

Effect of Resonances on the Coherent Control of the Photoionization and Photodissociation of HI and DI

Langchi Zhu,¹ Kunihiro Suto,² Jeanette Allen Fiss,¹ Ryuichi Wada,² Tamar Seideman,^{3,*} and Robert J. Gordon¹

¹Department of Chemistry (m/c 111), University of Illinois at Chicago, 845 W Taylor Street, Chicago, Illinois 60607-7061

²Department of Molecular Engineering, Kyoto University, Kyoto 60601, Japan

³Institute for Theoretical Atomic and Molecular Physics, Harvard-Smithsonian Center for Astrophysics, 60 Garden Street, Cambridge, Massachusetts 02138

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The phase lag between the ionization and dissociation products of HI and DI molecules obtained by one- and three-photon excitation was measured as a function of excitation energy in the vicinity of a molecular resonance. The phase lag was observed to have an asymmetric profile with a deep minimum near the center of the resonance. A strong isotope effect was also observed. A theoretical analysis, using projection operators to partition the Lippmann-Schwinger equation, showed that a molecular phase is responsible for the energy dependence of the phase lag. [S0031-9007(97)04599-7]

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Controlling the branching ratios of atomic and molecular processes by quantum mechanical interference has received much recent attention [1,2]. Control has been achieved in various experiments by using two excitation paths to populate a continuum state and varying either the energy detunings or the relative phase of the two paths. An example of the first method is a three-level Λ excitation scheme to induce structure in the continuum of Na_2 [3]. One laser was used to transfer population to the continuum, while a second, more intense laser was used to drive a transition from an unpopulated level to the continuum. The control parameter in this case is the energy difference between the final states reached by the two paths. A second example is the use of a four-level diamond configuration to reach a common upper level of Ba via two different resonant states [4]. The control parameter in this case is the detuning from the intermediate states. In both experiments, control is achieved by using energy tuning to alter the profile of a continuum state, which in turn affects the dissociation or ionization product ratio [5]. The phases of the two paths play no role in these experiments.

A different control strategy is to fix the energies and vary the relative phase of the two paths. For example, the two paths may consist of excitation from a discrete initial state to a degenerate continuum by either three photons of frequency ω_1 (wavelength λ_1) or by one photon of frequency $\omega_3 = 3\omega_1$ [6]. The excited continuum state correlates asymptotically to two product channels, A and B. The probability of forming product channel S by a single photon is $p_3^S = I_3 F_3^S$, whereas the probability of obtaining the same product channel by three photons is $p_1^S = I_1^3 F_1^S$, where I_3 and I_1 are the intensities of the two laser beams, and F_3^S and F_1^S are frequency-dependent constants. When both fields are present, the total probability of forming each product channel is [6]

$$p^S = p_1^S + p_3^S + 2(I_3 I_1^3)^{1/2} |F_{13}^S| \cos(\phi + \delta_{13}^S), \quad (1)$$

where F_{13}^S is a frequency dependent constant, and $\phi =$

$\phi_3 - 3\phi_1$ is the phase difference of the two fields. We define the phase lag of the two channels by the difference $\Delta\delta = \delta_{13}^A - \delta_{13}^B$. The yield of product channel S can be maximized (or minimized) by setting $\phi = -\delta_{13}^S$. The physical origin of δ_{13}^S will be discussed later.

Recently Zhu *et al.* [7] reported using this scheme to control the branching ratio of ionization and dissociation of HI molecules at energies slightly above the first ionization potential (~ 10.6 eV). The principle finding of that study was that at wavelengths λ_1 near 355 nm $\Delta\delta$ was $\sim 150^\circ$, whereas for λ_1 near 354 nm $\Delta\delta$ was ~ 0 . Although the source of the phase lag was unknown, the observation that the two wavelength regions correspond to transitions to different molecular Rydberg states suggests that resonances might play a role in the energy dependence of $\Delta\delta$. We report here data and theory which clarify the control mechanism.

The experimental method has been described previously [8]. Briefly, a pulsed (10 Hz) molecular beam of HI or DI was injected between repeller and extractor electrodes in a vacuum chamber. The molecular beam was crossed with a XeCl excimer-pumped UV dye laser (4 mJ/pulse). The third harmonic of the dye laser was generated by focusing the laser beam into a cell containing 2–10 Torr of Xe. The phase between the fundamental (UV) and the vacuum ultraviolet (VUV) third harmonic was varied by passing the two beams through a chamber containing several Torr of H_2 gas, which has very different refractive indices at the two wavelengths [9]. The two beams were focused into the reaction chamber by a pair of Al/MgF₂-coated spherical mirrors ($f = 20.3$ cm). Ions produced in the reaction region were accelerated into a 1 m flight tube and detected by a microchannel plate and a boxcar signal averager. HI^+ (or DI^+) and I^+ mass peaks were monitored as the H_2 pressure was slowly increased, so that ϕ varied by at least five cycles of 2π . From previous work [7,10] we know that the I^+ peak is produced by the ionization of neutral I atoms by one or more ω_1

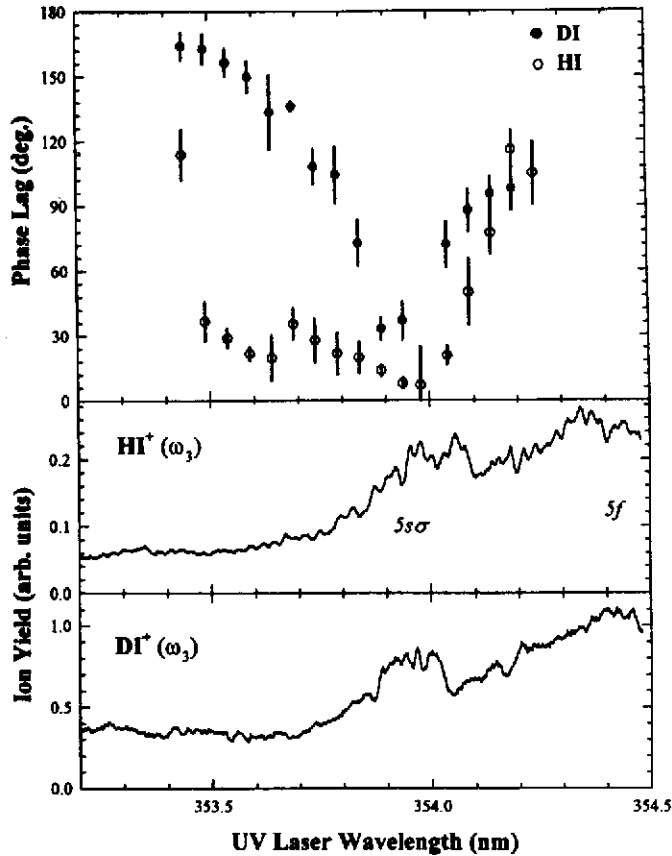


FIG. 3. Phase lag between the dissociation and ionization channels for HI (open symbols) and DI (closed symbols). The VUV ionization spectra of HI and DI are shown in the bottom panels.

brief outline is presented below, where, for simplicity of presentation, we treat the field-matter interaction within the golden rule approximation. We consider a bound eigenstate $|g\rangle$ that is coupled to several continua of general nature, both directly and via a manifold of (possible coupled) resonances (see Fig. 4).

Following partitioning of the eigenstates $|ES\hat{k}^- \rangle$ of the field free Hamiltonian H_M we obtain for the matrix element of $D^{(j)}$

$$\begin{aligned} \langle g|D^{(j)}|ES\hat{k}^- \rangle &= \langle g|D^{(j)}|ES\hat{k}_1^- \rangle \\ &+ \langle g|D^{(j)}[I + (E^- - PH_M P)^{-1}PH_M] \\ &\times QGQH_M|ES\hat{k}_1^- \rangle. \end{aligned} \quad (2)$$

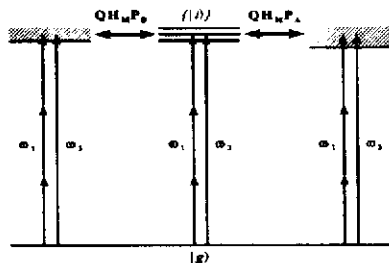


FIG. 4. Schematic drawing of the model.

where E is the total energy, and \hat{k} denotes the photofragment or photoelectron scattering angles. In Eq. (2) the $-$ superscript signifies incoming wave boundary conditions, serving as a reminder that $|ES\hat{k}^- \rangle$ evolves in the limit where the unbound coordinate q^S [the internuclear separation vector in the case of (pre)dissociation and the electron-core separation vector in the case of (auto)ionization] approaches infinity to an eigenstate of an asymptotic Hamiltonian $\lim_{q^S \rightarrow \infty} H_M$, with well-defined channel index S and scattering angles \hat{k} . We denote by P a joint projector onto the continua, $P = \sum_S P_S$ with $P_S H_M P_{S'} = 0$ for $S' \neq S$ and $Q|ES\hat{k}^- \rangle$ is the portion of the eigenstate that vanishes asymptotically. Thus, $P_S H_M P_S$ is the scattering Hamiltonian of the uncoupled S continuum, $|ES\hat{k}_1^- \rangle$ are the eigenstates of $P_S H_M P_S$, $P_S H_M Q$ is the nonadiabatic interaction coupling the excited and continuum manifolds, and $QGQ = [E^- - QH_Q]^{-1}$ is the Green operator corresponding to an effective Hamiltonian $QH_Q = QH_M Q + QH_M P(E^- - PH_M P)^{-1}PH_M Q$ (see Fig. 4). The extension of this formalism to the case of a nonperturbative field and details of the theory will be published separately [18].

Constructive and destructive interference between the direct (first) and resonance-mediated (second) terms of Eq. (2) produce a Fano-type line shape [19] for both the one- and three-photon processes [p_1^S and p_3^S in Eq. (1)]. The total transition probability for channel S is given by

$$p^S = |\langle g|D^{(1)}|ES\hat{k}^- \rangle + e^{i\phi} \langle g|D^{(3)}|ES\hat{k}^- \rangle|^2, \quad (3)$$

which has the same form as Eq. (1). The origin of the phase lag may be seen explicitly by identifying the phase of each of the terms in Eq. (2). The direct amplitude can be written as $f_j^{S,d} \exp(i\delta_j^{S,d}) = \langle g|D^{(j)}|ES\hat{k}_1^- \rangle$ where $f_j^{S,d}$ and $\delta_j^{S,d}$ are real. Assuming first that a single resonance dominates, $Q = |i\rangle\langle i|$, the resonance-mediated amplitude can be written as the product of two factors, $f_j^{S,r} \exp(i\delta_j^{S,r}) = \langle g|D^{(j)}[(E^- - PH_M P)^{-1}PH_M + I]|i\rangle\langle i|H_M|ES\hat{k}_1^- \rangle$ and $f \exp(i\delta) = [E - E_i - \Delta_i - i\Gamma_i/2]^{-1}$, where E_i is the eigenvalue of $QH_M Q$ corresponding to eigenvector $|i\rangle$, and Γ_i and Δ_i are the standard width and shift parameters [17] (modified to account for the presence of two continua [18]). The quantity δ is the familiar Breit-Wigner phase [20], and $\delta_j^{S,d}$ and $\delta_j^{S,r}$ may be described as "molecular phases," which arise from the fact that $|ES\hat{k}_1^- \rangle$ is complex. Another possible source of the molecular phase is a continuum intermediate state in the three-photon process [16]. With the above definitions the transition probability for channel S takes the form

$$\begin{aligned} p^S &= |f_1^{S,d} e^{i\delta_1^{S,d}} + f_1^{S,r} f e^{i(\delta_1^{S,r} + \delta)} \\ &+ e^{i\phi} (f_3^{S,d} e^{i\delta_3^{S,d}} + f_3^{S,r} f e^{i(\delta_3^{S,r} + \delta)})|^2. \end{aligned} \quad (4)$$

Several limiting forms of this equation are apparent, giving rise to different schemes for controlling the product branching ratio. One scheme may be achieved by

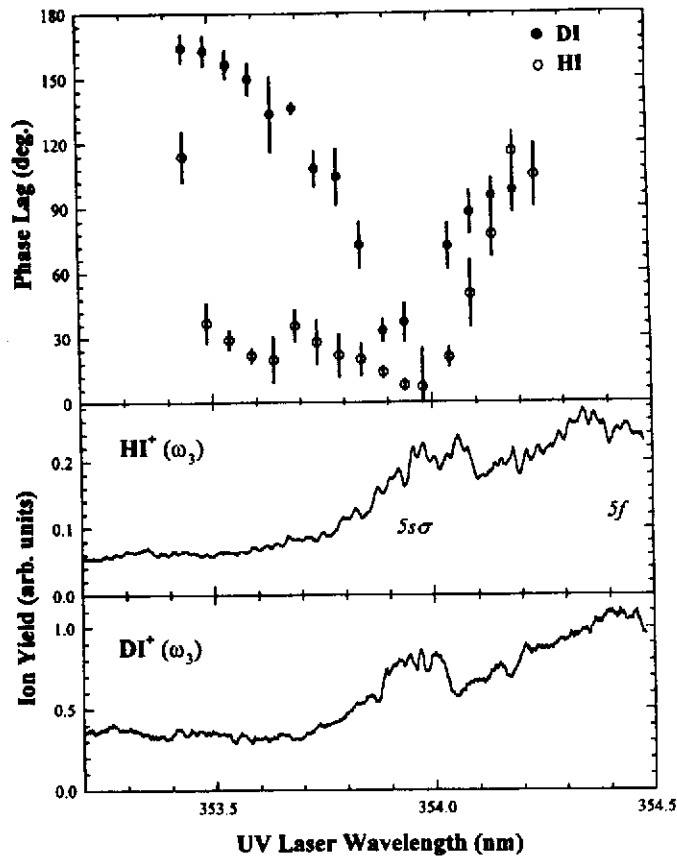


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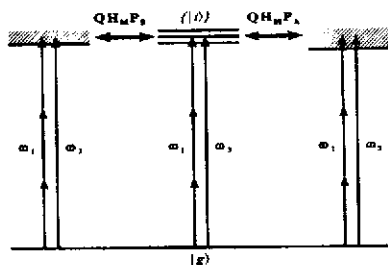


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Several limiting forms of this equation are apparent, giving rise to different schemes for controlling the product branching ratio. One scheme may be achieved by

scanning a single laser across the resonance (e.g., by setting $f_1^{s,d} = f_1^{s,r} = 0$ and varying δ) [21]. Control is possible because the relative flux into each channel depends on the interference between the direct and indirect paths, which varies with excitation energy. Another, more active form of control results from interference between the one- and three-photon paths. Control is achieved at a fixed energy by selecting ϕ to maximize the flux in either channel. The phase lag can arise from two sources. The first shown by Nakajima *et al.* [15,22] is a consequence of the fact that the shape of the Fano profile in a bichromatic field depends on the phase difference ϕ of the two fields. It follows from their work that the cross term in Eq. (1) can have a nonvanishing value of δ_{13}^S even if the molecular phases are all zero. In this case, however, one can show [18] that $\delta_{13}^S \rightarrow 0$ far from resonance, whereas the experimental phase lag has a minimum near resonance and rises strongly far from the center of the resonance.

The second source of $\Delta\delta$ is the molecular phase. From Eqs. (2) and (4) we find that, in the limit of a single resonance and no direct component, the molecular phase vanishes. With the first term of Eq. (2) omitted, p_1^S and p_3^S take Lorentzian line shapes, and the cross term

$$F_{13}^S = \int dk \times \frac{\langle g|D^{(1)}|i\rangle \langle i|H_M|ES\hat{k}_1^- \rangle \langle ES\hat{k}_1^- |H_M|i\rangle \langle i|D^{(3)}|g\rangle}{(E - E_i - \Delta_i)^2 + (\Gamma_i/2)^2} \quad (5)$$

is real. In the presence of a direct component F_{13}^S includes three additional terms, of which the product of the two direct components of the one- and three-photon elements dominates far from resonance, while the product of the resonant components, Eq. (5), dominates on resonance. Provided the direct component carries an at most weakly energy dependent molecular phase, δ_{13}^S has a minimum on resonance and rises asymmetrically as the frequency is red or blue detuned, saturating at the value corresponding to the direct process alone far from the resonance, as indeed was observed experimentally (Fig. 3). More generally, if several resonances contribute, F_{13}^S is no longer real and the energy dependence of the phase lag is more structured. The qualitative behavior of the phase is, nevertheless, similar, having minima at the resonance energies and rising to the asymptotic plateau. A nonnegligible direct component and/or overlapping resonances would produce a net phase also on resonance [18].

In conclusion, we have shown that the phase lag between dissociation and ionization is caused by a molecular phase, and that the energy dependence of the phase lag is caused by a resonance coupled to two continua. In the center of the resonance the molecular phase for the one- and three-photon paths nearly cancels, producing a minimum in the phase lag. In this region the phase in-

formation contained in the preparation step is lost in the eventual decay of the molecule. Off resonance, the phase lag can be used to determine the relative phase of two outgoing wave functions, a quantity not readily obtained by other methods. Moreover, our experimental data and theoretical formalism point to the possibility of determining resonance lifetimes by measurement of the phase lag, even in the presence of a direct component.

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*Permanent address: Steacie Institute, National Research Council, 100 Sussex Drive, Ottawa K1A0R6, Canada.

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