the first two states both have strong decays to the ^{20}Ne 4^+ state and the 21.3-MeV state also has a weak γ-decay branch that feeds the ground-state band. It is possible that the high-spin selectivity and the enhancement of the 21.3-MeV state at $\theta_{lab} = 10°$ and 20° reflect a ^{28}Si structure (the inelastic channel also exhibits structure[20] at $E_{c.m.} \simeq 26.7$ MeV). Further experiments are needed to make definitive spin assignments, and the enhanced cross section exhibited at this anomaly should facilitate studies.

The authors wish to thank N. R. Fletcher, J. Philpott, and D. E. Gustafson for their comments and acknowledge support of the National Science Foundation.

(a) Present address: Noise Pollution Division, Department of Transportation, Tallahassee, Fla.

[1] Y. Akiyama, A. Arima, and T. Seke, Nucl. Phys. A138, 273 (1969).
[2] D. Branford, M. J. Spooner, and I. F. Wright, Part. Nucl. 4, 231 (1972).
[3] L. K. Fifield, R. W. Zurmühle, and D. P. Balamuth, Phys. Rev. C 8, 2217 (1973), and references therein.
[4] G. J. KeKelis and J. D. Fox, Phys. Rev. C 10, 2613 (1974).
[5] P. Sperr, D. Evers, K. Rudolph, W. Assman, E. Spindler, P. Konrad, and G. Denhöfer, Phys. Lett. 49B, 345 (1974).
[6] E. R. Cosman, T. M. Cormier, K. Van Bibber, S. Sperduto, G. Young, Y. Erskine, L. R. Greenwood, and O. Hansen, Phys. Rev. Lett. 35, 265 (1975), and references therein.
[7] N. R. Fletcher, J. D. Fox, G. J. KeKelis, G. R. Morgan, and G. A. Norton, Phys. Rev. C 13, 1173 (1976).
[8] A. Watt, D. Kelvin, and R. R. Whitehead, Phys. Lett. 63B, 385 (1976).
[9] A. H. Lumpkin, G. R. Morgan, K. W. Kemper, and D. E. Gustafson, Bull. Am. Phys. Soc. 21, 966 (1976).
[10] K. R. Chapman, Nucl. Instrum. Methods 124, 299 (1975).
[11] J. L. C. Ford, Jr., J. Gomez del Campo, R. L. Robinson, P. H. Stelson, and S. T. Thornton, Nucl. Phys. A226, 189 (1974).
[12] M. Sakai, Nucl. Data Tables 8, 323 (1970).
[13] P. Charles, F. Auger, I. Badawy, B. Bertheier, M. Dost, J. Gastebois, B. Fernandez, S. M. Lee, and E. Plagnol, Phys. Lett. 62B, 289 (1976).
[14] D. R. James, G. R. Morgan, N. R. Fletcher, and M. B. Greenfield, Nucl. Phys. A274, 177 (1976); D. R. James, Ph.D. Dissertation, Florida State University, 1977 (unpublished).
[15] G. Young, K. Van Bibber, E. Cosman, A. Sperduto, W. Thoms, and L. Herreid, Bull. Am. Phys. Soc. 21, 679 (1976).
[16] T. M. Cormier, J. Applegate, B. M. Berkowtiz, P. Braun-Munzinger, P. M. Cormier, J. W. Harris, C. M. Jachcinski, L. L. Lee, Jr., J. Barrette, and H. E. Wegner, Phys. Rev. Lett. 38, 940 (1977).
[17] H. V. Klapdor, G. Rosner, H. Reiss, and M. Schrader, Nucl. Phys. A244, 157 (1975).
[18] E. Almqvist, D. A. Bromley, and J. A. Kuehner, Phys. Rev. Lett. 4, 515 (1960).
[19] G. J. KeKelis, A. H. Lumpkin, D. W. Kemper, and J. D. Fox, Phys. Rev. C 15, 664 (1977).
[20] D. Shapira, R. G. Stokstad, M. W. Sachs, A. Gobbi, and D. A. Bromley, Phys. Rev. C 12, 1907 (1975), and references therein.

Effect of an Electric Field upon Resonances in the H^- Ion

P. A. M. Gram, J. C. Pratt, and M. A. Yates-Williams
Los Alamos Scientific Laboratory, Los Alamos, New Mexico 87545

and

H. C. Bryant, J. Donahue, H. Sharifian, and H. Tootoonchi
The University of New Mexico, Albuquerque, New Mexico 87131

(Received 18 October 1977)

Two resonances in the photodetachment cross section of H^- have been studied in a variable dc electric field by exploiting a relativistic colliding-beam technique. The Feshbach resonance, at a photon energy of 10.930 eV, splits into two and, possibly, a smaller third component for field strengths of about 10^5 V/cm, disappearing entirely for fields greater than 2×10^5 V/cm. The shape resonance, at a photon energy of 10.98 eV, begins to diminish in amplitude and broaden only for fields above 5×10^5 V/cm.

We have observed the effects of large barycentric electric fields on the photodetachment cross section of H^- ions in an intense laser beam in the region of two resonances near 11 eV. Our unusual experimental method[1] makes use of the relativisitic kinematics of the 800-MeV H^- beam

0031-9007/78/4002-0107/$0.00

at the Clinton P. Anderson Meson Physics Facility (LAMPF).

Variation of the angle of intersection between the beam of H^- ions ($\beta = 0.842$) and an N_2 laser beam ($h\nu = 3.678$ eV) from 20° to 160° produces a Doppler-shifted photon energy ranging from 1.4 to 12.2 eV. The photodetached electrons, confined to a 4-mrad cone with a laboratory energy of 435 keV, are separated from the main H^- beam by a weak "collection magnet" some 25 cm downstream from the interaction region and are focused onto a solid-state counter. Their distinct energy signal in coincidence with the laser flash identifies them as photodetached.

The near-luminal velocity of the ions simplifies the task of applying a strong electric field to the interaction region. A small electromagnet, the "Stark magnet," located at the intersection of the beams, produces a magnetic field of up to 1200 G normal to the plane defined by the two beams. In the rest frame of the ions this field transforms to 2200 G accompanied by an electric field of 5.6×10^5 V/cm. With these fields, the energy of interaction between the ion and the electric field is about 100 times that of the interaction with the magnetic field; the effects observed are therefore presumed due to the electric field. The region of uniform field extends about 1 cm along the flight path of the ion, centered on the interaction region; laser-excited resonances with widths in excess of 0.06 meV would spend more than one half life in the uniform field.

A second magnet gap 4 cm downstream from the interaction region, sharing the yoke and coils of the Stark magnet, counters the effect the gap centered on the interaction region has upon the trajectories of the photodetached electrons, so that they still reach the detector. Further details of the apparatus are the same as described in Ref. 1.

The resonances we have studied occur near 11-eV photon energy.[2,3] Figure 1 shows the data (points with error bars), taken with no applied field in the interaction region, the theory of Broad and Reinhardt[4] (solid line), and their theory folded with a 10-meV Gaussian resolution function. The integrated cross section of the Feshbach resonance is taken to be (8 meV)a_0^2, where a_0 is the Bohr radius. Although the zero of the experimental energy scale is uncertain (±20 meV) because of uncertainty in the H^- energy and the absolute angle, relative energies are quite accurate. Our scale has been shifted to agree with the theoretical Feshbach energy.[4]

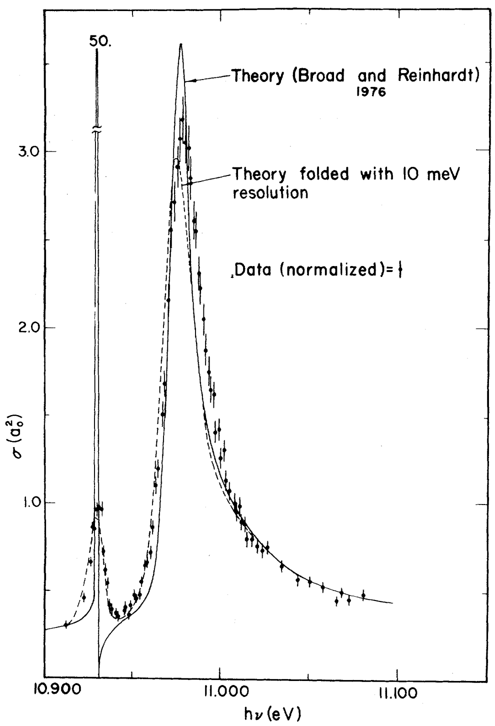

FIG. 1. Comparison of "zero-field" data with the calculated cross sections of Broad and Reinhardt (Ref. 4). Error bars reflect counting statistics only. The set of experimental points has been normalized to be fitted by the theory in the continuum. The zero point on the experimental energy scale has also been adjusted (slightly) within our experimental uncertainty to match the theory.

Figures 2(a) and 2(b) show what occurs with the application of a small field. As the field is raised from its residual value, 1.2×10^4 V/cm, to 2.2 $\times 10^5$ V/cm, the Feshbach resonance line disappears entirely while the shape resonance is not visibly affected. In Fig. 3 the behavior of the Feshbach resonance is followed as the electric field is raised in small steps to 1.3×10^5 V/cm. The "zero-field" data, Fig. 3(a), where there was a small residual field, represent the best resolution that we have attained so far, 5.5 meV, full width at half-maximum (FWHM). At 9.2 $\times 10^4$ V/cm, Fig. 3(d), the narrow line is visibly splitting into two components which become clearly resolved with the next two increments of field.

These data are analyzed by testing the hypothesis that the cross section $\sigma(E)$ is fitted by three Gaussian peaks of the same standard deviation, δ, with variable central energy, E_B, and amplitudes A, B, and C which lie on a background described by at most a second-order polynomial in energy; namely,

$$\sigma(E) = A\exp[-(E - E_B + \Delta)^2/2\delta^2] + B\exp[-(E - E_B)^2/2\delta^2] + C\exp[-(E - E_B - \Delta)^2/2\delta^2] + D + G[E - E_S] + H[E - E_S]^2,$$

where Δ is the energy gap between adjacent peaks, E_S is fixed at 10.975 eV, and E is the barycentric photon energy.

Table I presents the results of our fitting program. The outcome suggests that the small third component is significant. However, the most striking result is that the separation of the major peaks, Δ, depends linearly upon the electric field. A quadratic dependence is normally expected, unless the atomic configurations being mixed by the field are degenerate in energy. The energy shift in the linear electric-field effect is given by $E = Kea_0F$, where a_0 is the Bohr radius, F is the electric field strength, e is the electronic charge, and K is a dimensionless con-

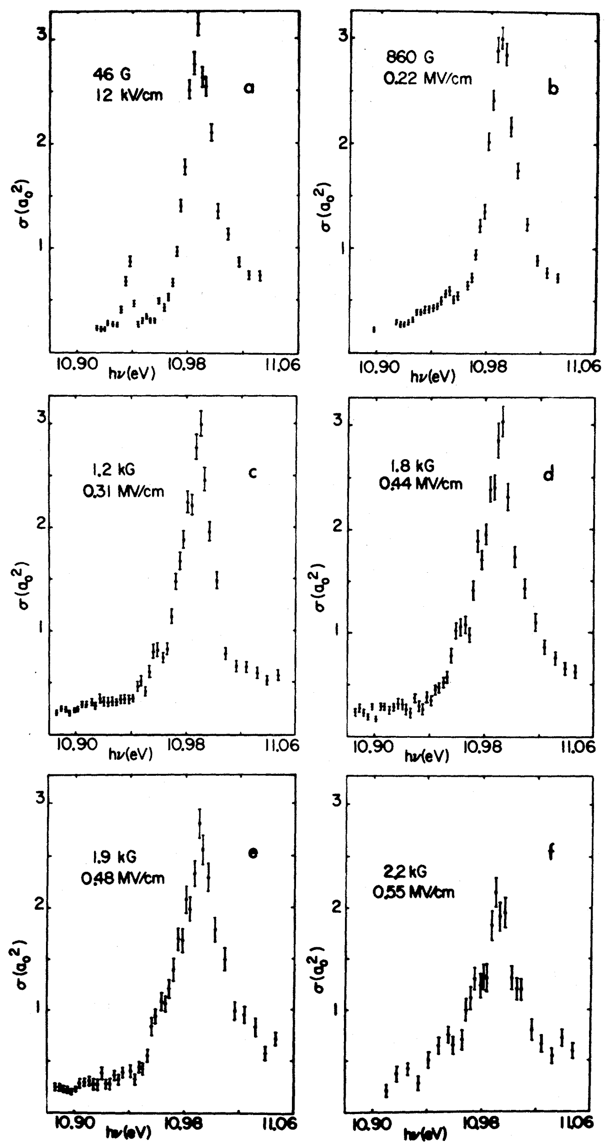

FIG. 2. Behavior of the resonance region for large steps in the applied dc field. Values are given in each case for the barycentric magnetic and electric fields.

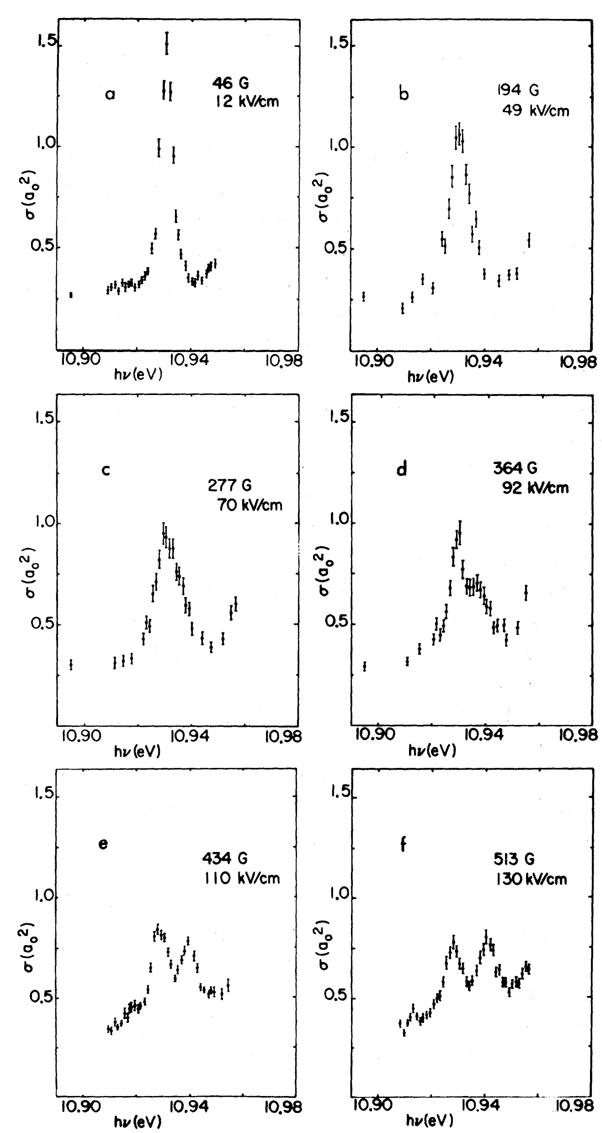

FIG. 3. Behavior of the Feshbach resonance for small steps in the applied field. Note that the energy scale is twice as large as in Fig. 2.

TABLE I. Results of fitting three Gaussian line shapes and a background to the data shown in Fig. 3. The Gaussians are assumed to have the same standard deviation, δ (in meV), which is essentially the energy resolution; their amplitudes are A, B, and C (in units of a_0^2). The field F is given in units of 10^5 V/cm. The energy at which peak B occurs is given by E_B (eV, with errors of ~ 0.4 meV), with A and C being spaced an energy Δ (in meV) below and above B, respectively. The background is assumed linear with terms D (in a_0^2) and G (in a_0^2/eV). In the case of 3(d) a small quadratic term in the energy has also been assumed for which $H = (14 \pm 31)a_0^2/\text{eV}^2$. In case 3(f), the possibly significant peak 15 meV below B is narrower than our resolution, and a fit, with present constraints, results in an insignificant (and negative) amplitude. J give the χ^2 per degree of freedom. The errors in the parameters correspond to changes which would cause χ^2 to increase by 1. We estimate that the systematic errors in the field F may be as high as 10%.

Case	F	A	B	C	E_B	Δ	δ	G	D	J
3(a)	0.12	0	1.11 ± 0.02	0	10.9305	...	2.33 ± 0.04	2.4 ± 0.3	0.47 ± 0.02	2.86
3(b)	0.49	0.16 ± 0.03	0.75 ± 0.03	0.24 ± 0.03	10.9304	5.5 ± 0.4	2.4 ± 0.1	2.1 ± 0.7	0.42 ± 0.03	2.65
3(c)	0.70	0.11 ± 0.03	0.55 ± 0.02	0.28 ± 0.03	10.9299	6.1 ± 0.2	2.9 ± 0.1	2.9 ± 0.7	0.50 ± 0.04	0.87
3(d)	0.92	0.10 ± 0.03	0.54 ± 0.03	0.28 ± 0.03	10.9296	8.1 ± 0.3	2.6 ± 0.2	10 ± 4	0.75 ± 0.09	1.28
3(e)	1.10	0.05 ± 0.02	0.43 ± 0.02	0.30 ± 0.02	10.9288	10.4 ± 0.2	3.0 ± 0.1	4.3 ± 0.4	0.64 ± 0.02	0.52
3(f)	1.30	−0.01 ± 0.02	0.30 ± 0.02	0.27 ± 0.02	10.9280	12.4 ± 0.3	3.1 ± 0.1	4.7 ± 0.1	0.69 ± 0.01	0.68

stant that depends upon the wave functions and is a measure of their spatial extent. For reference, $K=3$ for a well-known simple case, the degenerate $n=2$ states of hydrogen.[5] The slope of the line observed in our data representing energy difference between the components (Δ, Table I, using only the four highest-field cases) versus field yields a value of $K=20\pm 1$, confirming the expectation that the wave function of the resonance is very large.

In our measurements the laser was unpolarized; the geometry is such that we expect that the $m=0$ and the $m=\pm 1$ states are equally excited. Thus, although at least one other state of opposite parity is being mixed with the ^1P to form the observed Feshbach multiplet, we cannot distinguish between the presence of a ^1S and of a ^1D from our data. Indeed the Stark multiplet seen in our data can be interpreted either as a triplet, which requires the admixture of two other states of opposite parity, or, ignoring the weak peak A, a doublet, which requires only one other state. Data taken with a linearly polarized laser beam could clarify the picture, since the orientation could be chosen to excite the $m=\pm 1$ states only or to excite mainly the $m=0$ state.

Although the splitting between the components of the multiplet is fitted well be a linear function of the Stark field F, the energy of the peak B is fitted well by a quadratic function of F:

$$E_B(\text{meV}) = 10930.41 + 0.95F - 2.16F^2,$$

where F is in units of 10^5 V/cm. One half of the coefficient of F^2 can be interpreted as the difference between the polarizability of the Feshbach resonance and the H^- ground state. The large value of this coefficient also supports the notion that the radial wave function of the Feshbach resonance is spread over a large region.

Figures 2(c)–2(f) display the behavior of the shape resonance as the field is increased to large values. Finally, at 5.5×10^5 V/cm the shape resonance has clearly diminished in area and broadened from 21 to 28 meV. The growth of the errors indicated in the figure is partly due to an experimental problem encountered with the use of strong magnetic fields at the interaction point. The "beam" of forward-going photodetached electrons is apparently defocused by fringe fields to the extent that the collection magnet cannot bring all of them to the detector. This effect is corrected by a normalization based on the relative counting rates observed at energies far from any resonance.

Detailed calculations of the electric-field broadening of both of the resonances have been undertaken for the first time (known to us) by Wendoloski and Reinhardt.[6] They predict, in their preliminary exploration of the effect, that the Feshbach resonance should disappear under the influence of a small field, just as is observed, and that the shape resonance should spread from 14 meV at zero field to 26 meV at 4×10^5 V/cm. These numbers are in reasonable agreement with observation. Further work, both experimental and theoretical, is needed to clarify the behavior of this interesting two-electron system.

We thank R. Carlini, T. Rupp, and W. Thomas for their help in running the experiment and the personnel of LAMPF, especially R. Werbeck,

J. Gomez, J. van Dyke, G. Krause, and C. Dalton, for their contributions to its success. We acknowledge with thanks helpful discussions with C. D. Lin, W. P. Reinhardt, and J. J. Wendoloski. This work was supported by the U. S. Energy Research and Development Administration.

[1]H. C. Bryant *et al.*, Phys. Rev. Lett. 38, 228 (1977).

[2]C. D. Lin, Phys. Rev. Lett. 35, 1150 (1975).

[3]J. Macek and P. G. Burke, Proc. Phys. Soc. London 92, 351 (1967); J. Macek, Proc. Phys. Soc. London 92, 365 (1967).

[4]J. T. Broad and W. P. Reinhardt, Phys. Rev. A 14, 2159 (1976).

[5]Nils Ryde, *Atoms and Molecules in Electric Fields* (Almquist and Wiksell, Stockholm, Sweden, 1976).

[6]J. J. Wendoloski and W. P. Reinhardt, private communication (unpublished), and Phys. Rev. A (to be published).

New Photon-Correlation Effects in Near-Resonant Multiphoton Ionization

S. N. Dixit and P. Lambropoulos
Physics Department, University of Southern California, Los Angeles, California 90007
(Received 23 March 1977)

It is shown that in near-resonant N-photon ionization, chaotic light can be more efficient than purely coherent light by more than the usual factor of $N!$ (N factorial). This occurs when the resonant intermediate state interferes with another nearby state. Specific results are presented for two-photon ionization.

Multiphoton processes depend on the photon correlation (coherence) properties of the radiation and not simply on the average intensity, as single-photon processes do.[1-4] In short,[5,6] the yield of nonresonant N-photon processes with chaotic (incoherent) radiation is larger by a factor of $N!$ than with purely coherent (Glauber state) radiation. More generally, the process is proportional to the Nth-order correlation function of the radiation, a result intimately related to the nonresonant nature of the process which leads to the factorization of the field-correlation function from the atomic parameters. As shown in two recent theoretical papers,[7,8] in the presence of a resonant intermediate atomic state, this dependence on photon correlations changes considerably. For example, in two-photon ionization via one intermediate state, all dependence on photon correlations disappears when the transition from the initial to the (resonant) intermediate atomic state is completely saturated, which occurs in the limit of large field intensities. Well below saturation, one has the usual factor of 2! between chaotic and purely coherent radiation. Similar results are obtained for resonant processes of higher order. Under saturation conditions the difference between the effect of chaotic and coherent light diminishes. Thus the factor $N!$ has so far been considered as the maximum ratio between chaotic and coherent radiation for multiphoton transitions into a continuum.

In this Letter, we report a surprising new result. We show that under near-resonance conditions, the ratio of the yields for chaotic to that for purely coherent radiation can be significantly larger than $N!$ (N factorial). This occurs when a near-resonant intermediate state interferes with another nearby state or even with a background due to more than one distant nonresonant level.

The problem can be formulated in more than one way. We have chosen the resolvent-operator[9] formalism which has been used in a number of recent papers on resonance processes. Consider an initial atomic state $|g\rangle$ and a single-mode photon state $|n\rangle$ of frequency ω. The initial state of the system "atom plus field" is $|I\rangle = |g\rangle|n\rangle$. Consider in addition the system states $|A = |a\rangle|n-1\rangle$ and $|B\rangle = |b\rangle|n-1\rangle$, where the atomic states $|a\rangle$ and $|b\rangle$ are both connected to the state $|g\rangle$ by a dipole single-photon transition. Let $|F\rangle = |f\rangle|n-2\rangle$ be the final state for two-photon ionization, where the atomic state $|f\rangle$, assumed here to be in the continuum, is connected to both $|a\rangle$ and $|b\rangle$ by a single-photon electric dipole transition. The energies of the system states are denoted by ω_I, ω_A, ω_B, and ω_F and are measured in inverse seconds, as all Hamiltonians have been divided by $\hbar$. In terms of atomic and field energies we have $\omega_I = \omega_g + n\omega$, $\omega_A = \omega_a + (n-1)\omega$, $\omega_B = \omega_b + (n-1)\omega$, and $\omega_F = \omega_f + (n-2)\omega$. We shall be interested in photon frequencies near and around the resonance frequency $\omega_{ag} \equiv \omega_a - \omega_g$. It is rather straightforward to write a set of equations governing the relevant matrix elements of the resolvent opera-