provide for any substitutional invariance if $f_1 \neq f_2$, nor does it allow gauge transformations of any kind resulting in assignments of different quantum numbers to the different members of a given charge multiplet other than those already ascribed to them by their transformation properties in the charge space. This last point, of course, is a necessary requirement which must be imposed on any charge multiplet if the introduction of additional symmetries is to be avoided. $\{B\text{-}K\}$, however, is invariant under the hypercharge gauge transformation. Since $\vec{T}^2$ and T_3, where $\vec{T} = \vec{I} + \vec{K}$, are now the only quantities conserved, the charge operator may be defined as $Q = T_3 + \frac{1}{2} U$, in the manner of d'Espagnat and Prentki,[6] where U, the hypercharge quantum number, has the values 0 for π, N_2, and N_3, +1 for N_1 and K, and -1 for N_4 and K^c, with $U = S + N$. In addition to the conservation of $\vec{T}^2$, T_3, and U (or S), $\{B\text{-}K\}$ also imposes the selection rules $\Delta\vec{K} = 0, \pm 1$; $\Delta I = 0, \pm 1$, which, however, will be satisfied with the former.

In virtue of the fact that our $\{B\text{-}K\}$ imposes no more than charge independence (in the conventional sense) and the strangeness rule, simple amplitude relations of the type derived by Pais can no longer be easily obtained. However, certain processes are now favored with more channels than are others, and the implications of this fact are presently being investigated.

Finally we remark that our $\{B\text{-}K\}$ corresponds to taking $F_1 = -3\ F_2$ in the notation of Pais.[7]

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ELASTIC SCATTERING OF PHOTONS BY TANTALUM*

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The neutron yield cross section for the deformed nucleus Ta^{181} has recently been measured[1,2] and found to consist of two maxima at 12.45 and 15.45.[1] This cross section has been fitted by the superposition of two Lorentz lines. These have been associated with oscillations of the nuclear charge along the one long and two short axes of the ellipsoid according to the suggestions of Okamoto[3] and Danos.[4] The present note describes the results of a measurement of the differential cross section at 120° for the elastic scattering of photons by the tantalum nucleus. The experimental methods employed were described in a previous paper.[5] The results are given in Fig. 1.

For comparison with these results the differential elastic scattering cross section, $d\sigma_S(\theta)/d\Omega$, associated with electric dipole absorption, has been calculated assuming (a) that the nuclear dipole polarizability is a tensor, i.e., that the polarizability has a different value associated with the major and minor axes of the nucleus; and (b) that the nuclear dipole polarizability is

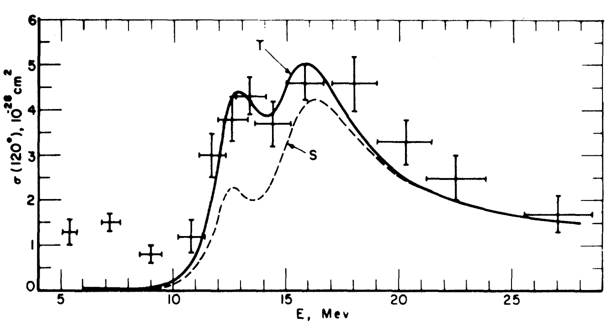

FIG. 1. The differential elastic scattering cross section for tantalum at 120° as a function of photon energy. The smooth curves are calculated using the dispersion relation and the resonance parameters given in reference 1 with the peak cross sections reduced by 10%. The solid curve is the result obtained assuming a tensor electric dipole polarizablity [Eq. (5)]; whereas the dashed curve results from assuming a scalar polarizability [Eq. (6)].

a scalar, i.e., that it is independent of the orientation of the nucleus with respect to the photon's polarization.

In terms of the forward scattering amplitude, f, the relationships between the photon absorption and elastic scattering cross section, σ and σ_S, are

$$\sigma = 4\pi \lambda\!\!\!^{-} \operatorname{Im} f,$$
$$d\sigma_S(0°)/d\Omega = |f|^2. \quad (1)$$

Whenever σ can be represented by a Lorentz line, and the nuclear polarizability associated with σ is along the direction of polarization of the incident photon, the forward scattering amplitude is given by

$$f = \frac{\sigma_0 \Gamma}{4\pi \hbar c} \frac{E^2(E_0^2 - E^2) + iE^3\Gamma}{(E_0^2 - E^2)^2 + E^2\Gamma^2}, \quad (2)$$

where σ_0 is the peak absorption cross section, E_0 the resonance energy, and Γ the full width at half maximum of the resonance.

Assuming a tensor electric dipole polarizability, the scattering amplitude for a given nuclear axis, z, arbitrarily oriented with respect to the polarization of the incident photon, is

$$f(\theta) = f_z \cos\eta \sin\rho, \quad (3)$$

where θ is the scattering angle, η is the angle between the photon's polarization and z, ρ is the angle between z and the direction of observation, and f_z is given by Eq. (2) where the resonance parameters associated with the absorption of photons polarized along z are used. Danos has shown that, for an ellipsoidal nucleus, the relationship between the cross sections σ_a and σ_b associated with the absorption of photons polarized along the major and minor axes and the total absorption cross section for an unpolarized photon beam by an unoriented sample is

$$\sigma = \sigma_a/3 + 2\sigma_b/3. \quad (4)$$

For such an ellipsoidal nucleus the scattering cross section at an angle θ will be given by the absolute square of the sum of the scattering amplitudes associated with the three orthogonal nuclear axes, averaged over all orientations of the nucleus with respect to the incident photon's polarization and finally averaged over all polarizations of the photon. If A and α, B and β, are the real and imaginary parts of the scattering amplitudes associated with σ_a and σ_b, and D is the nuclear Thomson scattering amplitude, $-Z^2e^2/AMc^2$, the resulting expression for the differential scattering cross section is

$$\frac{d\sigma_S(\theta)}{d\Omega} = \frac{3}{15}(A^2+\alpha^2) + \frac{11}{30}(B^2+\beta^2) - \frac{1}{15}(AB+\alpha\beta) + \frac{1}{3}AD + \frac{2}{3}BD + \frac{1}{2}D^2 + \left[\frac{1}{15}(A^2+\alpha^2) + \frac{7}{30}(B^2+\beta^2) + \frac{1}{5}(AB+\alpha\beta) + \frac{1}{3}AD + \frac{2}{3}BD + \frac{1}{2}D^2\right]\cos^2\theta. \quad (5)$$

If, on the other hand, the two resonances found for the neutron yield cross sections for tantalum are not associated with different axes, i.e., if the nuclear polarizability is assumed to be a scalar, the dipole moment induced in the nucleus is always along the polarization of the incident photon. The scattering amplitude for the two resonances is then the sum of the amplitudes for the two resonances with an angular dependence determined by the sine of the angle between the polarization of the incident photon and the direction of observation. The differential scattering cross section is then given by the average of the absolute square of this scattering amplitude over all orientations of the photon's polarization:

$$\frac{d\sigma_S(\theta)}{d\Omega} = \frac{1+\cos^2\theta}{2}\left[\left(\frac{A+2B}{3}+D\right)^2 + \left(\frac{\alpha+2\beta}{3}\right)^2\right], \quad (6)$$

where $A/3$ and $2B/3$, $\alpha/3$ and $2\beta/3$ are the real and imaginary parts of the forward scattering amplitudes associated with the two resonances found in the absorption cross section.

Equations (5) and (6), evaluated at 120° and using the parameters of reference 1, are plotted in Fig. 1. The cross section obtained from the scalar expression [Eq. (6)] shows the effects of strong destructive interference between the two resonances. This interference is essentially missing in the tensor polarizability case since the absorption is now associated with the orthogonal axes of the nucleus. The experimental cross section as a function of energy is in better agreement with the result calculated from the tensor expression than that obtained assuming a scalar polarizability. This result tends to confirm the spatial correlation of the dipole polarizability for a deformed nucleus as proposed by Danos and Okamoto.

In order to achieve the agreement in magnitude displayed in Fig. 1, the peak cross sections of the resonances of reference 1 were reduced by 10%. Since the absolute cross sections of reference 1 are known to only about ±15% and since there are possible systematic errors (as large as ±7%) in the absolute magnitude of the scattering cross section, this adjustment is well within the errors in the absolute sizes of the two cross sections being compared. The difference between the experimental points and the smooth curve at the higher energies is probably significant. It represents more absorption than is given by the two Lorentz lines used to predict the scattering cross section. The neutron yield cross section of reference 1 is also higher than the Lorentz lines in this energy region.

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DOUBLING OF FERMIONS?

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The recent discovery that the long sought π-e mode of decay does indeed exist,[1,2] and that its intensity[3] relative to the π-μ mode is compatible with the assumption of equal interactions for electrons and μ mesons,[4,5] removes a stumbling block to further thinking about the remarkable similarity between these two particles. Except for their mass difference, the electron and μ meson of a given electric charge behave, as far as their known interactions are concerned–i.e., electromagnetic, as well as weak interactions–like two states of one particle. Tentatively we shall assume that they differ by a hitherto unknown internal parameter. This parameter is presumably connected with their mass difference, but does not affect their known interactions which may be held to be invariant with respect to the internal parameter. Many examples of invariance of interactions with respect to an internal parameter are well established, e.g., the charge independence of nucleon interactions. On this view, one might refer to the μ meson as a "heavy electron"–a name by which it was once known–and use the symbols e_1 and e_2 for e and μ, respectively.

One might go further and ask, again assuming the view outlined to be correct, whether the heavier fermions, the baryons, might not also appear as doublets of which only the lighter members have so far been observed. Although the baryons possess strong interactions, in addition to the electromagnetic and weak ones, it is conceivable that whatever the cause of the doubling of electrons into light and heavy ones might be, it might also be responsible for a doubling of baryons. In the absence of definite knowledge on this point, it would seem worthwhile to start out by inquiring why the "heavy baryons," if they actually do exist, might have so far escaped detection.

Take the case of the proton, where besides the ordinary proton p_1, a "heavy proton," p_2, might exist. Simple considerations show that it would have been difficult to recognize such a particle without a deliberate search. What would be some possible modes of production of p_2? For guidance, let us remind ourselves of the processes which give rise to μ mesons. Two essentially different processes are known: μ mesons may appear as decay products of π or K mesons, or they may be produced by photons in pairs: $\mu^+ + \mu^-$. The pair production is difficult to observe, and it was only discovered as a result of a deliberate and careful search.[6] In analogy with these phenomena, we might look for heavy nucleons both in decay processes and in pair production. Depending on the mass of the heavy nucleons, it might or might not be energetically possible for a heavy nucleon to appear as a decay product of one of the known hyperons. However, a process of this type would only have been found with a reasonable probability if the mass difference $\Delta = m(p_2) - m(p_1)$ were considerably smaller than the Q value of hyperons. There remains, as a more systematic approach, the deliberate search for the pair-production process $p_2 + \bar{p}_2$. Here we have the advantage, compared with the case of μ-pair production, that we can make use of the strong interactions, besides the electromagnetic ones.

What would be the characteristics of a p_2, or its antiparticle $\bar{p}_2$ (apart from their mass), by