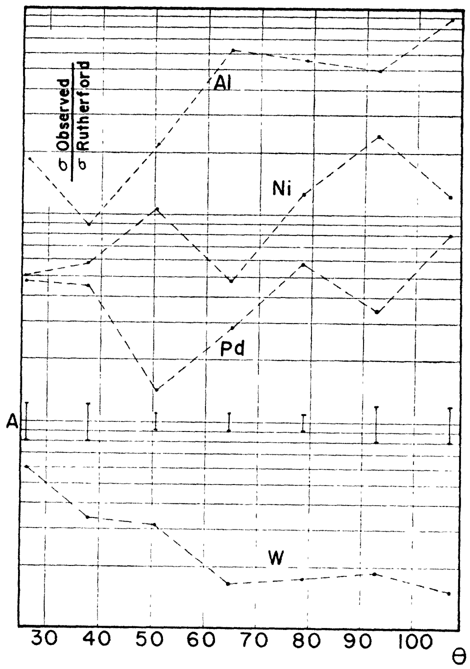

FIG. 1. A plot of the logarithm of the ratio of the observed elastic scattering to Rutherford scattering *vs* the scattering angle θ for 18.6-Mev. protons. The relative vertical position of the four curves has no significance.

through which the protons are scattered. No scale for the ordinates is indicated, since none would be significant. The relative vertical position of the four curves is not significant. They are presented on the same figure merely as a matter of convenience. The vertical lines along the horizontal line A indicate the probable errors in the various observations. The procedure of overlapping sets of observations leads to a larger probable error for the smallest and largest angles.

The W elastic scattering decreases monotonically below the Rutherford dependence; the Ni scattering has an average dependence the same as that expected for Rutherford scattering; and the Al definitely decreases less rapidly than would be expected for the case of Rutherford scattering. In addition, there are considerable deviations for Pd, Ni, and Al about their average angular dependence. A certain amount of progress has been made toward the description of the Al results in terms of a nuclear model which gives a sticking probability dependent on the angular momentum.[2]

By using a pair of cameras very near one another, the cross section for elastic scattering at 78.7±0.5° of a number of elements

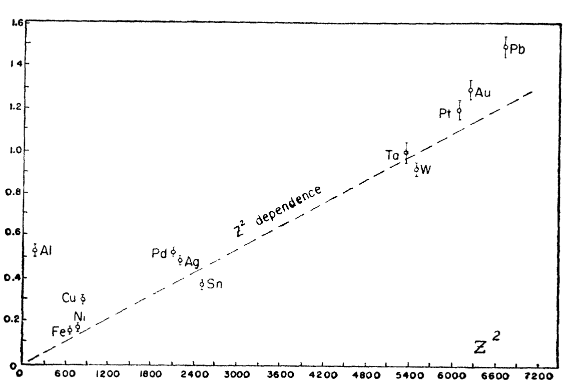

FIG. 2. For scattering at 78.7°, a plot of the ratio of the scattering by various elements relative to W *vs* the square of the atomic number Z.

relative to W was measured. The mass per unit length of the scattering wires was obtained with a quartz torsion microbalance. The reference wire was W in all cases except for four cross check runs. The results of this experiment are shown in Fig. 2. The ratios have probable errors varying from six to nine percent. In view of the marked angular dependence indicated in Fig. 1, the results of the scattering at a fixed angle, shown in Fig. 2, would appear to have little chance of interpretation at the present time.

* This work was supported in part by the joint program of the ONR and AEC.

[1] H. W. Fulbright and R. R. Bush, Phys. Rev. **74**, 1323 (1948).
[2] R. LeLevier and D. S. Saxon (private communication).

Relative Back-Scattering of Electrons and Positrons

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RECENT measurements by Seliger[1] have indicated a significant difference in the back-scattering coefficients, β^- and β^+, for electrons and positrons. For isotropic sources of electrons and positrons with energies of the order of 0.5 Mev, and several scattering media with Z values ranging from 4 to 82, it was found that $\beta^-/\beta^+ \sim 1.3$.

Seliger suggested that this different behavior of electrons and positrons might result simply from their different cross sections for elastic single scattering at relativistic energies.[2] The ratio of the electron and positron cross sections for 90° scattering at 1.7 Mev is as large as 4, but the ratio approaches 1 at lower energies and for small angle scattering. The cumulative effect of small angle single scatterings constitutes the main source of the diffusion and back-scattering of electrons. Therefore, it is not immediately apparent that the difference of electron and positron scattering, while obviously in the right direction, is actually adequate to account quantitatively for the observed effect.

No detailed theory of back-scattering is available, but Bothe[3] has evaluated the back-scattering coefficient β approximately by considerations modeled on the neutron "albedo" theory. In the treatment of Bothe, the main parameter controlling the value of β is the ratio of the "true" range R of a particle in the scattering medium to the "scattering length" λ_s (inversely proportional to $\int_0^\pi \sigma(\theta)(1-\cos\theta)\sin\theta d\theta$), that is, $\beta=\beta(R/\lambda_s)$. Bothe computed λ_s from a nonrelativistic differential single-scattering cross section $\sigma(\theta)$, i.e., in essence a Rutherford cross section, equal for electrons and positrons.

A new evaluation based on relativistic cross sections has been made for the purpose of estimating the magnitude of the effect under consideration. The theoretical data of Bartlett, Watson, and Massey[2] on the cross section for single scattering of 500-kev electrons and positrons in mercury were inserted in the Bothe formulas. The λ_s values for electrons and positrons were calculated by approximating the $\sigma(\theta)$ by simple functions and doing the integral analytically. Since there is no reason to expect an appreciable difference in the ranges of the two types of particles, the range R for 500-kev electrons and positrons in mercury was taken as 0.030 cm.

The resulting values of R/λ_s for erectrons and positrons are 18.1 and 12.4, respectively. Table I shows a comparison of the values of λ_s, R, R/λ_s from the present calculation and of the corresponding values derived from Bothe's nonrelativistic for-

TABLE I. Comparison of calculated values with those of Bothe's theory.

		$1/\lambda_s$ (cm^{-1})	R (cm)	R/λ_s	β
Present calculation	Electrons	$605=1.39Z^2\rho/A$	$0.03=0.163A/Z\rho$	18.1	0.59
	Positrons	$414=0.95Z^2\rho/A$		12.4	0.51
Bothe		$(400/V)^2Z^2\rho/A$ $=0.64Z^2\rho/A$	$1.25\times10^{-6}V^2A/Z\rho$ $=0.31A/Z\rho$	$0.2Z$ $=16$	0.56

mulas. (The Bothe values for λ_s and R are actually not relevant to 500-kev particles. Nevertheless, Bothe was able to apply his method to 680-kev electrons with success, because, as shown in Table I, the Bothe value for the ratio R/λ_s departs from the improved estimate much less than the separate values of R and λ_s.[4])

Estimates of the back-scattering coefficient β were derived from the various values of R/λ_s according to the Bothe theory. The resulting asymmetry, $\beta^-/\beta^+ \sim 1.16$, is roughly comparable to the observed effect. This result is taken as an indication that the differences in single scattering alone may well be responsible for the observed differences in back-scattering. Much better agreement with the experimental results should not have been expected in view of the approximations involved. In particular, it should be noted that the observed values of β correspond to an isotropic source at the boundary of the back-scatterer, whereas the values calculated on Bothe's theory do not apply to an isotropic source but rather to one intermediate between that and a normally incident one. For this reason the observed and calculated values of β cannot be expected to agree in an absolute sense.

I would like to thank Dr. U. Fano for suggesting this investigation and for many helpful discussions.

[1] H. H. Seliger, Phys. Rev. **78**, 491 (1950).
[2] See, e.g., N. F. Mott and H. S. W. Massey, *The Theory of Atomic Collisions* (Oxford University Press, New York, 1945), p. 81ff.
[3] W. Bothe, Ann. Physik **6**, 44 (1949).
[4] See also C. H. Blanchard and U. Fano, Bull. Am. Phys. Soc. **26**, No. 2, G10(A) (1951).

Concerning Certain Anomalous Small Angle Diffraction Effects

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HALPERN and Gerjuoy[1] have pointed out that radiation scattered from a parallelopiped will (on the basis of the usual theory) extend to anomalously large angles when the incident beam is very nearly parallel to any face. Recently, Forrester and Mittenthal[2] have attempted to observe this effect in the diffraction of light by a glass cube immersed in a liquid of nearly equal index. They show that extremely large intensities at these angles should be expected, providing the optical path length through the glass cube is many wavelengths larger than in the surrounding medium. Their experiments fail to reveal this intense scattering, and they suggest several explanations, one of them being that the Born approximation (single scattering) underlying Halpern and Gerjuoy's formulas is inapplicable here.

We wish to point out that smallness of the phase shift through the scatterer is a common criterion for the validity of the Born approximation,[3] and that, further, for the case in which the refractive index of the scatterer differs from the surrounding index by much less than unity, a different type of approximate calculation can be made which is valid for both large and small phase shifts and which reduces to the Born approximation in the latter case. In this method, the radiation field is calculated on a plane situated just beyond the scatterer and perpendicular to the incident beam; from this, the far-field is found by the Kirchhoff form of Huygen's principle. The field on this intermediate plane is approximated sufficiently well by (a) the undisturbed incident field at all points outside the geometrical shadow of the scatterer, and (b) a field inside the shadow having the same amplitude but with its phase modified at each point according to the optical path length on a line straight through the scatterer and reaching that point. These approximations seem rather drastic at first sight, but their validity is demonstrated by work of van de Hulst,[4] who considers spherical scatterers and shows that the rigorous Mie solution for electromagnetic radiation agrees with such computations under similar limitations. It is also easy to show directly that the Born approximation emerges (for scattering angles small compared with 1 radian) when the phase shift is small. This can be shown for any shape of scatterer.

We have carried through such a calculation for the case of a cube with two opposite faces parallel to the incident beam, and we have calculated the differential cross section for scattering in a plane perpendicular to these faces. When the phase shift through the center of the cube is large compared with unity, the differential cross section never exceeds that of an equal sized rectangular aperture by more than a factor of 4. Also, at intermediate values of the phase shift, the cross section is always of the same order of magnitude as that of an equal sized aperture. These results can be readily understood on the basis of a physical interpretation of this mode of calculation, also owing to van de Hulst.[4] The contribution from the radiation in the geometrical shadow corresponds closely, for large cubes, to rays refracted by geometrical optics; the remainder correspond to radiation diffracted around the cube. Here, the field arising from the "ray" portion cannot exceed that of a simple aperture, and the field arising from the diffraction portion equals that of a simple aperture. Their combined intensity, then, depends on relative phase and cannot exceed 4 times the intensity scattered by an aperture.

The negative results of Forrester and Mittenthal are thus to be explained entirely on the basis of the breakdown of the Born approximation in their case.

[1] O. Halpern and E. Gerjuoy, Phys. Rev. **76**, 1117 (1949).
[2] A. T. Forrester and L. Mittenthal, Phys. Rev. **81**, 268 (1951).
[3] D. Bohm, *Quantum Theory* (Prentice-Hall, Inc., New York, 1951), p. 553.
[4] H. C. van de Hulst, *Optics of Spherical Particles* (J. F. Duwaer en Zonen, Amsterdam, 1946). This excellent work deserves wider circulation than it seems to have attained.

Ionization by the Collision of Pairs of Metastable Atoms

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MICROWAVE techniques are used to measure the variation in electron density following the interruption of a discharge.[1] It is found that the electron density actually increases for approximately a millisecond after the electric field is removed from a helium or neon discharge (see Fig. 1). Analysis of the data indicates that the delayed ionization results from the collisions of pairs of metastable atoms, as suggested by T. Holstein.

After the maintaining field is removed from the discharge, ionization by electron impact ceases; however, electrons continue to be produced by metastable-metastable collisions. The electrons, which quickly (<100 μsec) come into thermal equilibrium with the gas, diffuse to the walls ambipolarly.[2] The metastable atoms produced during the discharge are lost by diffusion to the walls, by de-exciting collisions with gas atoms, and by metastable-metastable ionizing collisions. Experimental data indicate that the ionization loss is small compared to the diffusion and collision loss; hence, the metastable concentration is given by

$$M \simeq M_0 \exp(-t/T_m), \tag{1}$$

where M_0 is the metastable concentration immediately following the discharge and

$$1/T_m = (D_m/\Lambda^2) + \nu_d. \tag{2}$$

Here T_m is the metastable mean decay time, D_m is the metastable diffusion coefficient, ν_d is the frequency of de-exciting collisions with gas atoms, and Λ is the characteristic diffusion length of the container.[2] The electron density is given by

$$n = A \exp(-t/T_D) - B \exp(-2t/T_m), \tag{3}$$

where

$$T_D = \Lambda^2/D_a, \tag{4}$$

D_a being the ambipolar diffusion coefficient.[2]

Figure 1 shows typical data for helium. At late times, the pro-