Photodissociation dynamics of the $A^2\Sigma^+$ state of SH and SD radicals


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(Received 30 October 2008; accepted 5 December 2008; published online 20 January 2009)

Atomic sulfur products from predissociation of the lowest rotational states of SH/SD $A^2\Sigma^+$ ($v' = 0, 1, 2$) are studied using velocity map imaging. The dissociation process, which is slow compared to rotation, is dominated by interference effects due to predissociation of states with low rotation quantum numbers prepared by photoexcitation using overlapping transitions of different parities. The measured product angular distributions can be modeled using the methods presented recently by Kim et al. [J. Chem. Phys. 125, 133316 (2006)]. The $S(3P_J) (2+1)$ resonance enhanced multiphoton ionization scheme used in the detection step of the experiment is sensitive to the angular momentum polarization of the atomic fragments. $S(3P_J), J = 2, 1, 0$, fine-structure yields, angular distributions, and atom polarization parameters are reported. Strong polarization of the $S(3P_J)$ products was observed along with a weak sensitivity of the branching ratio to excess energy and a full insensitivity of the atomic product polarization to excess energy. None of the data fit the predictions of either adiabatic or diabatic photodissociation, emphasizing the need for a fully quantum treatment. © 2009 American Institute of Physics. [DOI: 10.1063/1.3056570]

I. INTRODUCTION

The past decade has witnessed tremendous progress in the understanding of the photodissociation of closed-shell diatomic molecules. For the hydrogen halides (HX, X = F, Cl, Br, I), for example, fully quantum calculations combined with advanced imaging experiments have unraveled the intricate dynamics of a molecule evolving from its optically excited molecular orbital configuration to its final product atoms. Using this combined approach of experiment and theory, it is now possible to determine the amplitudes and phases of the fragment matter waves that describe a complete photochemical experiment.

For a few molecules, fully quantum mechanical calculations are available to guide the interpretation of photodissociation dynamics. For the rest, photodissociation is described by limiting case models based on adiabatic, sudden, or statistical approximations. Detailed investigation of diatomic molecule photodissociation is thus crucial in evaluating the validity of these limiting case models, which are at present the best tools available for interpreting the photodissociation dynamics of complex polyatomic molecules.

Open-shell molecules such as the covalent hydroxyl and mercapto radicals, OH and SH, tend to absorb at longer wavelengths and more readily produce highly reactive fragments such as O($^1D$) and S($^1D$) than most closed-shell molecules. Owing to their relatively simple electronic structure, a fully quantum treatment of the photodissociation properties of OH (Refs. 7–10) and the isovalent SH (Ref. 11) molecules should soon be tractable. Experimentally, however, the production and isolation of sufficient concentrations of these species for further study are challenging. Although OH has been the subject of numerous theoretical and experimental investigations, only a few collision-free studies of OH photodissociation have been reported. Velocity map imaging detection of O and H atom products was previously used by Radenovic et al. to study direct one-photon photodissociation of vibrationally excited OH and OD $X^2\Pi(v')$ via the repulsive $1^2\Sigma^-$ electronic state, while Zhou et al. used H-Rydberg atom photofragment translational spectroscopy (HR-PTS) to study predissociation of OH excited to the $A^2\Sigma^+$ ($v' = 3$ and 4) levels.

Reports of the photodissociation dynamics of SH and SD also remain sparse, although the spectroscopy of the SH and SD $A^2\Sigma^+ - X^2\Pi$ band systems is well explored and understood. In early studies, broadening of spectroscopic lines was observed as the SH($A^2\Sigma^+$) vibrational quantum number increased. The line broadening is a consequence of predissociation, which reduces the lifetime of the excited state. Laser induced fluorescence and cavity ring-down spectroscopy were subsequently used to measure the lifetimes of SH($A^2\Sigma^+, v' = 0–2$) and SD($A^2\Sigma^+, v' = 0–2$) levels, and their dependence on rotational quantum number ($N'$). The reductions in lifetimes as $v'$ increases and the $N'$-level dependence were quantitatively modeled by a series of calculations using ab initio and experimentally fitted potential energy (PE) curves and Fermi golden rule (FGR) calculations. The $A^2\Sigma^+$ state correlates diabatically to $H(2S) + S(1D)$ atomic fragments but is crossed by three repulsive PE curves, $1^2\Sigma^-, 1^2\Sigma^+$, and $1^2\Pi$, that correlate to the ground dissociation limit $H(2S) + S(3P)$ as illustrated in Fig. 1. Coupling from the $A^2\Sigma^+$ state to one or more of these repulsive PE curves causes predissociation. The FGR calcula-
In short, a mixture of 2.
s26 potentials including the 1 22 - state. Provided estimates of the strengths of the couplings between levels of the 1 22 - state; only for higher vibrational levels do couplings to the 1 22 - and 1 42 states become competitive or dominant.

These various spectroscopic and computational studies identified the states participating in the predissociation of low vibrational levels of the A 2Σ+ state of SH and SD and provided estimates of the strengths of the couplings between the bound and the dissociative excited states. They did not, however, provide any information on the dynamics that occur following coupling into one or more of the repulsive electronic states: for example, the associated bond breaking may occur adiabatically on a single PE curve or via nonadiabatic transitions between these states (as is known to arise for HF, HCl, and HBr). Zhang et al.25 used HR-PTS to study direct photodissociation via the 1 2Σ+ repulsive state, and Dillon and Yarkony undertook computational studies of the dynamics to address these uncertainties. Direct dissociation via the repulsive wall of the bound A 2Σ+ state was observed by Janssen et al.27 when starting from ν"=2—6 vibrational levels of the ground X 2Σ state and exciting above the H(2S)+S(1D) dissociation limit. There have also been analyses of the dynamics of secondary photodissociation of SH radicals produced from H2S at UV wavelengths ranging from 193 to 244 nm,28—31 at these wavelengths, excitation of the SH is to one or more repulsive potentials including the 1 2Σ+ state.

In our previous study of direct SH and SD photodissociation and in the current work, which focuses on predissociation dynamics of the A 2Σ+ state at excitation wavelengths from 292 to 324 nm, we employed a pulsed electric discharge supersonic expansion as a source of the mercapto radicals. To derive detailed information about photofragment angle and speed distributions, we used the velocity mapping technique, coupled with polarized laser photodissociation and resonance enhanced multiphoton ionization (REMPI) detection of S(3P_j) photofragments. The predissociation of mercapto radicals that have been selectively prepared in specified rovibrational levels of the A 2Σ+ excited electronic state occurs on a time scale comparable to the rotational period. Under such circumstances, the spatial anisotropy of the internuclear axis in the transient excited state varies with the rotational transition due to interference effects,32,33 and therefore so does the space-fixed recoil anisotropy. By measuring values of the anisotropy parameter βr across several rovibrational transitions, detailed insights are obtained of the dissociation in the molecular frame. Nuclear spin angular momentum can also couple to the prepared rotational and electronic angular momenta of the excited state through hyperfine interactions on the time scale of the experiment, potentially affecting the measured angular distributions. Furthermore, the energy splitting of the e, f parity states of the SH free radical is sufficiently small that there is simultaneous excitation of different parity states, leading to a second level of interference effects, which can be simulated, together with the effects of molecular rotation, using the treatment by Houston and co-workers. From velocity map images, resolved at the S(3P_j) spin-orbit level, and their dependence on the polarization of the probe laser, we derive anisotropy parameters, S(3P_j) spin-orbit branching ratios, and m-state propensities. The outcomes of these measurements enable us to make detailed deductions about the dynamics on coupled, dissociative PE curves that cross the A 2Σ+ state and correlate to H(2S)+S(1D) photofragments.

II. EXPERIMENTAL

A detailed description of our velocity map imaging apparatus has been given in previous publications,13,14,27,35 and a schematic is presented in Fig. 2. In short, a mixture of 25%—30% H2S (D2S) in Xe was expanded into a vacuum chamber through a pulsed (10 Hz) Jordan valve (0.5 mm orifice diameter) and a pulsed ring-type stainless-steel discharge electrode (4 mm diameter).36 The pulsed-discharge electrode was mounted 2 mm downstream from the valve orifice, and SH (SD) radicals were produced in the discharge during the supersonic expansion. The voltage and the width and delay of the pulse applied to the discharge device were optimized for the expansion conditions (~2 bar backing.
The results were obtained by pumping the second harmonic of a neodymium doped yttrium aluminum garnet (Nd:YAG) laser (Quanta Ray DCR-3). $S(3P, J)$ atoms from SH dissociation via $\text{SH}(A, v=0)$ (top), $\text{SH}(A, v=1)$ (middle), and $\text{SH}(A, v=2)$ (bottom). Images shown in the columns from left to right are for $S(3P, J)$ with $J=0$, 1, and 2.

![FIG. 3. (Color online) Raw velocity map images of $S(3P, J)$ from SH photodissociation via $\text{SH}(A, v=0)$ (top), $\text{SH}(A, v=1)$ (middle), and $\text{SH}(A, v=2)$ (bottom). Images shown in the columns from left to right are for $S(3P, J)$ with $J=0$, 1, and 2.](image)

III. RESULTS AND DISCUSSION

Representative velocity map images of $S(3P, J)$ atoms from SH dissociation are shown in Fig. 3. The results derived from analysis of the intensities and angular dependence of such images are presented in this section as branching ratios for the $J=0$, 1, and 2 spin-orbit levels and as $m$-state populations and recoil anisotropy parameters. The data provide clear evidence for nonadiabatic dissociation pathways following coupling from the $A \rightarrow X$ state onto one or more repulsive PE curves and are compared with the expectations of adiabatic and diabatic (sudden) models for the dissociation dynamics.

A. Branching into $S(3P, J)$ spin-orbit levels

Integration of the intensities of images recorded with the dissociation laser tuned to the $P_1(1.5)$ transition of a selected...
SH or SD A 2Σ+→X 2Π vibrational band and the detection laser tuned to the frequencies of the (2+1) REMPI transitions for the S(3PJ) levels with J=0, 1, and 2 allows the branching between spin-orbit states of this photofragment to be deduced. The image intensity ratios depend not only on the relative populations of the three spin-orbit levels, however, but are also influenced by the transition strengths for the two-photon excitations used for the state-specific product detection. Table I shows measurements of the branching into S(3PJ) levels following 232 and 226 nm excitations of SH and comparison with values derived by Zhou et al. using the HR-PTS technique, from which the REMPI transition strengths are deduced to be approximately equal. The multi-photon ionization detection scheme used in the current work is resonance enhanced at the two-photon level by the S(3p 34p P J) levels and has been employed previously by Hsu et al. and Brouard et al. The former study demonstrated that the line strengths for the individual transitions are very similar, in agreement with our measurements. The integrated image intensities can depend on the direction of polarization of the linearly polarized probe laser, but this effect, which depends on the electronic alignment of the S atoms, is minimized for the VH geometry. Thus VH image intensities were used to derive spin-orbit branching ratios without further correction.

To avoid bias in the branching ratio measurements, efforts were made to keep the experimental conditions constant for each series of measurements. Particular care was taken to ensure that the pressures of H2S were kept constant (both above the dissociation limit) in Fig. 5 can be used to account for the greater propensity for S(3P2) production. Some of the details concerning the construction of this diagram are addressed in Sec. III B 2 and the Appendix. The model proposed by Wheeler et al. attributes the predissociation of the v=0→2 vibrational levels of the SH or SD A 2Σ+ state to a spin-orbit mediated interaction almost exclusively with the Ω=½ component of the 1 2Σ+ state; only for higher vibrational levels of SH(A) and SD(A) does coupling to the 1 2Σ− and 1 4Π states become competitive with or dominant over the coupling to the 1 2Σ− state. The adiabatic correlation of the 1 4Π state is to H(2Σ+)S(3P2) fragments, and our results thus suggest that the majority of dissociative flux follows this adiabatic pathway or crosses to the 2Π state of the ground electronic state, which also dissociates to this limit (see Sec. III B 2). Branching to the other S(3Pj) spin-orbit states might arise from couplings from the A 2Σ+ state directly to the 1 2Σ− state or the Ω=½ component.

Table II shows the S(3Pj) spin-orbit branching ratios derived from the experimental data for predissociation of SH(A 2Σ+, v=0, 1, and 2) and SD(A 2Σ+, v=0, 1, and 2), and the values are plotted as a function of total kinetic energy release (TKER) of the photofragments (i.e., excess energy above the dissociation limit) in Fig. 4. The ratios are obtained from averages of three to five sets of images and the uncertainties account for the reproducibility of the measurements and include systematic errors that derive from the assumption of equal line strengths for the REMPI transitions used. From the data in the table and figure, it is evident that the dominant products are S(3P0) atoms, with minor branching into the S(3P1) and S(3P2) spin-orbit levels. For SH(A, v=0→2), the branching to S(3P0) is generally observed to be greater than to S(3P1) atoms, but for SD(A, v=0→2), any preference for S(3P0) over S(3P1) is not so clear-cut. There is an apparent increase in the branching to S(3P3) at the expense of S(3P2) for SH(A, v=2) predissociation.

The correlation diagram for SH (and SD) shown in Fig. 5 can be used to account for the greater propensity for S(3P2) production. Some of the details concerning the construction of this diagram are addressed in Sec. III B 2 and the Appendix. The model proposed by Wheeler et al. attributes the predissociation of the v=0→2 vibrational levels of the SH or SD A 2Σ+ state to a spin-orbit mediated interaction almost exclusively with the Ω=½ component of the 1 2Σ+ state; only for higher vibrational levels of SH(A) and SD(A) does coupling to the 1 2Σ− and 1 4Π states become competitive with or dominant over the coupling to the 1 2Σ− state. The adiabatic correlation of the 1 4Π state is to H(2Σ+)S(3P2) fragments, and our results thus suggest that the majority of dissociative flux follows this adiabatic pathway or crosses to the 2Π state of the ground electronic state, which also dissociates to this limit (see Sec. III B 2). Branching to the other S(3Pj) spin-orbit states might arise from couplings from the A 2Σ+ state directly to the 1 2Σ− state or the Ω=½ component.

### Table I

<table>
<thead>
<tr>
<th>Wavelength (nm)</th>
<th>S(3P0)</th>
<th>S(3P1)</th>
<th>S(3P2)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Current work</td>
<td>0.44(0.02)</td>
<td>0.35(0.03)</td>
<td>0.21(0.02)</td>
</tr>
<tr>
<td>Zhou et al.</td>
<td>0.50(0.03)</td>
<td>0.36(0.01)</td>
<td>0.14(0.02)</td>
</tr>
<tr>
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<td>0.30(0.03)</td>
<td>0.18(0.02)</td>
</tr>
<tr>
<td>Zhou et al.</td>
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<td>0.36(0.05)</td>
<td>0.14(0.07)</td>
</tr>
</tbody>
</table>

### Table II

<table>
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<th>S(3P1)</th>
<th>S(3P2)</th>
</tr>
</thead>
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<td>0.36(0.05)</td>
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</tr>
</tbody>
</table>

FIG. 4. Spin-orbit branching ratios for S(3Pj) plotted against TKER. Experimental fractions of J=2 (circles), J=1 (triangles), and J=0 (squares) for SH (filled symbols) and SD (open symbols) are compared to the predictions at the adiabatic (solid line) and the diabatic (dashed lines) limits.

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components of the $1^4\Pi$ state, which correlate adiabatically with the other spin-orbit components. The evidence from previous FGR calculations of predissociation rates suggests, however, that this is not the case, with the possible exception of the SH(A, $v=2$) radicals. It is more likely that the branching to the $S(^3P)$ and $S(^5P)$ photoproducts occurs at larger internuclear separations, following bond extension on the repulsive $1^4\Sigma^-$ state. The diabatic $4\Sigma^-$ and $4\Pi$ states are spin-orbit coupled, but $4\Sigma^- - 4\Sigma^-$ interaction is forbidden via first-order spin-orbit coupling. There are thus likely to be regions along the $S-\Pi$ bond extension coordinate where nonadiabatic (derivative) couplings exist between the adiabatic $\Omega = \frac{1}{2}$ components of the $4\Sigma^-$ and $4\Pi$ states that mix dissociative flux onto the higher lying of these two PE curves. Further mixing to the adiabatic $2^3\Sigma^-$ state is then possible. These deductions are broadly consistent with the arguments presented by Zhou et al. to account for the branching between $S(^3P)$ fine-structure states following direct photoexcitation to the repulsive $1^2\Sigma^-$ state. The extent to which $S(^3P)$ population is favored is, however, greater for predissociation of the low-lying vibrational levels of the $A^2\Sigma^+$ state than for the direct excitation of the $1^2\Sigma^-$ state, most likely reflecting the different asymptotic adiabatic correlations, as shown in Fig. 5.

The increase in the branching to $S(^3P)$ products for SH(A, $v=2$) predissociation is consistent with the onset of coupling from the $A^2\Sigma^+$ state to the $1^4\Pi$ state either directly or via initial $A^2\Sigma^+-1^2\Sigma^-$ coupling. As Fig. 5 shows, one $\Omega = \frac{1}{2}$ component of the $1^4\Pi$ state correlates adiabatically to the $S(^3P)$ state. Wheeler et al. estimated direct $A^2\Sigma^+-1^2\Pi$ coupling to be negligible, but coupling to the $1^2\Sigma^-$ state amounts to $\leq 10\%$ of the main channel to the $1^4\Sigma^-$ state at this vibrational level. More recent calculations by Brites et al. suggested that direct coupling to the $1^2\Sigma^-$ state contributes about $25\%$ of the total predissociation rate for SH(A, $v=2$) [but it remains a negligible pathway for SH(A, $v=0$) and SD(A, $v=0$, 1, and 2), in accord with the deductions of Wheeler et al.]. As Fig. 1 shows, the crossing of the $1^4\Pi$ state with the outer wall of the $A^2\Sigma^+$ state lies too high in energy to influence the predissociation dynamics of these low vibrational levels.

The observation of S atoms populating spin-orbit levels other than expected from adiabatic correlation arguments is clear evidence for nonadiabatic dynamics between the three sets of repulsive potentials correlating to the $H(^2S)+S(^3P)$ limit. Comparisons can be made with the expectations of a “sudden” or diabatic model in which the electronic wave functions of the molecular states are projected onto a basis of atomic states for this dissociation limit. The results of such calculations, the procedures for which are explained in detail elsewhere, are shown in Table III and included in Fig. 4. While the fully diabatic model for the $1^2\Sigma^+_{1/2}$ state dissociation agrees qualitatively with the experimental observations, for example, in favoring the $J=2$ spin-orbit level, with near equal propensities for the $J=1$ and 0 levels, the experiments show a greater branching into $S(^3P)$ products. The dynamics following coupling from the $A^2\Sigma^+$ state to the $1^4\Sigma^-$ state can thus not be viewed as occurring either in the purely adiabatic or diabatic limits, and an intermediate picture must be invoked.

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**TABLE III.** $S(^3P)$ spin-orbit branching ratios $P(J)$ for predissociation of SH(A) via the $1^4\Sigma^-$ and $1^2\Sigma^-$ states calculated in the sudden or diabatic limit and the associated population distribution $P_m$ over available $m$ states, where $P_m$ is the total population of the $+m$ and $-m$ states when $m$ is nonzero.

<table>
<thead>
<tr>
<th></th>
<th>$1^4\Sigma^+_{1/2}$</th>
<th></th>
<th>$1^2\Sigma^+_{1/2}$</th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>$J$</td>
<td>$P(J)$</td>
<td>$m=0$</td>
<td>$m=\pm 1$</td>
<td>$m=\pm 2$</td>
</tr>
<tr>
<td>0</td>
<td>0.222</td>
<td>1.000</td>
<td>...</td>
<td>...</td>
</tr>
<tr>
<td>1</td>
<td>0.167</td>
<td>0.000</td>
<td>1.000</td>
<td>...</td>
</tr>
<tr>
<td>2</td>
<td>0.611</td>
<td>0.727</td>
<td>0.273</td>
<td>0.000</td>
</tr>
</tbody>
</table>
B. Recoil velocity anisotropy and angular momentum polarization

The velocity map images of the photofragments show angular dependence to their intensities because of a combination of the effects of spatial anisotropy in both the recoil velocities and electronic angular momenta of the \( S(3P_J) \) atoms. The experiments are sensitive to the angular momentum polarization because the \( 2+1 \) REMPI detection scheme employs linearly polarized light. The lifetimes of the vibrational levels of the \( A^{2S+} \) state studied here range from a few picoseconds to many nanoseconds and are thus comparable to or much longer than the rotational period of the molecules. Nevertheless, the predissociation can result in an anisotropic distribution of photofragment recoil velocities (\( p_{\nu} \)).

Within the axial recoil approximation, the recoil velocity vector for dissociation of a diatomic molecule lies parallel to the internuclear bond axis, and a preferred distribution of the planes of rotation of the ensemble of excited state molecules in space, prepared by absorption of polarized light, can thus maintain an overall anisotropy of the recoil velocity vectors of the atomic photofragments.

The state-specific detection of \( S(3P_J) \) atoms with \( J=0, 1, \) and \( 2 \) means that each velocity map image contains information on only a single fragmentation channel. Analysis of the variation of intensity with angle for the outermost few pixels of the 2D velocity map images is thus equivalent to analysis of a slice through the three dimensional (3D) reconstruction of the Newton sphere of recoil velocities of the nearly monoenergetic photofragments and might be expected to show an intensity variation of the form described by:

\[
p_{\nu}(\theta) = \frac{1}{4\pi}[1 + \beta_2 P_2(\cos \theta)].
\]

Here, \( p_{\nu}(\theta) \) is the product angular distribution with respect to \( \theta \), the angle between the linear polarization of the photolysis laser (\( e \)) and the direction of product recoil (\( v \)) in the laboratory frame. \( P_2(\cos \theta) = \frac{1}{2}(3\cos^2 \theta - 1) \) denotes a second Legendre polynomial, and the anisotropy parameter \( \beta_2 \) takes limiting values of \( +2 \) for a parallel transition and \( -1 \) for a perpendicular transition followed by prompt dissociation. In practice, this intensity distribution is further modulated by the sensitivity of the probe laser polarization to the angular momentum alignment of the photofragments [the detection efficiency \( I_{\text{det}}(\theta) \)], resulting in a general expression for the intensity variation with image angle:

\[
I(\theta) \propto p_{\nu}(\theta) I_{\text{det}}(\theta) = \frac{1}{4\pi}[1 + \beta_2 P_2(\cos \theta) + \beta_4 P_4(\cos \theta) + \beta_6 P_6(\cos \theta)].
\]

The values of the parameters \( \beta_2, \beta_4, \) and \( \beta_6 \) depend on the recoil anisotropy, the angular momentum alignment of the products, the polarizations of the photolysis and probe lasers, and the plane of the imaging detector. In the following analysis of image angular variation to extract quantitative information on photofragment velocity anisotropy and angular momentum alignment, we follow the procedures employed by Coroiu et al., which were based on the work of Mo and Suzuki. The resultant velocity anisotropy parameters are compared with expectations for a predissociative system using the method of Kim et al. Angular momentum alignment is presented in terms of populations of magnetic sublevels of the \( S(3P_J) \) atoms with \( J=1 \) or \( 2 \), which are indicative of the degree of nonadiabatic dissociation dynamics. The general analysis procedures are first reviewed and then the \( \beta_e \) parameter values and \( S(3P_J) \) spin-orbit resolved \( m \)-state populations are discussed separately.

The outermost rings (two to three pixels wide) of the velocity map images were fitted to the angular function given in Eq. (2). In all cases, incorporation of a \( P_4(\cos \theta) \) function made no discernible difference to the quality of the fits, and the expansion in Eq. (2) was thus truncated at fourth-order terms, giving two fit parameters, \( \beta_2 \) and \( \beta_4 \), for each image. The detection efficiency of the photofragments with aligned electronic angular momentum can be expressed as:

\[
I_{\text{det}}(\theta) = \sum_k \rho_k^{(k)} P_k(J_f, J_i),
\]

where \( P_k(J_f, J_i) \) are line-strength factors for the two-photon transition between levels with angular momentum quantum numbers \( J_i \) (initial state) and \( J_f \) (final state). The \( \rho_k^{(k)} \) are the components with \( q=0 \) of rank-\( k \) multipole moments of the density matrix describing the populations and coherences of the magnetic sublevels with angular momentum quantum number \( J_i \) and projection quantum numbers \( m \). The projections are defined with respect to the probe laser linear polarization vector, and the subscript \( FF \) thus denotes the probe frame. The index \( k \) takes values of \( 0, 2, \) and \( 4 \) for a probe transition induced by linearly polarized light and resonance enhanced at the two-photon level. When \( J_i \neq J_f \), Mo and Suzuki demonstrated that the ratios of the line-strength factors with \( k=2 \) and \( 4 \) to the \( k=0 \) factor simplify to the convenient form (expressed here for linearly polarized light):

\[
\tilde{P}_k(J_f, J_i) = \frac{P_k(J_f, J_i)}{P_0(J_f, J_i)} = 5 \sqrt{2k+1} \sqrt{2J_i+1} (-1)^{l_f+l_i} \sqrt{2} \frac{2}{0} \frac{2}{0} \frac{k}{0} \times \begin{pmatrix} J_i & J_i & k \\ 2 & 2 & J_f \end{pmatrix}.
\]

Evaluation of the \( 3-j \) and \( 6-j \) symbols in this expression for the REMPI transitions employed gives \( \tilde{P}_2(1,2) = \sqrt{5}/14 \) and \( \tilde{P}_4(1,2) = -2\sqrt{7}/21 \) for the \( S(3p^34p) \) two-photon excitation and \( \tilde{P}_2(2,1) = 1/\sqrt{2} \) and \( \tilde{P}_4(2,1) = 0 \) for the \( S(3p^34p) \) transition. For transitions from the \( ^3P_0 \) level, the \( k=2 \) and \( 4 \) line-strength factors are, by definition, zero.

Two geometries of laser polarizations were employed for the measurements; in both cases the photolysis laser was polarized vertically in the laboratory (denoted \( V \), and parallel to the detector face), but the probe laser was polarized either vertically (specified hereafter as \( VV \) geometry) or horizontally (\( VH \) geometry). If we take a laboratory frame of reference defined by Cartesian axes with \( Z \) parallel to the vertical polarization of the photolysis laser, \( X \) the axis of propagation of the two laser beams, and \( Y \) the axis along the TOF tube.
toward the detector (see Fig. 2), the probe laser polarization lies either along the Z or Y axes (V or H, respectively) and the detector lies parallel to the XZ plane. Rotation from the frame defined by the photofragment recoil velocity vector [the recoil frame (RF)] to the probe laser frame transforms the multipole moments of the angular momentum in accord with

\[
\rho_{0,RF}^{(k)} = \rho_{0,RF}^{(k)} + \sum_{q'=-k}^{k} D_{q'q}^{(k)}(\Phi, \Theta, \chi) f_q^k(q', \Phi, \Theta, \chi),
\]

where \(D_{q'q}^{(k)}(\Phi, \Theta, \chi)\) is a Wigner rotation matrix, and the Euler angles \((\Phi, \Theta, \chi)\) link the two frames of reference. In the case of the VV polarization geometry, we identify \((\Theta, \Phi) = (\theta, \phi)\), the polar and azimuthal angles of the recoil velocity vector in the PF (which is equivalent to the frame defined by the photodissociation laser polarization), and \(\theta\) also corresponds to the polar angle in the image plane. For the VH geometry, we make the identification that \(\Phi = \pi/2\), \(\Theta = \pi/2\), and \(\chi = \theta\) (considering that the outermost ring in the image arises from recoil velocities in the XZ plane). The analysis is further simplified if we make the assumptions that the angular momentum polarization is cylindrically symmetric about the recoil velocity vector and independent of direction. A consequence is that only \(q' = 0\) components of \(\rho_{0,RF}^{(k)}\) are nonzero, and we neglect the effects of coherences between \(m\) states but can derive information on fractional populations of these \(m\) states \((f_m)\) in the RF from

\[
\rho_{0,RF}^{(k)} = \sum_{m=-J_i}^{J_i} (-1)^{J_i-m} \sqrt{2J_i+1} \binom{J_i}{J \; m \; J_i} f_m.
\]

With the above approximations, Eq. (5) reduces to

\[
\rho_{0,RF}^{(k)} = \rho_{0,RF}^{(k)} P_k(\cos \theta)
\]

for the VV polarization geometry and

\[
\rho_{0,RF}^{(k)} = \rho_{0,RF}^{(k)} P_k(\cos \pi/2)
\]

for the VH geometry. Substitution of Eqs. (4) and (7) or Eq. (8) into Eq. (3) and re-expression of the result as a single expansion in Legendre polynomials allows us to evaluate Eq. (2) as

\[
I(\theta) \approx [1 + \beta_v P_v(\cos \theta)] \sum_k \rho_{0,RF}^{(k)} \tilde{P}_k(J_f, J_i) P_k(g(\theta))
\]

\[
= \sum_{m=0}^{J_f} \beta_m P_m(\cos \theta),
\]

where \(g(\theta)\) denotes \(\cos \theta\) for the VV geometry and \(\cos(\pi/2)\) for the VH geometry. The expansion coefficients used in Eq. (2) are obtained as

\[
\begin{pmatrix}
\beta_0 \\
\beta_1 \\
\beta_2 \\
\beta_3 \\
\beta_4 \\
\beta_5 \\
\beta_6
\end{pmatrix} =
\begin{pmatrix}
1 & \frac{1}{2} \beta_v & 0 & 0 & 0 & 0 & 0 \\
\beta_v & 1 + \frac{2}{3} \beta_v & \frac{2}{3} \beta_v & 0 & 0 & 0 & 0 \\
0 & \frac{2}{3} \beta_v & 1 + \frac{2}{3} \beta_v & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \frac{2}{11} \beta_v & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & \frac{2}{11} \beta_v & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & \frac{2}{11} \beta_v & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & \frac{2}{11} \beta_v
\end{pmatrix}
\]

for the VV geometry and

\[
\begin{pmatrix}
\beta_0 \\
\beta_1 \\
\beta_2 \\
\beta_3 \\
\beta_4 \\
\beta_5 \\
\beta_6
\end{pmatrix} =
\begin{pmatrix}
1 & -\frac{1}{2} \beta_v & \frac{1}{8} \beta_v & 0 & 0 & 0 & 0 \\
\beta_v & -\frac{1}{2} \beta_v & \frac{3}{8} \beta_v & 0 & 0 & 0 & 0 \\
0 & \frac{3}{8} \beta_v & \frac{3}{8} \beta_v & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & \frac{2}{11} \beta_v & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & \frac{2}{11} \beta_v & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & \frac{2}{11} \beta_v & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & \frac{2}{11} \beta_v
\end{pmatrix}
\]

for the VH geometry.

The image analysis thus involves fitting of the angular dependence to Eq. (2) to obtain experimental values of the coefficients \(\beta_i\) and \(\beta_v\) with normalization to \(\beta_0 = 1\) and neglect of the \(\beta_v\) term. These \(\beta_i\) values can be simultaneously forward simulated for both VV and VH geometries by calculating \(\rho_{0,RF}^{(k)}\) values from Eq. (6) by cycling over possible combinations of populations of the \(m\) states and combining the values with \(\beta_i\) values in the range from \(-1\) to \(+2\). For an aligned sample, \(f_m = f_{-m}\), and we thus imposed this restriction on the simulations, together with the requirement that the populations of all \(m\) states sum to unity. It proved useful to define \(p_m = f_m\) for \(m = 0\) and \(p_m = f_{-m} + f_m\) for \(m \neq 0\). The forward simulation procedure was carried out using a simple computer program with a step size of 0.01 for both \(p_m\) and \(\beta_i\) to find optimum values for these parameters to reproduce the experimental \(\beta_i\) and \(\beta_v\) data for both VV and VH geometries, with use of the line-strength factors appropriate for each \(S(3P_i)\) state studied. Thus, recoil velocity anisotropy parameters and angular momentum alignment information were deconvoluted from the experimental angular distributions. Such an analysis is not, however, necessary for \(S(3P_0)\) atoms because the \(J = 0\) level cannot show any angular momentum alignment effects, and its detection probability is independent of probe laser polarization.

The analysis method requires that \(J_f \neq J_i\) in the two-photon absorption step in the REMPI scheme. In the case of \(S(3P_1)\) detection via the \(3p^2^4p^2(3^P_1)\leftarrow \leftarrow 3p^2(3^P)\) excitation scheme, however, the transitions of the \(J_f = 0, 1, 2\) are only partially resolved within the linewidth of the laser. For \(S(3P_2)\) detection, the \(3p^2\leftarrow 3p^2(3P_2)\) component is stronger than for excitation to the \(3P_1\) and \(3P_0\) levels, and several images were thus accumulated using this \(J_f = 0\) transition. To test the robustness of our methods, measurements of angular distributions and analysis to derive \(\beta_i\) and \(m\)-state populations were therefore made at probe laser wavelengths on the sides of the broad REMPI feature corresponding to the \(3P_1\leftarrow 3P_2\) and \(3P_2\leftarrow 3P_2\) transitions. At the chosen wavelengths, we estimate that there are \(-15\%\) and \(<5\%\) contributions, respectively, from the overlapping transitions. The values of \(\beta_i\) obtained showed no dependence on the choice of REMPI wavelength within this two-photon excitation feature, and \(m\)-state populations required only a small correction (by factors of 1.3, 0.9, and 1.0 for \(m = 0, \pm 1, \pm 2\)) to make the values derived from data.
obtained in the $^3P_2\rightarrow^3P_2$ wing of the feature match those from the $^3P_1\rightarrow^3P_1$ wing. A similar investigation of the $S^2(^3P_1)$ REMPI detection scheme, for which the $^3P_2$ and $^3P_1$ upper state components are also incompletely resolved, showed that within experimental error, no correction was required for either $\beta_0$ or $m$-state populations derived from images obtained at excitation wavelengths across the spectral feature.

In most cases of fitting to the angular dependence of the outer parts of velocity map images, the fits are good, and Fig. 6 compares the experimental data with simulations obtained using Eq. (2) and the values of $\beta$ and $m$-state populations from the data analysis procedure described in the text. The gray lines indicate the differences between experimental data and the simulations.

![ FIG. 6. Experimental angular distributions (open circles) from which anisotropy parameters and alignment information are obtained: (a) $S^2(^3P_2)$ images and (b) $S^1(^3P_1)$ images for the SD($v=1$), $Q(1,5)$ line. The upper graphs correspond to VV geometry images and the lower graphs to VH geometry. The solid lines are simulations of the angular dependence obtained using Eq. (2) and the values of $\beta$ and $m$-state populations from the data analysis procedure described in the text. The gray lines indicate the differences between experimental data and the simulations.](image)

The treatment of the anisotropy parameters resulting from predissociation of states with lifetimes comparable to or longer than their rotational periods has been the subject of recent studies by Kim et al. 32 and Kuznetsov and Vasyutinskii. 33 Here, we compare the outcomes of our experimental measurements of $\beta$; parameters for predissociation of different vibrational levels of SH(A) and SD(A) with the predictions of the model proposed by Kim et al. 32 using the computer program BETAOFNU. 34 This model accounts for the polarization of the internuclear axis caused by absorption of linearly polarized light, the effects of the rotational motion of the excited state, and interference effects in the spatial distribution of photofragments caused by overlapping $P$, $Q$, and $R$ branch absorption features. The model does not explicitly deal with the state-specific detection of atomic fragments in individual spin-orbit states nor does it take into account possible depolarization of the spatial distribution of the internuclear axis caused by coupling of the rotational and electronic angular momenta of the SH or SD to the nuclear spin of the H or D atom (hyperfine depolarization effects). The consequences of hyperfine depolarization on $\beta_e$ values are discussed further below.

Figure 7 shows experimentally determined $\beta_e$ values for SH($v=2$) obtained by detection of $S^2(^3P_2)$ and $S^1(^3P_0)$ in the VH geometry. Equation (11) demonstrates that no alignment effects are expected in this geometry and thus $\beta_2=\beta_0$; experimental fits to Eq. (2) largely confirm this, as was noted at the end of the previous section. The short lifetime of SH(A,$v=2$) means that spectroscopic lines are sufficiently broadened that a number of $\beta_e$ values could be obtained experimentally for photoexcitation across the wavenumber range from 34 190 to 34 260 cm$^{-1}$ to investigate the frequency dependence of $\beta_e$. The data in Fig. 7 show an oscil-
The agreement between experimental and calculated values was poorer for these data than for the SH(A,v=2) results, with the measured \( \beta_v \) values consistently smaller in magnitude than the expectations of the theoretical model [with the exception of the \( P_{1}(1.5) \) line for which the calculated \( \beta_v=0.0 \) was as expected for an upper state with \( J=\frac{3}{2} \) and \( N=0 \)]. For SH(A,v=0) and SD(A,v=0), the measured anisotropy parameters all lie close to zero. The lifetimes of the SH(A,v=0) and SD(A,v=0) levels are, respectively, 3.2 and 250 ns for low rotational levels \(^{18,20,21}\) and these lifetimes are not only considerably longer than the rotational periods but are also comparable to or longer than the characteristic time scales for coupling of the nuclear spins of the H or D atoms to the rotational and electronic spin angular momenta of the SH or SD radicals. The nuclear hyperfine splittings for \( SH(A,v=0) \) were reported by Ubachs et al.,\(^ {18} \) and, for \( J \) values up to \( S/2 \), range from 289.9 to 898.6 MHz. Taking a representative value of \( \nu_{HF}=500 \text{ MHz} \), the time scale for hyperfine coupling of \( J \) with \( I_H \) to give the resultant total angular momentum \( F \) is \( \tau=2\pi/\nu_{HF}=12 \text{ ns} \) and about an order of magnitude longer for SD.\(^ {46} \) The nuclear spin of \( ^{32}\text{S} \) is \( I_S=0 \), and it thus does not play a part in the following discussion.

We consider the hyperfine coupling interactions further because they can cause depolarization of the optically prepared spatial alignment of rotational angular momentum (de-
noted here by a quantum number \( N \) or internuclear axes of the SH or SD. Depolarization of \( N \) by \( S \) (the electronic spin angular momentum) is treated implicitly in the computer simulations described so far, but hyperfine depolarization is not incorporated. Values of \( \beta_\nu \) can thus degrade from the expectations from calculations using the method of Kim et al.\(^{32}\) because of precession of the nuclear framework rotational angular momentum \( N \) about \( J \) and of \( J \) about the total angular momentum \( F \).\(^{39}\) The procedures for calculation of the time-dependent depolarization of \( N \) are well established:\(^{50-52}\) a rotational angular momentum alignment parameter \( A_q^{(k)}(0) \) at \( t=0 \) is multiplied by a time-dependent hyperfine depolarization factor \( G^{(k)}(t) \) of the same rank to give a time-dependent alignment parameter:

\[
A_q^{(k)}(t) = A_q^{(k)}(0)G^{(k)}(t). \tag{12}
\]

To account for hyperfine depolarization of \( \beta_\nu \), Zhang et al.\(^{53}\) simply multiplied \( \beta \) values by \( G^{(2)}(t) \) factors in their study of photodissociation of aligned HF molecules. To verify this approach, we computed the alignment parameters \( A_q^{(2)}(N) \) for one-photon excitation of a diatomic molecule via \( P_i^1Q \) and \( R \) branch transitions using linearly polarized light, and, using Eq. (12), applied the hyperfine depolarization appropriate for SH calculated from\(^{51}\)

\[
G^{(k)}(t) = \sum_{F,F'} \cos \left[ (E_{P_i} - E_{P_F})t/\hbar \right] \\
\times \frac{(2F' + 1)(2F + 1)}{(2J + 1)} \left[ \begin{array}{ccc} F' & F & k \\ J & J & 1 \end{array} \right]^2 \tag{13}
\]

and knowledge of splittings of hyperfine energy levels, \( E_{P_i} - E_{P_F} \). Figure 9 shows the time dependence of the \( G^{(2)}(t) \) parameters for excitation of SH via the \( Q_1(1.5) \) and \( R_1(1.5) \) transitions. The resultant depolarized angular momentum alignment parameters were used to calculate the associated internuclear axis alignment parameter in the presence of hyperfine coupling. As a test of the method, these latter parameters, derived in the absence of hyperfine interactions, were compared with values of \( A_0 \), the parameter introduced by Zare\(^{54}\) to describe the spatial distribution of internuclear axes of a diatomic molecule in a long-lived excited state following the absorption of linearly polarized light, and the two calculations gave identical results. In the limit of an excited state that predissociates on a time scale much longer than its rotational period, with the fragmentation occurring under conditions of axial recoil, the internuclear axis alignment parameter \( A_0 \) is equivalent to the photofragment recoil anisotropy parameter \( \beta_\nu \) and we thus deduce that the time-dependent hyperfine depolarization of \( \beta_\nu \) is correctly described by

\[
\beta_\nu(t) = \beta_\nu G^{(2)}(t), \tag{14}
\]

in agreement with the method of Zhang et al.\(^{53}\) Here, the \( \beta_\nu \) values are those computed using the theory of Kim et al.,\(^{32}\) and the time chosen in Eq. (14) corresponds to the predissociative lifetime of the SH or SD radicals in their \( \Sigma + \Sigma^+ \) state and a particular \( \nu \) and \( N \) level.

From these calculations, we conclude that hyperfine depolarization can reduce \( \beta_\nu \) values by a factor of 0.25–1 following excitation of SH to \( \nu=0 \) of the \( \Sigma + \Sigma^+ \) state via the \( Q_1(1.5) \) line and a factor of 0.66–1 for the \( R_1(1.5) \) line. The precise values depend on the predissociative lifetimes, but the respective time-averaged reduction factors are 0.68 and 0.83. We do not make a more quantitative comparison with the experimental data, however, because of the blended nature of the spectroscopic absorption lines. For higher vibrational levels of the \( \Sigma + \Sigma^+ \) state with lifetimes on the picosecond time scale, little or no hyperfine depolarization of \( \beta_\nu \) is expected.

2. Populations of \( m \) levels for \( S(3P_1) \) and \( S(3P_2) \) photoproducts

The angular dependences of the velocity map images for \( S(3P_1) \) and \( S(3P_2) \) photofragments show characteristic signatures of electronic angular momentum polarization in the S atoms which can be quantified in terms of unequal populations, \( P_{m\nu} \), of \( m \) sublevels. The outcomes of the data analysis procedure described above are presented in Figs. 10 and 11 and Table IV; we do not distinguish \( \pm m \) level populations because the experiments are sensitive only to even rank alignment moments. The data presented in Table IV are averages of the populations of the \( m \) sublevels for the \( S(3P_2) \) and \( S(3P_1) \) atoms following excitation of SH or SD to each of the \( \nu=0, 1, \) and 2 levels of the \( \Sigma + \Sigma^+ \) state via \( P(1.5), Q(1.5), \) and \( R(1.5) \) transitions. As is evident from the standard deviations of the populations presented in Table IV and as will be discussed further below, these \( m \) sublevel populations are insensitive to the use of SH or SD in the experiments, the vibrational level of the \( \Sigma + \Sigma^+ \) state from which predissociation occurs, or the rotational level that is excited.

There are two further clear outcomes of the measurements. The first is that the electronic angular momentum of the \( S(3P_1) \) atoms is polarized, with the \( m=\pm 1 \) levels possessing nearly four times the population of the \( m=0 \) level. The second is that the \( S(3P_2) \) atoms are also strongly polarized, with the \( m=\pm 2 \) levels exhibiting very little population but the \( m=\pm 1 \) and 0 levels both significantly and nearly
equally populated. We note that these results derive from an analysis that neglects coherence effects among the \( m \) levels and that the angular momentum polarization is not subject to hyperfine depolarization because \( \ell_S = 0 \).

Figures 10 and 11 show plots of \( m \)-state populations of the \( S(1^2P_2) \) and \( S(3^2P_1) \) atom products, respectively, for the \( P_1(1.5) \) and \( Q_1(1.5) \) and \( P_{1/2}(1.5) \) lines plotted against TKER of the atomic products. The consistency of the outcomes of the \( m \)-state populations regardless of the isotopolog, vibrational, or rotational state is clearly evident in these plots. For comparison, the predicted \( m \)-state populations for dissociation from the \( ^1S_{1/2} \) state within the adiabatic and diabatic models (see Sec. III A and Table III) are also plotted. For \( S(3^2P_2) \) atoms it is clear that neither the adiabatic or diabatic model captures the \( m \)-state populations quantitatively. For \( S(3^2P_1) \), which adiabatically should not be observed as photoproducts, the \( m \)-state populations approach the diabatic limit with the \( m = \pm 1 \) levels exhibiting the majority of the population, but there is also non-negligible population of the \( m = 0 \) level.

In the adiabatic model, a correlation diagram is constructed to determine the photodissociation product states. In the correlation diagram, the states are labeled by \( \Omega = m_S + m_H \), with \( m_S \) and \( m_H \), respectively, denoting the projections of the electronic angular momenta of the S and H atoms onto the internuclear axis. This axis corresponds to the RF if, as we expect, the axial recoil approximation applies. To derive the fragment polarization, the separate \( m_S \) and \( m_H \) quantum numbers must be assigned to the states, as was done for the \( S(1^2D_2)+H(2S_1/2) \) atomic limit in Ref. 27. For the \( S(3^2P_2)+H(2S_1/2) \) limit, we must take into account the non-zero spin-orbit coupling in the \( S(3^2P_2) \) atom. When \( \Omega \) takes the highest possible value for a given atomic fine-structure state, \( |m_S|, |J_S|, |m_H|, \) the only possibility is \( m_S = J \) and \( m_H = J_H \) (we adopt the convention that \( \Omega > 0 \)). For the other substates, there are two possibilities: \( m_S = \Omega \pm \frac{1}{2} \). These two sublevels are degenerate asymptotically and the assignment is based on the energy ordering at long range. As before, the leading long-range term that lifts the degeneracy is the dispersion interaction. The energy ordering may be derived from the anisotropy of the polarizability of the sulfur atom, as derived in the Appendix, and the result is shown in Fig. 5.

In the following discussion, we focus on states with \( \Omega = \frac{1}{2} \) because, after excitation to the \( A^2 \Sigma^+ \) state, subsequent couplings to states with \( \Omega = \frac{3}{2} \) or greater require rotation-induced interactions which will be weak for the conditions of low rotational angular momentum in our experimental study. The \( \Omega = \frac{1}{2} \) components of the \( X^2 \Pi \) and \( 1^4 \Sigma^+ \) states exhibit the long-range correlations:

\[
X^2 \Pi_{1/2} \rightarrow S(3^2P_2; m_S = 0) + H(2S_1/2; m_H = \pm \frac{1}{2}),
\]

\[
1^4 \Sigma^+_1/2 \rightarrow S(3^2P_2; m_S = \pm 1) + H(2S_1/2; m_H = \pm \frac{1}{2})).
\]

Adiabatic dissociation on the \( 1^4 \Sigma^+_1/2 \) state should thus produce exclusively \( S(3^2P_2; m_S = \pm 1) \) atoms, with mixing at long range with the \( X^2 \Pi_{1/2} \) state as degeneracy is approached resulting in population of \( S(3^2P_2; m_S = 0) \). There are no routes to formation of \( S(3^2P_2; m_S = \pm 2) \) unless an \( \Omega = \frac{5}{2} \)

![FIG. 10. m-state populations for S(1^2P_2) atoms as a function of TKER of the S and H atom photodissociation products. The experimentally determined populations of m=0 (circles), m=\pm 1 (triangles), and m=\pm 2 (squares) are compared to the predictions of the adiabatic (solid line) and diabatic (dashed lines) calculations for predissociation via the 1^1 \Sigma^+ potential (see also Table III). Open and filled symbols are data obtained, respectively, via the Q_2(1.5)+P_2(1.5) and P_2(1.5) transitions. The errors represent the quality of the fits and are \( \frac{1}{2} \) of the root mean square differences between the experimental \( \beta_1 \) and \( \beta_2 \) values obtained from fits to VV and VH images and the calculated best fit values.](image)

![FIG. 11. m-state populations for S(3^2P_1) atoms as a function of TKER. The experimentally determined populations of m=0 (circles) and m=\pm 1 (triangles) are compared to the predictions of diabatic (dashed line) calculations for predissociation via the 1^1 \Sigma^+ potential (see Table III). Open and filled symbols are data obtained, respectively, via the Q_1(1.5)+P_1(1.5) and P_1(1.5) transitions. The errors represent the quality of the fits and are \( \frac{1}{2} \) of the root mean square differences between the experimental \( \beta_1 \) and \( \beta_2 \) values obtained from fits to VV and VH images and the calculated best fit values.](image)

### Table IV

<table>
<thead>
<tr>
<th>( p_m )</th>
<th>( P(J) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( m=0 )</td>
<td>0.06</td>
</tr>
<tr>
<td>( m=\pm 1 )</td>
<td>0.79 (0.08)</td>
</tr>
<tr>
<td>( m=\pm 2 )</td>
<td>0.01 (0.06)</td>
</tr>
</tbody>
</table>
state is populated, but as noted above, this requires unfavorable Coriolis interactions. The $\Omega = \frac{1}{2}$ and $\frac{1}{2}$ components of the $1^4\Sigma^-$ state will be effectively mixed by molecular rotation because of the weak coupling of the electron spin to the internuclear axis, giving a possible mechanism for formation of the small fraction of $S(^3P_J)$ $m_S = \pm 2$ atoms suggested by our data analysis. The observed populations of the $m$ levels provide firm evidence for nonadiabatic dynamics between the $1^4\Sigma^-$ and $X^2\Pi_{1/2}$ states which were alluded to in the discussion of spin-orbit branching ratios in Sec. III A.

One $\Omega = \frac{1}{2}$ component of the $1^4\Pi_{1/2}$ state correlates to the $S(^3P_J) + H(^2S_{1/2})$ asymptote (Fig. 5) and was implicated in the production of $S(^3P_J)$ atoms following nonadiabatic transfer of flux from the $1^4\Pi_{1/2}$ state (see Sec. III A). Dissociation on the $1^4\Pi_{1/2}$ state is expected to give rise to $S(^3P_J)$ atoms with $m_S = 0$ and mixing with the $1^2\Sigma^-$ state must be invoked to account for the $S(^3P_J)$ $m_S = \pm 1$ atoms observed experimentally.

Figures 10 and 11 include predictions for $m$ sublevel populations within the approximations of adiabatic and diabatic (or sudden) limits for the dynamics. In the diabatic limit, $m_S = 0$ products should dominate the $S(^3P_J)$ channel, with population of $m_S = \pm 1$ a minor pathway. For $S(^3P_J)$ the diabatic model predicts formation of $S$ atoms exclusively with $m_S = \pm 1$. In both cases these predictions are qualitatively (but not quantitatively) reproduced by the experimental data. We thus conclude that the observed dynamical behavior is intermediate between purely diabatic and adiabatic models, which is reinforced by the observed formation of $S(^3P_J)$ and $S(^3P_J)$ photofragments, as discussed in Sec. III A.

IV. CONCLUSIONS

The predissociation dynamics of SH and SD from the $v=0$, 1, and 2 vibrational levels of the $A^2\Sigma^+$ state were investigated by velocity map ion imaging of the $S(^3P_J)$ atom products. The measured spin-orbit branching ratios into $J=0$, 1, and 2 products show that $S(^3P_J)$ atoms are the main product in all cases, with minor branching into $S(^3P_J)$ and $S(^3P_J)$ products. This branching is consistent with predissociation via the $1^4\Sigma^-$ potential which correlates adiabatically to the $S(^3P_J) + H(^2S_{1/2})$ asymptote and is in agreement with deductions from spectroscopic measurements of the SH($A_v$) and SD($A_v$) lifetimes and a model for the predissociation mechanism based on FGR calculations. The nonzero population of $S(^3P_J)$ and $S(^3P_J)$ products for predissociation of all vibrational states of SH($A$) and SD($A$) investigated is indicative of deviations from adiabatic dissociation via the $1^4\Sigma^-$ state, but branching ratios are also inconsistent with a fully diabatic model for dissociation. An intermediate picture thus prevails for the dissociation dynamics following coupling from the $A^2\Sigma^+$ to the $1^2\Sigma^-$ state in which there are nonadiabatic (spin-orbit mediated) couplings between the $1^2\Sigma^-$ and the $1^4\Pi$ potentials at large internuclear separations (and perhaps also subsequent couplings from the $1^4\Pi$ to the $1^2\Sigma^-$ potential). The evidence from measurements of $m$ sublevel populations for the $S(^3P_J)$ channel also points to nonadiabatic dynamics involving mixing of flux from the $1^4\Sigma^-$ state to the $X^2\Pi_{1/2}$ state. The near absence of $S(^3P_J)$ atoms with $m = \pm 2$ is consistent with weak or negligible Coriolis couplings during the dissociation.

An enhanced branching into $S(^3P_J)$ following excitation to SH($A_v=2$) is attributed to the onset of coupling from the $A^2\Sigma^+$ potential to the $1^2\Sigma^-$ potential and subsequent transfer of flux to the $1^2\Pi$ potential which correlates adiabatically to the $S(^3P_J) + H(^2S_{1/2})$ asymptote. This observation is in accord with the FGR calculations of predissociation rates by Wheeler et al. and Brites et al. that showed negligible coupling to the $1^2\Sigma^-$ state for SH($A_v=0, 1$) and SD($A_v=0, -2$) but an enhancement of this pathway for SH($A_v=2$), which was predicted to account for up to 25% of the predissociation mechanism of this vibrational level.

The measured anisotropy parameters $\beta_m$ show a strong frequency dependence which is well described by calculations, based on the work of Kim et al., that take into account the lifetimes of the excited state levels and interference effects resulting from excitation of the SH or SD molecules via overlapping $P$, $Q$, and $R$ branch rotational transition. Evidence is presented that the anisotropy parameters for predissociation of the longer lived SH and SD ($A^2\Sigma^+$) vibrational levels are further reduced by hyperfine interactions, with the coupling of the electronic and rotational angular momenta to the nuclear spin of the H (or D) atom on a time scale of a few nanoseconds partially degrading the prepared optical alignment of the internuclear axes in the excited state.

ACKNOWLEDGMENTS

The Nijmegen group gratefully acknowledges partial financial support by the Council for Chemical Sciences of the Netherlands Organization for Scientific Research, CW-NWO (ECHO Grant No. 700-55-025) and from the Research Training Network Molecular Universe (Grant No. MRTN-CT-2004-512302). We also thank Andre van Roij and Leander Gerritsen for expert technical assistance. The theoretical work has been financially supported by the Council for Chemical Sciences of the Netherlands Organization for Scientific Research (CW-NWO). The Bristol group thanks the University of Bristol for a postgraduate scholarship (R.A.R.) and EPSRC for funding through the LASER Portfolio Partnership Grant No. GR/S71750/01. A.J.O.-E. is grateful to the Royal Society and the Wolfson Foundation for a Research Merit Award.

APPENDIX: FRAGMENT POLARIZATION IN THE ADIABATIC MODEL

The dispersion interaction between two atoms with nonzero angular momentum is discussed in Ref. 55. As the spin-orbit coupling in the H atom is zero, however, we can ignore the H atom electron spin and write the dispersion interaction between $S(^3P_J)$ and $H(^2S_{1/2})$ as a function of the internuclear distance $R$ as

$$V_{m_S}(R) = \frac{C_6(J, m_S)}{R^6}. \quad (A1)$$

The dispersion coefficients $C_6(J, m_S)$ are given by...
\[ C_0(J, m_2) = C_{6,0}(J) - \frac{3m_2^2 - J(J + 1)}{(2J - 1)(2J + 3)} C_{6,2}(J), \]  
\[ \text{where } C_{6,0}(J) \text{ and } C_{6,2}(J) \text{ denote the scalar and rank-2 tensor components of the dispersion interactions, respectively. These two equations show that the energy ordering is determined by } m_2 C_{6,2}(J), \]

\[ C_{6,2}(J) = -\frac{3(2J + 3)}{2 \pi J} \int_0^\infty \alpha_2(J; i\omega) \tilde{\alpha}_1(i\omega) d\omega, \]

where \( \alpha_2(J; i\omega) \) is the tensor polarizability of the \( \Sigma^2P_1 \) state and \( \tilde{\alpha}_1(i\omega) \) is the dynamic polarizability of the \( \Sigma^2S_1/2 \) atom. As before,\(^{27,29}\) we assume that the sign of the integral is minus the sign of \( \alpha_2(J; i\omega) \) at \( \omega = 0 \). The polarizability of the fine-structure state is, to a very good approximation, given by\(^{58,59}\)

\[ \alpha_2(J) = (-1)^{3J + L + 1}(2J + 1) \left( \begin{array}{ccc} J & 2 & J \\ -J & 0 & J \\ L & 2 & L \end{array} \right) \times \left( \begin{array}{ccc} L & 2 & L \\ -L & 0 & L \end{array} \right)^{-1} \alpha_2(L) . \]  

The tensor polarizability \( \alpha_2(L) \) is related to the polarizability anisotropy \( \Delta \alpha(L,0) \) through\(^{27,29}\)

\[ \alpha_2(L) = -\frac{2}{3} \Delta \alpha(L,0) \frac{L(2L-1)}{L(L+1)}. \]  

Medved et al.\(^{59}\) reported a calculated value for \( \Sigma^2P \) of \( \Delta \alpha(L=1,0) = +4.49 \) a.u. Evaluating the \( 3-j \) and \( 6-j \) symbols for \( L=1 \) in Eq. (A4) gives \( \alpha_2(3J^2P_1) = \alpha_2(L=1) < 0 \) and \( \alpha_2(3J^2P_3) = -0.5 \alpha_2(L=1) > 0 \). This gives a positive anisotropic dispersion coefficient \( C_{6,2}(J=2) > 0 \) and hence \( m_2 = 0 \) is the lowest for the \( J=2 \) fine-structure state, while \( m_2 = 0 \) is the upper level for \( J=1 \).