly independent of A to the degree that could be tested in the limited subspace to which the variation was restricted.

It is true that spin-orbit splittings can result from the tensor force if its strength is unrealistically large. (This has been considered extensively by Ebenhöh.[4]) The relevant conclusion of this investigation is, however, that unless one is utilizing a force whose tensor part is nearly 100% stronger than the accepted strength, it is sufficient to restrict the single-particle wave functions to be eigenstates of the parity operator when performing Hartree-Fock calculations of nuclear ground-state energies.

*This work is supported in part through funds provided by the Atomic Energy Commission under Contract

No. AT(30-1)-2098.

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[1]K. Bleuler and C. Terreaux, Helv. Phys. Acta 30, 183 (1957).

[2]J. P. Amiet and P. Huguenin, Nucl. Phys. 46, 171 (1963).

[3]J. P. Amiet and P. Huguenin, Nucl. Phys. 80, 353 (1966).

[4]W. Ebenhöh, Z. Physik 195, 171 (1966).

[5]W. H. Bassichis, A. K. Kerman, and J. P. Svenne, to be published. For an exposition of the Hartree-Fock method see M. Baranger, Cargèse Lectures in Theoretical Physics (W. A. Benjamin, Inc., New York, 1963).

[6]F. Tabakin, Ann. Phys. (N.Y.) 30, 51 (1964).

[7]M. H. Hull and C. M. Shakin, Phys. Letters 19, 506 (1965).

[8]T. Hamada and I. D. Johnston, Nucl. Phys. 34, 382 (1962).

[9]C. Bressel, thesis, Massachusetts Institute of Technology (unpublished); C. Bressel, A. Kerman, and E. Lomon, Bull. Am. Phys. Soc. 10, 584 (1965).

[10]F. Villars, private communication.

ETA PHOTOPRODUCTION IN THE REGION FROM THRESHOLD TO 940 MeV*

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(Received 21 November 1966)

We have measured the photoproduction of the η^0 meson in the region from threshold to ~940 MeV. The differential cross section near 90° in the center of mass is characterized by a rapid rise near threshold reaching a maximum value of ~1 μb/sr approximately 40 MeV above threshold. The cross section then decreases to a value of 0.3 μb/sr at a laboratory photon energy of ~900 MeV. Several different angles have been measured at a laboratory photon energy of 790 MeV and the angular distribution is consistent with isotropy. This implies either an S or P wave with $T=\frac{1}{2}$, $J=\frac{1}{2}$ near threshold. Recent phase-shift analyses of pion-nucleon scattering suggest that the η^0 production near threshold is to be identified with the $N_{1/2}^*(1550\text{ MeV})\ S_{11}$ state.

Photoproduction cross-section measurements of the eta, a pseudoscalar meson with positive G parity and mass 548.6 ± 0.4 MeV, have previously been obtained at higher photon energies than the range covered by this experiment.[1] More recently, the Frascati group has extended its measurements to lower energies and there is thus some overlap between the present experiment and the Frascati results.[2] In addition, there are new data in the region from 950-1100 MeV from Heusch et al.[3] However, all these previous experiments measure the η^0 production for decay into a specific channel, viz., $\eta^0\rightarrow\gamma+\gamma$. Comparison with the present results thus requires a knowledge of $\Gamma_{\gamma\gamma}/\Gamma_{\text{total}}$, the branching ratio for the two-photon decay mode as compared with all other channels.

Observation of η^0 photoproduction in the reaction $\gamma+p\rightarrow\eta^0+p$ has been accomplished by using counter techniques with the Stanford Mark III 1.1-BeV linear electron accelerator.[4] Since this process has a two-body final state, a measurement of the proton angle and momentum is sufficient to determine uniquely the incident photon energy.

The experimental set up consists of equipment as follows: The linac electron beam is brought out into the experimental area and passes through two secondary emission monitors which are used to integrate the electron-beam current. These are followed by an air Čerenkov counter used to monitor the beam pulse shape. The electron beam is then incident upon a radiator foil typically ~0.01 radiation lengths to produce the photon beam. The photon and electron beam then pass through a sweeping magnet which deflects the electron beam away from the target. The photon beam then passes through a liquid hydrogen target.

The recoil protons are detected by counters at the exit focus of a 44-in. radius, 90° bend, $n=0$ magnet which has 0.1% momentum resolution up to 700 MeV/c. The detecting counters consist of three backing counters and seven momentum-defining counters which are 1% in momentum width. The pulse-height distributions of the backing counters were carefully monitored and used to reject coincidences due to π^+ and e^+. The momentum defining counters are oriented in the focal plane of the spectrometer such that they are parallel to lines of constant photon energy for the reaction under investigation.