Dynamic Screening of Projectile Charges in Solids Measured by Target X-Ray Emission*

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We report a strong dependence of the Al K x-ray cross section in solid Al targets, for thicknesses ranging from 0.5 to 25 $\mu g/cm^2$, on the charge state of the incident oxygen projectiles, in the energy range 12–68 MeV. The results are analyzed in terms of the dynamic screening of the projectile charge by the electrons in the dense target, yielding a dynamic screening rate $\lambda = 5.5\times10^{14}$ sec^{-1} which is consistent with this approach.

We report first observations on the dependence of the characteristic K-shell x-ray yield of atoms in a solid ($^{27}_{13}Al$) under heavy-ion bombardment on (i) the charge state n (3 to 8) of the incident ions ($^{16}_{8}O^{+n}$), (ii) the projectile energy E_1 (12 to 68 MeV), and (iii) the target thickness (0.5 to 25 $\mu g/cm^2$). Charge-state-dependent effects have just been reported independently in gas targets.[1,2] Our experiments were performed under conditions where Coulomb excitation occurs over large impact parameters (compared to the Al K-shell radius $a_{2K} = Z_{2K}^{-1}a_0$, where $Z_{2K} = Z_2 - 0.3$, Z_2 being the atomic number of the target atom, and a_0 the Bohr radius).[3,4] Therefore, our results have implications for the general problems, of long standing, of projectile charges in solids and their dynamic screening as they apply to stopping power, projectile stripping, x-ray production cross sections, and the question of inter-relating charge states inside solids to those observed after emergency from surfaces.[5]

Figure 1 summarizes our results. The relative Al K x-ray production yields $R(E_1)$ measured in thin-target transmission experiments[3] are shown as a function of the target thickness D for different incident oxygen charge states n. The data converge toward thickness-independent limiting values following penetration through some 10 $\mu g/cm^2$, corresponding to several hundred atomic layers.

Oxygen ions were accelerated by the Brookha-

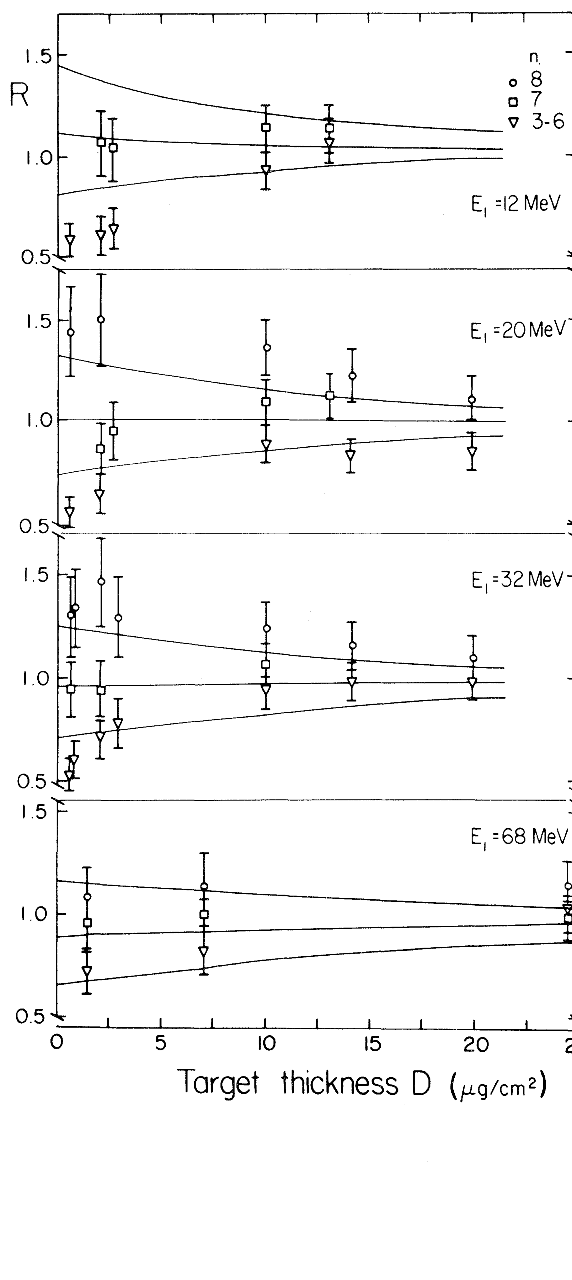

FIG. 1. Variations of the K x-ray yield expressed in terms of $R(n,D,E_1)$, Eq. (1), of Al targets under bombardment with O^{+n}, $n=3$ to 8, as a function of target thickness D for various particle energies E_1. The curves represent Eq. (3), with Eqs. (5) and (6), using the parameter values $\lambda = 5.5\times10^{14}$ sec^{-1} and $\gamma=1$. The upper curve for each energy is calculated for $n = q_1(x=0) = 8$, the intermediate curve for $n=7$, and the lower curve for $n=6$.

ven National Laboratory double tandem Van de Graaff accelerator, with the lowest charge states consistent with the available terminal voltages. In order to increase the projectile charge states, the beam passed through a 20-μg/cm^2 carbon-foil stripper located at the object point of the 90° accelerator analyzing magnet. This provided sufficient beam to stabilize the accelerator for each charge-state run. Moreover, in this mode the charge selection is essentially repeated by the following beam-switching magnet. Typical beam currents of ~1 nA reached the target placed in a 30-in.-diam scattering chamber.

The targets were prepared by vacuum evaporation of Al on ~20-μg/cm^2 carbon foils. Several targets were placed on a movable-target ladder with the Al layers facing the entering beam. A high-resolution Si(Li) x-ray detector (with a 0.25-μm Be window) viewed the target at 45° from a distance of 0.5 m. A set of polyethylene collimator aperatures ensured that the detector could only see Al x rays from the target. A lead cylinder surrounding the detector shielded it from the hard γ rays created at the beam collimators and the beam dump. The target was also viewed by a Si surface-barrier detector at 30°, which was set to register only oxygen particles elastically scattered from the Al target. The elastic-scattering yield provides a normalization basis that is independent of the beam charge state and target thickness. Still, from the comparison of the calculated Rutherford scattering cross sections with the measured beam intensities, the effective Al-target thicknesses could be calculated. They were in excellent agreement with the quartz evaporation microbalance measurements performed during target preparation.

Spectra of x-ray pulses and particle-detector signals were stored simultaneously in 1024 channels of a 4096-channel pulse-height analyzer. At each beam energy E_1 and charge state n, these sets of spectra were measured for several target thicknesses D and for a blank carbon foil. A small x-ray yield near the Al x-ray energy, and scattered particles in the energy range appropriate for kinematic shifts associated with target atoms of mass near 27, were observed with the blank carbon foils. These signals are presumably due to impurities in the carbon foils (perhaps Si), to x rays from the impurities, and possibly to broad-band x rays from the combined O-C atom.[6] They were subtracted as background by normalizing to the Al target measurements on the basis of the integrated beam current.

At each bombarding energy E_1, we calculated the ratio $R(n, D, E_1)$ of x-ray counts N_x and particle counts N_p:

$$R(n, D, E_1) \equiv C(E_1)\left(\frac{N_x}{N_p}\right)_D = \frac{(N_x/N_p)_D}{(N_x/N_p)_\infty}, \tag{1}$$

where the normalization constant $C(E_1)$ was determined as indicated by the last expression in (1) using runs of thicknesses so large ($D \approx$ "∞") that N_x/N_p is independent of n. In this definition, R depends on D only via changes in the effective charge states of the particles as they pass through the target. The trivial linear thickness dependence of the x-ray yield N_x is removed through the division by the elastic scattering yield N_p. The background correction to R is always < 10%. Errors in R vary from 15% for the thinnest targets to 10% for the thicker targets. They are a result of the background correction and the uncertainty in the determination of $C(E_1)$. We have measured the energies of the target K-shell x rays, $E(\mathrm{Al}, K)$, at each target thickness. The data show no dependence on the incident charge state n within the uncertainty of the energy analysis corresponding to ± 5 eV in $E(\mathrm{Al}, K) = 1.48$ keV. From this we conclude that the fluorescence yield at each target thickness is insensitive to the projectile charge states investigated. There is a trend toward higher values of $E(\mathrm{Al}, K)$ in the thinnest targets, by the same amount for all n, if compared with the thicker targets where $E(\mathrm{Al}, K)$ is a constant. Part of this rise is accounted for by the x rays emerging from the carbon backing which becomes unimportant for Al targets of thickness $\gtrsim 10$ μg/cm^2. When we subtract the background, a small residual rise may remain, but our data on these thinnest targets are not sufficiently reproducible to ascertain unequivocally whether this is a real effect. We repeated the experiment with a proton beam and, as expected, found $R = 1$ for all target thickensses.

In our thin targets of atomic density ρ and x-ray self-absorption coefficient μ, the thickness D is so small that $D \ll \mu^{-1}$ and the projectile-energy loss is negligible compared to the incident energy E_1. The x-ray yield per steradian is then given by

$$Y(E_1) = (\rho/4\pi)\int_0^D \sigma_X(Z_1, E(x))\,dx \simeq (\rho/4\pi)\sigma_X(Z_1, E_1)Z_1^{-2}\int_0^D q_1^{\,2}(n, x, E_1)\,dx, \tag{2}$$

where $\sigma_X(Z_1, E_1)$ is the x-ray production cross section by Coulomb excitation proportional to Z_1^2, Z_1 being the charge number of the bare projectile. The mean-square effective projectile charge for this process, q_1^2, depends on the incident charge state n, on the depth of penetration x, and on the projectile energy $E(x) \simeq E_1$. Our experiments measure, in effect, $(1/D)\int_0^D q_1^2(n, x, E_1)\,dx$ because

$$R(n, D, E_1) = \left[\frac{Y(D)/D}{Y(\infty)/D_\infty}\right]_{E_1} = \frac{\int_0^D q_1^2(n,x,E_1)\,dx}{D q_1^2(\infty, E_1)}, \quad (3)$$

where $q_1^2(\infty, E_1)$ is the asymptotic n-independent charge state of the projectile at large depths of target penetration. We note that under conditions where the high-velocity Z_1^3 effect is important,[4] the yield depends on q_1 as

$$Y(E_1) \propto \int_0^D (q_1^2 + \alpha q_1^3)_{E_1}\,dx, \quad (4)$$

where $\alpha(E_1) \ll 1$ is known theoretically[4] for $q_1 \ll Z_2$. When $q_1 \sim Z_2$, higher powers of q_1 may contribute.[7]

One can account for all the observed data, Fig. 1, in terms of only two parameters which have values very close to theoretical estimates, namely, a dynamic screening rate λ and an electron-density enhancement constant γ. We assume that the target electrons form a Fermi gas; that, in constrast to the situation in dilute gases, screening and collision broadening by the electron gas in a solid make the distinction unimportant between the screening of the projectile charge by bound states undergoing electron capture and loss processes, and the screening by the dynamic polarization of the target electron gas; and that a statistical linear-response approach suffices to characterize the dynamic response of the Fermi gas to the moving projectile. In this frame of reference, the appearance of the projectile in the target at $x = 0$ creates a charge imbalance at time $t = 0$. As the particle penetrates into the solid, the particle-target system approaches steady-state conditions at a rate λ. In linear response theory, λ is related to the inverse relaxation time of the electron gas[8] and can be given in terms of the plasma frequency ω_p as $\lambda = (\beta/2\pi)\omega_p$, where β is a number of the order of 0.1. The approach to equilibrium along the trajectory of the particle, then, follows the form

$$q_1(x) = q_1(0)\exp\left(-\frac{\lambda x}{v_1}\right) + q_1(\infty)\left[1 - \exp\left(-\frac{\lambda x}{v_1}\right)\right], \quad (5)$$

where $q_1(0) = n$ is the incident charge state, $q_1(\infty)$ the steady-state charge state, and $v_1 = (2E_1/M_1)^{1/2}$ the projectile velocity. This linear approach neglects asymmetries in λ between $q_1(0) > q_1(\infty)$ and $q_1(0) < q_1(\infty)$. Moreover, we observe no differences between R measured for $n = 3, 4, 5, 6$, presumably for two reasons: The outermost, loosely bound electrons are stripped off at depths much smaller than our thinnest target (0.5 μg/cm^2),[9] and the residual L-shell electron density distribution is screened by the target electron gas. As the projectile penetrates to depths $x > v_1/\lambda$, a steady-state charge distribution establishes it self through a competition between the rate of pertubation set up by the particle moving in the solid with velocity v_1 and the rate of enhancement $\sim\lambda$ of the electron density near the particle.[10] One estimates for $v_1 > v_F$ (v_F being the Fermi velocity of the target electron gas, on the order of the Bohr velocity v_0) that

$$q_1(\infty, E_1) \simeq Z_1(1 - \gamma v_F/v_1), \quad v_1 > v_F, \quad (6)$$

where γ is a constant of the order of 1.

We insert (5) and (6) into (3) and calculate the ratio R for our experimental conditions by fitting the two parameters of the theory. This yields values that are consistent with our approach: $\lambda = 5.5\times10^{14}$ sec^{-1} and $\gamma = 1$. Setting $\lambda = (\beta/2\pi)\omega_p$, where $\hbar\omega_p(\mathrm{Al}) = 15$ eV,[8] we obtain $\beta = 0.15$. Since $\hbar\omega_p$ is nearly the same for most solids,[8] the equilibration rate should be insensitive to the target material. It follows that the equilibrium projectile charge state is reached in solid targets of thickness $D \gg v_1/\lambda \simeq 4\times10^2 v_1/v_0$ Å, where $v_1/v_0 = [40E_1(\mathrm{MeV})/M_1(\mathrm{amu})]^{1/2}$.

The results of the experiments and the theory agree within the uncertainties, as shown in Fig. 1. Use of (4) with theoretical $\alpha(E_1)$ values[4] improves the fit somewhat by increasing the curvature for $x \lesssim 5$ μg/cm^2 closer to the trend exhibited by the data.

In summary, inner-shell excitations by heavy charged particles in solids at high velocities v_1 (compared to $v_{2K} = Z_{2K}v_0$) depend on an effective charge of the projectiles which must be considered in extracting cross sections for inner-shell Coulomb excitation from studies of the ensuing characteristic x rays or Auger cascades. The effective charge states in a solid can be described in terms of the dynamic screening of the moving particle charge by the target electron gas. After a response time of $\sim 2\times10^{-15}$ sec, a steady state is reached which, for the experiments reported here, is given by $Z_1(1 - v_F/v_1)$. This result is consistent with the notion of transient fields acting on ions in solids.[10] It predicts, for example, that the stopping power of thin targets for heavy projectiles should be thickness dependent as long as $D \lesssim 4\times10^2 v_1/v_0$ Å.[11] The steady-state charge states of oxygen producing Al K x rays are somewhat higher than the average charge states deduced empirically from stopping-power measurements.[12] In light of the different domains of impact parameters relevant for these two processes, this trend is reasonable. It suggests that the effective charge states of projectiles in solids reflect the importance of different averages over the bound and dynamic electron cloud that screens the particle charge, depending on the phenomena studied.

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Plasma Diffusion across a Magnetic Field Due to Thermal Vortices

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A convergent cross-field diffusion coefficient due to flute-type thermal fluctuations (vortices) is derived, taking into account the nonlinear, orbit-diffusion effect on the dynamic screening of the interaction and the fluctuation spectrum. Because of the finite-gyroradius reduction of the correlation, there is a region where the diffusion coefficient depends on B only logarithmically.

In a magnetized plasma, the magnetic field line can be charged uniformly ($k_{\parallel}=0$) due to thermal fluctuations. In a strong magnetic field, these flute-type fluctuations, termed thermally excited vortices, can contribute significantly to the plasma diffusion across the magnetic field in systems with finite length L along B. For gyrofrequency $\Omega = eB/mc$ much greater than the plasma frequency $\omega_p = (4\pi ne^2/m)^{1/2}$, the diffusion coefficient is inversely proportional to the magnetic field strength B.[1,2] In the computer simulation, Dawson, Okuda, and Carlyle[3] found that, for $1 < \omega_p/\Omega < (n\lambda_D{}^2 L)^{1/3}$ where λ_D is the Debye length, there is a novel region where the diffusion coefficient is essentially independent of B (plateau region). In this paper, we derive the cross-field diffusion coefficient of a three-dimensional plasma due to the vortices from the kinetic theory of a hot plasma, adopting the test-particle ap-