Observation of Motional-Field–Induced Ripples in the Photodetachment Cross Section of H^-

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A ripplelike structure in the photodetachment cross section of H^- near threshold, arising from autocorrelation in the wave function of the photoelectron in the presence of motional electric fields, is observed.

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We have observed structure in the photodetachment cross section near threshold of H^- in the presence of a motional electric field. This structure, which may be described as "ripples," appears under π-polarized laser light—that is, when the plane of the electric field of the incident laser beam is parallel to the direction of the motional electric field in the rest frame of the H^- system. The photodetachment cross section under σ polarization, when the electric field in the laser beam is perpendicular to the direction of the motional electric field, rises smoothly with energy, much like the zero-field case.[1]

Ripples in the cross section in the presence of an electric field have been predicted by Reinhardt[2] and discussed by Overman[3] on the basis of a time-dependent autocorrelation in the outgoing wave function.

The technique used in observing the ripples is similar to earlier work at the Clinton P. Anderson Meson Physics Facility (LAMPF).[4] We used to our advantage the relativistic kinematics of a laser beam incident upon an 800-MeV H^- beam. In particular, smooth Doppler tuning was possible by changing the intersection angle, α, between the two beams, defined such that $\alpha=0°$ when the laser beam is head on to the ion beam. The photon energy, E, in the ions' rest frame is related to the laboratory photon energy E_L by

$$E=\gamma E_L(1+\beta\cos\alpha), \tag{1}$$

where γ and β are the usual relativistic parameters, 1.852 and 0.842, respectively, for an 800-MeV H^- beam. At this energy, for a very modest laboratory magnetic field, $\mathbf{B}$, a very large motional electric field, $\mathbf{F}$, is present in the frame of the ions.

After the interaction the photodetached H^0 atomic beam was magnetically separated from the parent H^- beam and directed into a scintillation counter. A relative cross section, σ, was calculated from R, the rate at which H^0's were detected, by use of the relationship

$$\sigma=GR\beta\sin\alpha/IJ(1+\beta\cos\alpha), \tag{2}$$

where I and J are the ion and photon currents, respectively. G, a geometrical factor, which represents the spatial and temporal overlap of the two beams, was treated as an arbitrary constant.

There are two important differences from our earlier work[4]: (1) The ytrrium-doped aluminum garnet laser was operated in a non–Q-switched mode for which roughly the same number of photons is released in 150 μs that, in the Q-switched mode, are dumped in 10 ns. This change greatly reduced saturation and timing problems at the expense of a decrease in the signal-to-noise ratio. (2) The laser beam, incident on the ion beam, was parallel to the imposed magnetic field. To keep the magnetic field parallel to the laser beam as the angle α was changed, the magnet was mounted on the rotating platform that held the laser beam-bending mirrors. A schematic of the intersection region is shown in Fig. 1. In this geometry

$$F=\gamma\beta cB\sin\alpha. \tag{3}$$

The laser operated at its fundamental wavelength, 1.06 μm, about 1 J per pulse, 10 pulses per second, and was timed to overlap LAMPF macropulses with an average beam current of about 1.0 nA. Because the inter-

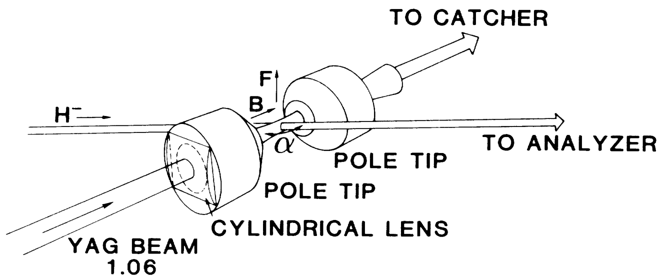

FIG. 1. Schematic representation of the interaction region. The magnetic field **B** is kept parallel to the laser direction so that states for which the laser light is pure π and σ in the center-of-mass system can be prepared.

section angle α was near the Doppler-free angle, $\cos^{-1}(-\beta)$ (147° at 800 MeV), the energy resolution was limited by angular divergence of the laser beam, and was determined, by the use of transitions in the H^0 Paschen and Brackett ($n=3$ and $n=4$) series, to be about 1 meV, FWHM.

Photodetachment data under π polarization for two values of the imposed magnetic field are shown in Fig. 2. A regular pattern of ripples can be clearly seen. For corresponding fields under σ polarization there is no reproducible structure and little change from the zero-field case.

A simple stationary-state picture can explain the presence of ripples. We assume, in what follows, that the relatively weak magnetic field in the barycentric system plays no essential role. In π polarization, the electrons tend to be ejected along the direction of the external field, whose potential forms a sloping barrier from which the wave function can reflect and interfere with the part of the wave at the origin heading away from the barrier. The result is essentially a two-beam interference pattern. In σ polarization, the waves are directed parallel to the sloping barrier so that no reflections occur and, hence, no ripples.

The ripple pattern can be understood in more detail by considering the final state ψ of the electron to be the solution of the one-dimensional Schrödinger equation in a constant electric field F, so that

$$(\hbar^2/2m)\,d^2\psi/dx^2+(E-eFx)\psi=0. \tag{4}$$

By our writing

$$x=bz+E/eF, \tag{5}$$

with $b^3=\hbar^2/2meF$, we have

$$d^2\psi/dz^2-z\psi=0. \tag{6}$$

If we require that ψ vanish as $z\rightarrow\infty$, the solution to this equation is just the Airy function $\mathrm{Ai}(z)$, shown plotted in Fig. 3.

The matrix element M from the initial state $\phi(x)$ to the final $\psi(x)$ is proportional to

$$\int\psi^*(x)x\phi(x)\,dx. \tag{7}$$

Since $\phi(x)$ is a well-localized symmetric function peaked around $x=0$, M is small when ψ is symmetric near $x=0$, namely when the extremum values (maxima and minima) of $\mathrm{Ai}(z)$ occur at $x=0$. The values of z when $\mathrm{Ai}(z)$ is maximum or minimum are denoted a_s' where $s=1,2,\ldots$.[5]

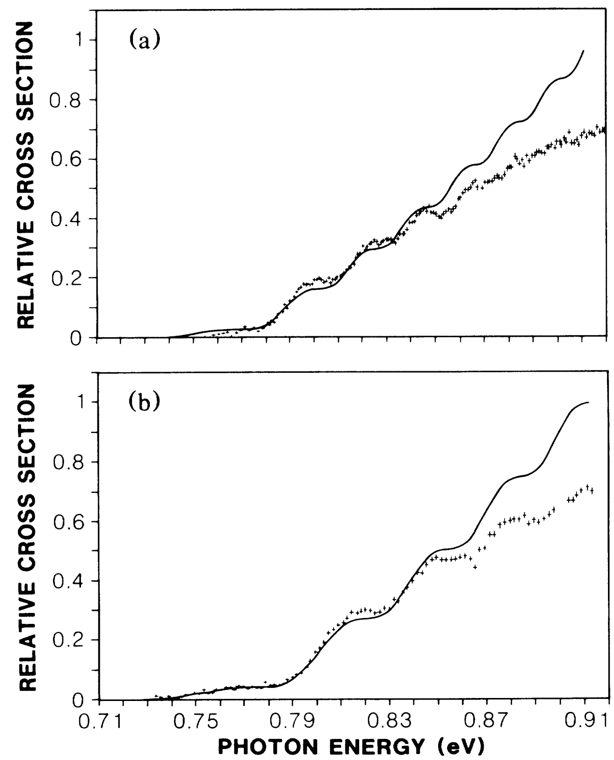

FIG. 2. Relative cross sections for π polarization compared with the theory of Reinhardt and Overman (solid curve). (a) Laboratory field of 300 G. (b) Laboratory field of 470 G.

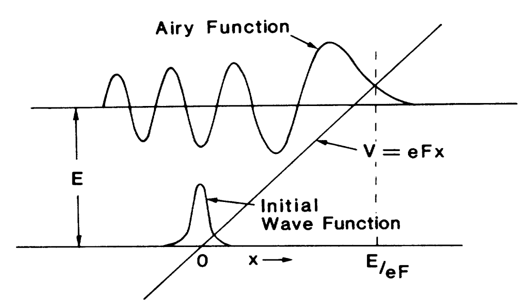

FIG. 3. The absorbed photon causes transitions from the initial bound wave function to that of an electron in a constant electric field. E is the energy above zero-field threshold. The classical turning point for the ejected electron in the constant field F is a distance E/eF from the center of the atom.

From Eq. (5) we can therefore write that E corresponding to the sth dip in the cross section is

$$E = -(eF)^{2/3}(\hbar^2/2m)^{1/3}a_s'. \tag{8}$$

The photon energy required to make a transition to a minimum is E plus the electron affinity of H^- (0.7542 eV). This simple picture predicts well the energies of minima in Fig. 2. Scaling of structure according to $F^{2/3}$ seems roughly to be borne out by the data.

The fading out of the ripples for photon energy $\gtrsim 0.88$ eV in Fig. 2(a) can be explained also in terms of the interference of a wave reflected from the barrier and one emerging from the atom (see Fig. 3). If the time to go from the atom to the classical turning point and back exceeds the coherence time for the detachment process, then there can be no interference. The time to go to the barrier and back is

$$\tau = 2\int_0^{E/eF} dx/v = (8mE)^{1/2}/eF. \tag{9}$$

For the case in Fig. 2(a), with $E=0.13$ eV and $F=100$ kV/cm, we compute $\tau \sim 2.4\times 10^{-13}$ sec. A coherence time of this magnitude corresponds to an energy uncertainty of 1.4 meV, which is consistent with our expected energy resolution. We see that this system behaves as an "atomic interferometer" from which one can determine the coherence time of the photon in the center of mass frame.

We note that the threshold energy under σ polarization is well defined and close to the zero-field theoretical value of 0.754 eV. Under π polarization [Fig. 2(b)], on the other hand, there is appreciable cross section for photodetachment as low as 0.720 eV, whose slope with energy changes before the ripples begin.

The success of this simple scaling law based on the free motion of a charged particle in a one-dimensional electric field, neglecting the final-state interaction between the outgoing electron and the residual H atom, suggests that a calculation of the photodetachment cross section itself at this same level of approximation would be worthwhile. The detachment cross sections were calculated as the Fourier transform of the dipole autocorrelation function,[2,3]

$$\sigma(\omega) \propto \int_{-\infty}^{\infty} dt\, e^{i\omega t}\langle \psi_{\text{bound}} | \boldsymbol{\mu}(t)\cdot\boldsymbol{\mu}(0) | \psi_{\text{bound}}\rangle = \int_{-\infty}^{\infty} dt\, e^{i\omega t}\langle \phi(t) | \phi(0)\rangle, \tag{10}$$

where $\boldsymbol{\mu}$ represents the dipole operator.

In the final expression $|\phi(t)\rangle$ is $|\phi(0)\rangle$ time evolved via the final-state Hamiltonian, including interaction with external dc fields. This is a convenient approach as closed-form expressions for the Feynman propagators for free motion in the y and z directions and motion in a linear potential in the x direction are well known.[6] The initial state, $\phi(0)$, which is the dipole operator times the assumed initial ground state, was taken to be of the generic form

$$x\exp[-(x^2+y^2+z^2)/a_0^2], \tag{11}$$

which is appropriate for an outgoing p wave with π polarization. These assumptions allowed closed-form evaluation of the dipole correlation function as

$$\langle \phi(t) | \phi(0)\rangle = e^3(t)e'(t)\exp\left[\frac{F^2t^2}{24}(-3-it)\right], \tag{12}$$

where

$$e(t) = (1+it)^{-1/2}$$

and

$$e'(t) = \left[\frac{-F^2t^3}{2(2t-i)} + \frac{1}{1+it} - \frac{F^2t^2}{4i(2t-i)}\right],$$

and where atomic units ($e^2=a_0=\hbar=m_e=1$) are used. The detachment cross sections were then found by direct numerical transformation of Eq. (12). This is a very efficient organization of the calculation when compared to calculation of the stationary-state differential angular cross section in the same Cartesian coordinates, where y and z motion would be described by plane waves, and the x motion by the appropriate Airy function. This latter calculation would yield the partial cross section for production of an outgoing electron with a specific vector momentum and would of necessity be followed by numerical integration to obtain the integral detachment cross section (photoabsorbtion cross section) desired here. The two methods are, of course, physically identical.

Results of calculation of the detachment cross section are shown in Fig. 2 for two representative values of the applied external magnetic field B. The results are seen to be in remarkable qualitative agreement with those experimentally observed. Theory and experiment were arbitrarily normalized to optimize agreement near threshold. Detailed comparison awaits the inclusion of the appropriate electron-atom final-state interaction and inclusion of a more appropriate initial bound-state wave function, which are both easily accomplished within the time-dependent formalism.[7]

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