bring the ϵ/β^+ ratios into accord with theory. The sum of the γ feeding in column 8 is about 10%. To this we add the 1% γ feeding[5] to the 329.5-keV state (not fed in β decay since the transition is second-forbidden) and the sum of the ground-state γ transitions. The latter can be estimated from a statistical model to be about 6%. The prediction for the intensity missing in the decay scheme of Ref. 1 is thus about 17%. This is in very good agreement with the value of about 20% found in the analysis of the present experiments, irrespective of the fact that the numerous γ rays, which enter this fraction, cannot be placed in a proper position in the level scheme.

It is concluded that the introduction of anomalous ϵ/β^+ ratios in the ^{145}Gd decay is not necessary for the understanding of the experimental findings of Ref. 3. This experiment also emphasizes the necessity of selecting more simple cases for rigorous tests of the theory used in the computation of ϵ/β^+ ratios.

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δ Rays from *K*-Shell Ionization Induced by Heavy Ions

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The spectrum of high-energy δ rays emitted in an 83-MeV (^{16}O, ^{197}Au) collision is measured. The data are compared with the scaled Born approximation and the adiabatic monopole model.

Inner-shell electrons ejected in atomic collisions with continuous energy distribution (known as δ rays) have been investigated since the early days of nuclear physics.[1] δ rays were observed with an energy far higher than the maximum energy which can be transferred classically to an electron at rest in a head-on collision.[2] These high-energy tails result from the high components of the momentum distribution of the initially bound electrons. So δ-ray spectroscopy may be used as a method for the investigation of the momentum distribution of strongly bound electrons. Until now, most experiments have been performed with fast, light projectiles[3-5] in conjunction with internal-conversion–electron measurements following Coulomb excitation or nuclear reactions, in which δ rays supply a strong background.

Our interest in studying δ-ray spectra started with the advent of the latest generation of heavy-ion accelerators, with which U-U collisions close to the Coulomb barrier can be studied. These collisions are adiabatic with respect to the velocity of the inner-shell electrons. Thus a quasiatom with charge as high as $Z = 184$ may be formed transiently. It is expected that the highest energy tails of the δ-ray spectra provide directly the momentum distribution of the strongest bound electrons of an atom in which quantum electrodynamics of very high fields plays an important role.

In order to establish δ-ray spectroscopy as a method to study electron wave functions in transiently formed quasiatoms we examined the high-energy δ rays from the *K* shell produced in an 83-MeV ^{16}O collision with ^{197}Au. In order to extract only *K*-shell δ rays the electron spectrum was measured in coincidence with the *K* x rays from gold.

δ rays were produced by bombarding a 150-μg/cm^2 gold foil with ^{16}O ions in the 8^+ charge state accelerated by the Garching tandem Van de Graaf facilities up to 83 MeV. The beam current was measured in a deep well-shielded Faraday cup and ran to 10 nA on the average. Also for the purpose of monitoring, ions elastically scattered were detected at 30° relative to the beam axis with a 100-μm-thick Si surface-barrier detector.

The target was a self-supporting ^{197}Au foil, positioned at 30° to the beam direction and perpendicular to the average emission direction of the detected electrons. The momentum of the δ rays was analyzed with an iron-free orange-type spectrometer[6] mounted with its axis perpendicular to the beam as described by Metag *et al.*[7] In order to define the δ-ray momentum direction the entrance aperture was completely covered by an 0.4-cm-thick aluminum plate, leaving only a hyperbolically shaped opening at 30° relative to the beam direction subtending angles of ± 5°. The transmission efficiency was 0.04 of the total detection efficiency of the spectrometer (of 0.16) and was determined before and after the runs with a Th-B and ^{207}Bi source, respectively. An exit aperture set in the focal plane allowed the detection of electrons within a momentum band $\Delta p/p = 0.03$ with an 0.01-cm-thick cone-shaped plastic scintillator connected via a short light pipe to an RCA8850 photomultiplier tube. The magnetic field of the spectrometer was swept repeatedly up and down from 645 to 1450 G cm in 1000 equal steps of 0.8 G cm. The measuring time after each field adjustment was normalized to the beam intensity by counting a constant rate of elastically scattered ions.

The δ electrons were detected in coincidence with the target K x rays which were measured by an 0.6-cm-thick NaI scintillation counter mounted with its surface 2.3 cm away from the target. In order to remove the target L x rays, Cu (0.02 cm) and Al (0.1 cm) absorbers were placed in front of the detector. The overall K-x-δ coincident detection efficiency and its dependence on electron momentum (i.e., on magnetic-field setting) in the region of interest was extrapolated from an off-beam measurement using a calibrated ^{207}Bi source in the target position and the identical setup and electronics.

To generate the coincident-electron momentum spectrum, windows were set digitally on the prompt peak of the time distribution (11 ns full width at half-maximum, due to the different length of the electron trajectories) and on the K-x-ray peak. The accidental coincidences and the γ background were carefully estimated and subtracted.

Another background contribution may result from a double collision (it is possible within the time resolution) where a K hole leading to K-x-ray emission is produced in the first collision and the detected electron originates from higher-shell ionization in the second collision. This background, however, can be neglected since the corresponding cross section is at most a few percent of the cross section for the one-collision process in the electron-energy region of interest.

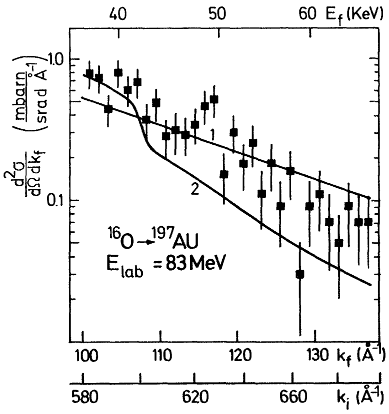

FIG. 1. Cross section for K ionization as a function of the momentum k_f (energy E_f) of the emitted electron at $\vartheta_{lab} = 30°$. Curve 1 is the scaled Born approximation, curve 2 the monopole model.

Figure 1 shows the double-differential cross section $d^2\sigma/d\Omega dk_f$ for K-δ-ray production as a function of the wave number of the emitted electron. The broad peak-shaped structure at 115 Å^{-1} is probably caused by the Auger KLL-transition group after double K-shell ionization. (A single internal-conversion line should have a smaller width.)

Since the ion velocity v is much smaller than the orbiting velocity of the target K-shell electron, the collision process can be described semiclassically within the adiabatic approximation. The transition amplitude for direct Coulomb ionization in a slow collision is given to first order by[8]

$$a_{fi} = -\int_{-\infty}^{\infty} dt \left\langle f \left| \frac{\partial}{\partial t} \right| i \right\rangle \exp\left[\frac{i}{\hbar}\int_0^t dt'(\epsilon_f - \epsilon_i)\right]. \quad (1)$$

Both final (f) and initial (i) states are time-dependent electron states of the transient molecule. The molecular-orbital (MO) energies ϵ_i and ϵ_f will also vary during the collision. Since the two levels are well separated the transition amplitude is very small. We evaluate (1) for δ electrons ejected from the K shell of the heavier collision partner.

The excitation of inner-shell electrons in slow ion-atom collisions takes place at small internuclear distances R. From the energy and momentum conservation laws, one obtains the momentum transferred to the electron,

$$q=\left[2K_i^2(1-\cos\theta)+\left(\frac{\Delta E}{\hbar v}\right)^2\cos\theta\right]^{1/2}, \tag{2}$$

valid for small changes of the initial relative momentum K_i of the nuclei. ΔE, which is the sum of ionization and kinetic energies, is the energy transfer to the electron, and θ is the c.m. scattering angle due to the electron emission. At $\theta=0$ the momentum transfer has its minimum at $q\equiv b_0^{-1}=\Delta E/(\hbar v)$. This implies that the ionization takes place at impact parameters $b\lesssim b_0$. Furthermore, only those parts of the path contribute where the internuclear distance is of the order of b_0, which is considerably smaller than the K-shell radius of the united atom (UA). Therefore, (1) can be evaluated within the UA configuration.

The coupling responsible for Coulomb excitation in an asymmetric collision is given[9] by

$$\frac{\partial}{\partial t}=(\epsilon_f-\epsilon_i)^{-1}\dot{\vec{R}}\nabla_R\frac{Z_Pe^2}{|\vec{r}-\alpha\vec{R}|}, \tag{3}$$

where $\alpha=Z_T/(Z_P+Z_T)$, and Z_P and Z_T are the charge numbers of the (light) projectile and the (heavy) target, respectively. It is easy to calculate the transition amplitude using the UA wave functions and energies. The result is identical with the usual Born approximation, except that a factor $1/\alpha^2$ appears which reflects the fact that the UA wave function correlates to the center of charge rather than to the heavy target. The cross section for the ejection of an electron into a continuum state with momentum $\hbar k_f$ and into the solid angle $d\Omega$ is then given by

$$\frac{d^2\sigma}{dk_f d\Omega}=\left(\frac{2Z_Pe^2}{\hbar v\alpha^2}\right)^2 k_f^2\int_{q_{\min}}^{\infty}\frac{dq}{q^3}\int_0^{2\pi}d\varphi_q|\langle f(\vec{k}_f)|e^{i\vec{q}\cdot\vec{r}}|i\rangle|^2, \tag{4}$$

with $q_{\min}=1/\alpha b_0$ and φ_q the azimuthal angle of $\vec{q}$. In deriving (4) we have assumed a straight-line ion-atom collision with a relative momentum much larger than $\hbar q_{\min}$. For $\alpha=1$ Eq. (4) reduces to the Born expression.[10] The UA initial state is normalized to unity, the UA final Coulomb continuum state to a δ function $\langle f(\vec{k})|f(\vec{k}')\rangle=\delta(\vec{k}-\vec{k}')$.

To include the molecular behavior of the electron during the collision we calculate the ionization cross section in the monopole approximation[11] where the transition operator (3) is replaced by the time derivative of a monopole potential with time-dependent nuclear charge $V(r)=-Z(R)e^2/r$, and the electronic states are eigenfunctions belonging to V. The ground-state energy $E_{1s}(R)$ can be calculated by use of the charge-cloud model[12,13]; the effective charge $Z(R)$ then follows from the relativistic hydrogen energy formula. In the above model the electron is described by a homogeneously charged cloud of radius r_0 centered a distance x apart from the target nucleus. $E_{1s}(R)$ is obtained by minimizing the energy of this system with respect to r_0 and x. For small R, x is given by the center of charge $Z_P/(Z_P+Z_T)$ which is consistent with the choice of the electron coordinate in (3). $E_{1s}(R)$ is in good agreement with perturbation theory (for small R) and with two-center calculations.

The monopole potential leads only to s states of the ejected electron which are important for large momentum transfer. Comparing the contribution to the cross section from all final partial waves with the s state, we find from (4) a factor of 4 at $E_f=35$ keV which decreases monotonically to 3 at $E_f=65$ keV for $\vartheta_{\text{lab}}=30°$ by which the monopole approximation underestimates the cross section.

Both calculations were performed with nonrelativistic hydrogen wave functions. We estimated the relativistic corrections by evaluating the matrix element $\langle f|\dot{Z}e^2/r|i\rangle$ with relativistic wave functions choosing $Z=Z(R\approx a_0/Z_T)$ with a_0 the Bohr radius, where $\dot{Z}(R)$ is peaked. We find an enhancement of the cross section[14] by a factor of 2 in the energy region under consideration.

Figure 1 shows the cross section for K-electron emission calculated with the scaled Born approximation defined by Eq. (4) and with the monopole model described above. Relativistic corrections are included and, in the monopole model, the enhancement factor from the $l\neq 0$ electronic states. By comparison with experiment, we find that the scaled Born approximation gives an overall description of the energy dependence, while the monopole approximation may account for the slope in the low-energy region of the spectrum. The energy dependence of the monopole model is correlated to the behavior of $\dot{Z}(R)$.

In conclusion, we would like to point out that δ rays from the high components of the momentum distribution of K-shell electrons can be observed in heavy-ion collisions. From the relation $\vec{k}_i=\vec{k}_f$

$-\vec{q}$ with q from (2) one can extract the initial momentum k_i of the electron from a measurement of $\vec{k}_f$. This is true since the dominant contribution to the cross section comes from $q \approx q_{\min}$ which means that $\vec{q}$ lies in the beam direction. We detected electrons with k_i in the region 580 $\text{Å}^{-1} \lesssim k_i \lesssim 690\ \text{Å}^{-1}$ which is approximately four times the Bohr momentum of a K electron for $Z = 87$. The measured spectrum can be explained by means of the K-shell momentum distribution of the quasimolecule.

Since the efficiency of the β spectrometer can be improved by mounting it axially to the beam and using a larger counter in the focal plane, the sensitivity of the setup would be sufficient for cross sections which are more than an order of magnitude smaller than those in the present measurement. Thus it might be possible to study the electronic properties of the U + U quasimolecule, as planned by a Heidelberg-Munich collaboration in order to study QED effects in high fields.[15,16]

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Nanosecond Time-Resolved Spectroscopy of the n=2 Levels in a High-Pressure He Discharge

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Nanosecond time-resolved absorption spectroscopy is used to study the populations of the 2^1P, 2^3P, 2^1S, and 2^3S levels in a high-pressure He discharge. At high pressures, the population of the 2^1P level decays primarily through three-body molecular formation involving two ground-state He atoms at a rate of $(3.5 \pm 0.8) \times 10^{-31}$ cm^6 sec^{-1}. The population of the 2^3P level decays primarily as a result of mechanisms other than collisions involving ground-state He atoms.

We report an experimental investigation of the time dependence of the populations of the 2^1P, 2^3P, 2^1S, and 2^3S levels in the afterglow of a pulsed He discharge at pressures of 50–3000 Torr. The populations of the metastable 2^1S and 2^3S levels have long been recognized to play an important role in the kinetics of the He afterglow.[1] The populations of the 2^1P and 2^3P levels are also quite important in the afterglow of a fast pulsed discharge. The fast pulsed discharge is characterized by relatively high pressures, a short voltage rise time, and a high value of the electric field divided by the pressure at breakdown. This leads to high initial densities of excited states and free electrons in the afterglow.

Immediately following the pulsed discharge, the populations of all the $n = 2$ levels are comparable. Radiation trapping can maintain significant populations of the 2^1P and 2^3P levels for many radiative lifetimes. The kinetics of these levels is not well understood, but is interesting and important for many reasons. For instance, it has been