The photodissociation dynamics of ozone at 193 nm: An $O(1D_2)$ angular momentum polarization study

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Polarized laser photolysis, coupled with resonantly enhanced multiphoton ionization detection of $O(1D_2)$ and velocity-map ion imaging, has been used to investigate the photodissociation dynamics of ozone at 193 nm. The use of multiple pump and probe laser polarization geometries and probe transitions has enabled a comprehensive characterization of the angular momentum polarization of the $O(1D_2)$ photofragments, in addition to providing high-resolution information about their speed and angular distributions. Images obtained at the probe laser wavelength of around 205 nm indicate dissociation primarily via the Hartley band, involving absorption to, and diabatic dissociation on, the $\tilde{B}^1B_2(3\,1A_1)$ potential energy surface. Rather different $O(1D_2)$ speed and electronic angular momentum spatial distributions are observed at 193 nm, suggesting that the dominant excitation at these photon energies is to a state of different symmetry from that giving rise to the Hartley band and also indicating the participation of at least one other state in the dissociation process. Evidence for a contribution from absorption into the tail of the Hartley band at 193 nm is also presented. A particularly surprising result is the observation of nonzero, albeit small values for all three rank $K=1$ orientation moments of the angular momentum distribution. The polarization results obtained at 193 and 205 nm, together with those observed previously at longer wavelengths, are interpreted using an analysis of the long range quadrupole-quadrupole interaction between the $O(1D_2)$ and $O_2(\Delta_g)$ species. © 2006 American Institute of Physics. [DOI: 10.1063/1.2210009]

I. INTRODUCTION

A. Atomic polarization studies

The electronic angular momentum polarization of the atomic fragments of molecular photodissociation has increasingly been a focus of attention over the past few years.\(^1\)\(-\)\(^4\) Its measurement involves the determination of the moments of the density matrix that characterizes fully the atomic product $M_J$ populations and any coherences between different $M_J$ states.\(^5\)\(-\)\(^6\) In many cases, the polarization of the electronic orbital angular momentum, $L$, is the quantity of particular interest, and its characterization can help identify which of the possible exit channel potentials are involved in dissociation. This information is distinct from that provided by the angular distribution of the photofragments [described by the $\beta(v)$ parameter], which provides information principally about the symmetry of the electronic state(s) initially excited in the parent molecule. Knowledge of the electronic angular momentum orientation and alignment, coupled with more familiar information about populations and angular distributions, thus provides insight about changes in orbital populations or surface crossings during dissociation.

There has been very significant progress recently on the interpretation and theoretical calculation of orientation and alignment parameters.\(^7\)\(-\)\(^9\) In the case of photodissociation of diatomic molecules, the electronic angular momentum polarization gives information about excitation of multiple electronic states and nonadiabatic effects.\(^6\)\(^3\)\(-\)\(^9\) These processes can lead to interferences between wave packets evolving on different potential energy curves that correlate with the same asymptotic product states, and the polarization moments provide a very direct probe of these phenomena.\(^10\) Although much of the theoretical work to date has been confined to diatomic molecules, extension to triatomic systems, albeit with more approximate methods, is by no means beyond the current scope. Simpler models that allow interpretation of electronic polarization information in more complex systems are also available.\(^8\)\(^11\)\(-\)\(^14\)

In the present work, the technique of velocity-map ion imaging is applied to a study of the molecular photodissociation of ozone at 193 and 205 nm, which has allowed determination of $O(1D_2)$ electronic orbital angular momentum polarization moments up to and including those of rank $K=4$. Together with information about the velocity distribution of the $O(1D_2)$ photofragments, we illustrate how the polarization data can be used to unravel the dissociation mechanism(s) operating in this important atmospheric molecule. Before we move on to describe the experimental and analytical procedures employed, we provide a brief résumé of previous work on ozone most relevant to the present study.

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B. Previous relevant work on ozone at 193 nm and other wavelengths

Ozone photochemistry has been the subject of extensive experimental and theoretical study due to its atmospheric importance and complex dissociation dynamics. The uv absorption spectrum of ozone is dominated by the intense Hartley band, which results in stratospheric O3 acting as a highly efficient filter of solar radiation in the wavelength range of 220–310 nm. At lower energies, the Hartley band is partially overlapped by the much weaker Huggins band. On the short wavelength side of the Hartley band, the absorption cross section shows a minimum at 201 nm ($3.1 \times 10^{-19}$ cm$^2$) and then increases again with shortening wavelength, indicating the onset of a new absorption feature. This short wavelength region has been rarely investigated due to the unfavorable absorption cross sections, and also because the solar flux at wavelengths around 193 nm is low enough that dissociation in this region was considered unimportant. However, around 190 nm, gaps in the Schumann-Runge bands of oxygen allow light to penetrate into the upper and middle stratosphere, and so ozone photolysis in this region, producing O($^1D_J$), remains of significance.

There are a total of six spin-allowed, energetically accessible product channels open in the photodissociation of ozone at 193 nm. Detailed product quantum yields for all of the accessible channels have not yet been determined, although Turnipseed et al. reported $\phi$(O($^3P_J$)) = 0.46±0.29 and $\phi$(O($^5P_J$)) = 0.57±0.14. The fact that the sum of the quantum yields is greater than unity was ascribed to the operation of a dissociation channel producing three O($^3P_J$) atoms. A quantum yield for O$_2$(b $^1Σ_g^+$) of 0.50±0.38 was also measured, and it was concluded that O($^3D_J$) and O$_2$(b $^1Σ_g^+$) were generated as coproducts. These results are in good agreement with the findings of a vacuum ultraviolet (vuv) laser induced fluorescence (LIF) experiment by Takahashi et al. and Nishida et al., who measured $\phi$(O($^1D_J$)) = 0.48±0.03. A high-resolution photofragment translational spectroscopy study carried out by Stranges et al. also allowed estimates to be made of the branching ratios for several of the energetically open channels; most notable in the context of the present work are those into the channels O($^1D_J$)+O$_2$(a $^1Δ_g$) and O($^3D_J$)+O$_2$(b $^1Σ_g^+$) (see Sec. III). Note that the former channel is dominant [e.g. $\phi$~0.9 (Ref. 16)] following excitation into the peak of the Hartley band. Evidence was also found for the three-atom dissociation channel. Recently, Takahashi et al. detected the O($^3S$) atoms from ozone photolysis at 193 nm, again using the vuv-LIF technique, and determined a quantum yield of 2.5 x 10$^{-5}$. Compared with the Hartley band, rather little attention has been paid to the photodissociation dynamics following 193 nm excitation. In the photofragment translational spectroscopy study of Stranges et al., both the speed distribution and translational anisotropy of the recoiling O-atom fragments were determined, with the latter found to be opposite in sign to that observed subsequent to excitation into the Hartley band. These measurements were corroborated by the vuv-LIF experiments of Takahashi et al., who obtained low-resolution speed and translational anisotropy parameters for both O($^1D_J$) and O($^5P$) photofragments separately from Doppler-resolved LIF measurements.

In addition to the above-mentioned work, a number of O($^1D_J$) atomic polarization studies at longer wavelengths are deserving of special mention. The first of these is the important angular momentum alignment study performed by Dylewski et al. They observed preferential population of $M_J=1$ levels at selected wavelengths across the Hartley band between 235 and 305 nm but also found evidence for contributions from coherence terms. The $M_J$ population results were interpreted as being consistent with dissociation on a state of $A'$ symmetry. More recently, Lee et al. carried out an elegant slice imaging study of ozone photofragmentation via the Hartley band, in which they observed strong, speed-dependent O($^1D_J$) orientation following absorption of plane polarized light. The origin of this orientation was ascribed principally to a “static” coherence effect, arising from coherent excitation of both parallel and perpendicular (in plane) components of the transition moment with respect to the recoil direction of the atom. It was argued that this coherence could persist, or be quenched, by nonadiabatic transitions at avoided crossings or in the long range region and could also be influenced by long range O–O–O bond angle dependent effects, similar to those described in the case of the photodissociation of N$_2$O Ref. 11 and SO$_2$. Finally, very recently Denzer et al. have published an in-depth resonantly enhanced multiphoton ionization (REMPI) time-of-flight study of the O($^1D_J$) atomic orbital polarization generated by photofragmentation in the long wavelength region of the Hartley band at 298 nm. They were able to determine moments of the angular momentum distribution of rank up to $K=4$, including $K=1$ orientation moments. Their results were in general agreement with those of both Dylewski et al. and Lee et al. However, compared with the former work, Denzer et al. observed a small coherence term, and it was suggested that this reflected the nuclear geometry of the dissociating molecule and the existence of possible nonadiabatic effects in the exit channel. Of particular note in the present context, and in agreement with the study by Lee et al., Denzer et al. found that of the three possible $K=1$ orientation moments, only that arising from the static coherence mechanism ($\chi'$) was nonzero. We return to a discussion of this point in Sec. IV E.

II. METHOD

A. Experimental procedures

The experiments were carried out using a standard velocity-map imaging apparatus that has been described in detail previously. The 15% of an ~1:1 O$_3$/O$_2$ mixture, prepared using a commercial ozonizer, was seeded in He at a backing pressure of ~1 bar and expanded through a pulsed nozzle (General Valve) with a 0.5 mm diameter orifice. The molecular beam was collimated by a 1 mm diameter skimmer. The rotational temperature of the skimmed molecular beam was believed from our previous work on the same apparatus to be ~50 K. Further downstream, the molecular beam was passed through a 2 mm hole in the repeller...
plate of the velocity mapping ion optics assembly and crossed 5 cm away from the nozzle exit by two counterpropagating laser beams. The 193 nm photolysis radiation was provided by a Lambda Physik EMG103 excimer laser, while the probe radiation, operating typically at pulse energies of around 500 μJ pulse⁻¹, was obtained by frequency doubling the output of a Nd:YAG (yttrium aluminum garnet)-pumped dye laser system (Photonic Solutions). The time delay between the two laser pulses was ~10 ns. Two lenses of 30 cm focal length were used to focus the radiation onto the molecular beam.

The O(¹D₂) photofragments were probed by (2+1) resonantly enhanced multiphoton ionization (REMPI) via the ¹F₁→¹D₂ and ¹P₁→¹D₂ (labeled in the following as the ¹F₁ and ¹P₁) transitions at 203.8 and 205.4 nm, respectively. During image acquisition, the probe laser wavelength was scanned over the Doppler profile of the O(¹D₂) transitions in order to ensure an equal detection sensitivity for all the product velocities. The oxygen ions were velocity mapped onto an imaging detector consisting of a pair of 40 mm chevron double microchannel plates (MCPs) coupled to a P47 phosphor screen (DelMar Ventures). The image on the phosphor was digitized by a frame grabber (DelMar Ventures) and captured by an intensified charge coupled device (CCD) camera (Photonic Science), electronically gated to the flight time of the detected ions, and subsequently sent to a personal computer (PC) for signal processing (thresholding, event counting, and accumulation). Images were averaged over 25 000 laser shots. Velocity calibration of the final images was achieved using images of O(³P J) from the photolysis laser. The analysis of angular momentum polarization effects on ion images was previously achieved by fitting basis functions, generated in terms of a semiclassical bipolar moment formalism, to Fourier moments of the ion images. In the present work, we proceed along similar lines but start from the quantum mechanical expression for the O-atom REMPI signal intensity in pump-probe geometry j given by Suits and co-workers, adapted for multiple velocities,

\[ I(\Theta,\Phi;v) = C \sum_{KQ} \rho_{KQ}(\Theta,\Phi;v) - C^{\ast}_{KQ}(\Theta',\Phi'). \]

In this equation C is a normalization constant, \( \rho_{KQ} \) are line strength factors for the REMPI transition, \( C_{KQ}(\cdot) \) is a modified spherical harmonic, and \( \Theta',\Phi' \) define the direction of the probe (analysis) laser reference vector in the laboratory frame (see Sec. IV A). The latter is defined such that the Z axis lies parallel to the time-of-flight (TOF) axis and the X axis lies along the propagation direction of the photolysis laser. The \( \rho_{KQ}(\Theta,\Phi;v) \) are speed-(and geometry j) dependent elements of the density matrix for the angular momentum distribution, or the angular momentum state multipoles of rank K and component Q at some laboratory frame scattering angle defined by \( (\Theta,\Phi) \). The moments have been defined by Siebbeles et al. and subsequently modified for application to ion imaging as,

\[ \rho_{KQ}(\Theta,\Phi;v) = \frac{3}{4\pi} \left( \frac{[P]}{[A]} \right)^{1/2} \sum_{k,q,q'} (-1)^{K+q'} \times \delta_{q,q'} \times \left( \frac{1}{q' - q - Q} \right) \times D_{Q'Q}(\Phi,\Theta,0) D_{q'Q}(\Phi,0). \]

In this equation, \([i]=2i+1\), the terms in large brackets are 3-J symbols, the \( D(\cdot) \) are Wigner rotation matrices, and \( Q'=q' - q, j_A \) is the angular momentum of the probed atom, which in the case of O(¹D₂) is equivalent to L. \( E_{k,q}(\cdot) \) are the matrix elements of the polarization tensor of the photodissociation light for geometry j and rank \( k_d \) (Ref. 5) and are defined as.
\[ E_{k,d}(e_j) = (-1)^n[k_d]^{1/2} \left( \begin{array}{cc} 1 & 1 \\ s & -s \\ 0 & 0 \end{array} \right) C_{k,d}(\Theta_d', \Phi_d'), \]
\[ \rho_{KQ}(q,q',i,n;j;r) \]
\[ = \frac{3}{4\pi} \left( \frac{[K]}{[J_{1/2}]} \right)^{1/2} \sum_{k_d} (-1)^{K+Q} \]
\[ \times E_{k,d}(e_j)[k_d]^{1/2} \left( \begin{array}{cc} 1 & 1 \\ q' - q & -Q' \end{array} \right) \]
\[ \times \sum_{K'} \left[ K K' \left( \begin{array}{cc} K & k_d \\ -Q & n \end{array} \right) \left( \begin{array}{cc} K & k_d \\ -Q' & 0 \end{array} \right) \right] \]
\[ \times \left[ 1 + (-1)^{K+n} \right] \int N e^{-\sigma(v_i - v)^2} \]
\[ \times C_{K'n}(\Theta(0))dv_z. \]

The experimentally acquired images can also be cast in a form analogous to Eq. (5) by performing a Fourier integral:
\[ I(r, \Phi) = C_0(r) + 2 \sum_{n=0}^\infty C_n(r) \cos n\Phi + C_n'(r) \sin n\Phi. \]

In fitting the simulated Fourier moments of the images to the experimental data, the quality of the fit \((\chi^2)\) was calculated by directly comparing experimental and fitted Fourier moments of equivalent order \(n\). In this work, the fitted parameters \(a_k(q,q';i)\) for a given number of Gaussians, \(i_{\text{max}}=25\), were selected using a genetic algorithm from a pool of parameters randomly generated within the constraints of the physically allowed values. A Monte Carlo procedure, described fully elsewhere,\(^{38}\) was then used to determine the errors on the fitted coefficients and on the speed-dependent moments of the angular distribution.

To reduce the number of fitting coefficients, the speed dependence of the polarization parameters was described in terms of a Legendre polynomial expansion. The fitted coefficients for all polarization parameters [excluding those for the speed distribution, characterized by the set of coefficients \(a_0(i) = a_0(0,0;i) + 2a_0(1,1;i)\)] were expressed as a convolution of the speed distribution parameters and a Legendre polynomial \([P_i(\cdots)]\) series with \(n\text{leg}\), the maximum order, set to 2:
\[ a_k(q,q';i) = a_0(i) \sum_{i=0}^{n\text{leg}} b_k(q,q';i) P_i(x_i), \]
where
\[ x_i = 2 \frac{i-1}{i_{\text{max}}-1} - 1. \]

This procedure allowed for the speed dependence of the polarization parameters while significantly reducing the basis set size.

Correct analysis of the ion images requires careful renormalization of the raw data, taking into account the integrated (over speed) angular momentum orientation and alignment. In the fitting procedure used in the present work, renormalization was done on the fly using the expressions for the integrated image intensity given in Appendix A.

In the case of diatomic molecules, the dynamical functions have been interpreted in terms of contributions to the
dynamics by transitions to the dissociation continuum with their transition dipole moment vectors lying either parallel \((q, q' = 0)\) or perpendicular \((q, q' = \pm 1)\) to the bond axis. Interpretation of the polarization parameters in the case of larger systems, \(23,39\) is more problematic and awaits a detailed theoretical treatment. In the present work we have used the polarization parameters to calculate the electron hole distribution of the atomic photofragment, representing the polarization of the electron hole, and the atomic product quantum mechanical angular momentum distribution, which represents the polarization of the electronic orbital angular momentum \(L\). These quantities provide a useful and rigorous means of interpreting the polarization parameters, which remain valid even for polyatomic systems such as ozone. The distributions can be calculated from a knowledge of the molecular frame state multipoles, \(\rho_{KQ}^{mol}(\theta, \phi; v)\), using the equations developed from those by Brouard et al.\(^{12}\) and de Miranda and co-workers.\(^{40,41}\) Equations for \(\rho_{KQ}^{mol}(\theta, \phi; v)\) are given in Appendix B. In the following, the molecular frame moments are integrated over the recoil speed, and these integrated moments will be referred to as the \(\rho_{KQ}^\text{mol}(\theta, \phi; v)\) molecular frame moments.

III. RESULTS

A. Images, moments, and fits

Representative ion images, their Fourier moments and the fits to those moments are shown in Fig. 1. For the pump-probe data at 193 nm, simultaneous fits were performed on 28 separate ion images, obtained using either different pump-probe geometries or different REMPI probe transitions. Of particular significance to the present work is the fact that these data also included the moments of ion images obtained in geometries sensitive to orientation (see further below). The probe-only data, obtained with just the REMPI laser firing at wavelengths around 205 nm, were also analyzed. These data were not sensitive to orientation moments of the angular momentum distribution but, in principle, could be used to determine angular momentum alignment. However, it proved impossible to obtain enough independent probe-only geometries to allow successful extraction of all of the polarization parameters, and thus, a reduced set was determined by setting all \(K=4\) moments to zero. This is a reasonable first approximation since these higher moments tend to have less influence on the image data than the low-order moments, particularly for the \(F_1\) transitions which have smaller \(K=4\) line strength factors (i.e., \(P_4 < P_2\)). The reduced set of probe-only geometries (HH, VV, LL, and RR) was also found to be insensitive to the \(\gamma_2\) alignment parameter, and equally good fits could be obtained when this parameter was set to zero. It should be emphasized that in spite of these uncertainties, the returned speed distribution and speed-dependent \(\beta(v)\) parameter obtained from the analysis were found to be fairly insensitive to the number of moments used in the analysis.

B. Speed distributions and \(\beta(v)\)

The speed distributions obtained at 193 and 205 nm, together with the associated speed-dependent translational anisotropy parameters \(\beta(v)\), are presented in Fig. 2. Although the speed distributions at the two wavelengths are similar, that obtained at the shorter wavelength shows increased intensity at low translational energies, and there is a hint of bimodality in the distribution that is completely absent at 205 nm, possibly indicative of the emergence of a second dissociation channel. As shown by the comparison in Fig. 2(b), the 205 nm data are in excellent agreement with the speed distribution obtained in the recent quantum mechanical (QM) study of the dissociation dynamics by Baloitcha and Balint-Kurti following excitation into the short wavelength region of the Hartley band.\(^{42,43}\) The comparison lends support to the notion that, notwithstanding the slight reduction in \(\beta(v)\) parameter (see below), even at wavelengths as short...
as 205 nm, the dominant pathway to formation of O(1D2) photofragments remains dissociation on the diabatic B-state potential energy surface (PES).

The speed and translational anisotropy distributions obtained in the present work are in reasonable agreement with those measured by Stranges et al.\textsuperscript{19} using the molecular beam, time-of-flight mass spectrometry technique [see Fig. 2(a)], although the results from the latter study show somewhat more pronounced structure. It is worth noting, however, that the previous experiments were insensitive to the electronic state of the atomic oxygen. In the kinetic energy region of interest, Stranges et al.\textsuperscript{19} tentatively associated the bimodality in the speed distribution to production of O2 in two electronically excited states: the a1Δg state (with a quantum yield of 48%), which is the dominant coproduct subsequent to excitation into the peak of the Hartley band\textsuperscript{16} and the b1Σg+ state (with a quantum yield of 23%). A very crude integration of the speed distribution obtained here at 193 nm into fast (v > 3000 m s\textsuperscript{-1}) and slow (v < 3000 m s\textsuperscript{-1}) components yields branching fractions of 0.59 (fast) and 0.41 (slow), similar to those reported by Stranges et al.\textsuperscript{19} Although the speed distribution obtained in the present work is insufficiently structured to allow definitive identification of the electronic state of the O2 cofragment, independent (albeit indirect) evidence for production of O2(b1Σg+) at 193 nm has been provided by Turnipseed et al.\textsuperscript{17} Interestingly, their measurements\textsuperscript{17} which involved the observation of the slow production of O(3P) from reaction of O2(b1Σg+) with O2, suggested production of O2(b1Σg+) at 193 nm, but not at 205 nm, consistent with the different forms for the speed distributions observed here at the two wavelengths. Finally, with respect to the molecular beam study by Stranges et al.,\textsuperscript{19} it should be mentioned that the present measurements seem to rule out assignment of the two peaks they observed in the translational energy release distribution at low kinetic energies (labeled 5 and 6 in Ref. 20) to channels involving O(1D) production.

The speed and spatial anisotropy data obtained at 193 nm are also in excellent agreement with the low-resolution Doppler resolved LIF measurements of Takahashi et al.\textsuperscript{18} Note that excitation of ozone at 193 nm yields a negative translational anisotropy (with a speed-integrated value of -0.21, as shown in Table I), whereas excitation into the peak of the Hartley band yields a positive β parameter, greater than unity.\textsuperscript{18,22} The β parameter obtained at probe-only wavelengths around 205 nm (+0.59) is significantly smaller than the value associated with absorption into the peak of the Hartley band, although unlike the value obtained at 193 nm it has the same sign. The reduced β parameter at 205 nm, compared with that obtained, for example, at 255 nm (where it is 1.29)\textsuperscript{22} might be taken as evidence for an increasing contribution from a perpendicular transition with decreasing wavelength. However, in the light of the excellent agreement between the photofragment speed distribution obtained here and the QM calculations of Balitcho and Balint-Kurti,\textsuperscript{42,43} which include only the diabatic B-state potential, it seems possible that the reduction in β at 205 nm might be associated with other factors, such as a change in the direction of the transition moment with geometry or an increase in the O2(a1Δg) rotational excitation at shorter wavelengths. The recent \textit{ab initio} calculations of the transition moment carried out by Balitcho and Balint-Kurti\textsuperscript{42} certainly support the former notion, with the ratio of the two in-plane components of the transition dipole moment chang-

### Table I. Speed-integrated angular momentum polarization parameters for O(1D2) products of O3 photodissociation at 193 nm. Numbers in parentheses give the 2σ experimental errors in the final decimal place(s). The final two columns refer to the same moments broken down into contributions from slow (<3000 m s\textsuperscript{-1}) and fast (>3000 m s\textsuperscript{-1}) O(1D2) photofragments, which contribute 41% and 59%, respectively, to the total product signal.

<table>
<thead>
<tr>
<th>Polarization parameter</th>
<th>This study</th>
<th>Slow</th>
<th>Fast</th>
</tr>
</thead>
<tbody>
<tr>
<td>β</td>
<td>-0.213(34)</td>
<td>-0.188</td>
<td>-0.230</td>
</tr>
<tr>
<td>α1</td>
<td>-0.014(9)</td>
<td>-0.005</td>
<td>-0.020</td>
</tr>
<tr>
<td>γ1</td>
<td>0.031(20)</td>
<td>-0.003</td>
<td>0.055</td>
</tr>
<tr>
<td>γ2</td>
<td>-0.018(17)</td>
<td>-0.039</td>
<td>-0.004</td>
</tr>
<tr>
<td>s1</td>
<td>-0.071(6)</td>
<td>-0.074</td>
<td>-0.068</td>
</tr>
<tr>
<td>α2</td>
<td>-0.007(4)</td>
<td>-0.023</td>
<td>0.005</td>
</tr>
<tr>
<td>γ2</td>
<td>-0.019(9)</td>
<td>0.006</td>
<td>-0.035</td>
</tr>
<tr>
<td>ζ1</td>
<td>0.013(7)</td>
<td>-0.018</td>
<td>0.035</td>
</tr>
<tr>
<td>α4</td>
<td>0.056(17)</td>
<td>0.029</td>
<td>0.075</td>
</tr>
<tr>
<td>α4</td>
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<td>0.021</td>
<td>-0.022</td>
</tr>
<tr>
<td>η1</td>
<td>0.002(24)</td>
<td>0.089</td>
<td>-0.057</td>
</tr>
<tr>
<td>η2</td>
<td>-0.005(17)</td>
<td>0.012</td>
<td>-0.016</td>
</tr>
</tbody>
</table>

FIG. 2. Speed distributions and β(v) parameters obtained at (a) 193 nm and (b) 205 nm. In (a) the dotted line corresponds to that portion of the O-atom speed distribution observed by Stranges et al. assigned by them to O(1D) atom production (Ref. 20). In (b) the experimental data are compared with the QM calculations of Balitcho Balint-Kurti (Ref. 43), which are shown as the shaded data. The latter were performed for excitation into the blue end of the Hartley band and employed the calculated B(3A') state potential, by the same authors (Ref. 42).
ing significantly with compression or extension of one of the O–O bonds away from C2v configurations. Indeed, contributions to the Franck-Condon region accessed at 193 nm from configurations with reduced O–O bond extensions of only a few tenths of a bohr, which seems reasonable based on the energetics and the published PESs, could significantly reduce β from the value observed at the maximum in the Hartley band.

To summarize the above evidence, it would appear that at 205 nm the predominant absorption is into the Hartley band, leading to similar photodissociation dynamics to that previously observed at longer wavelengths. In contrast, excitation at 193 nm appears to involve predominantly a transition to a state of either \(^1A_1(1^1A')\) or \(^1B_1(1^1A')\) symmetry, both of which could be associated with negative translation anisotropy parameters. The velocity distribution of the O atoms at this wavelength provides supporting evidence for the operation of two dissociation pathways leading to \(O_2\) molecular cofragments born in either the \(a\) or the \(b\) states. These two pathways could either reflect photon absorption via a mixed transition or curve crossings along the dissociation pathway. The similarity of the high speed component of the \(O(1D_3)\) speed distribution with that predicted and observed for the Hartley band suggests that some contribution from excitation into the tail of the Hartley band should not be ruled out, even down to wavelengths as low as 193 nm.

![FIG. 3. O(1D3) alignment polarization moments for (a) rank \(K=2\) obtained at 205 nm, (b) rank \(K=2\) at 193 nm, and (c) rank \(K=4\) at 193 nm. The symbols employed, which are not to be confused with individual data points, are as follows: \(s_{K=2} \cdots \), \(a_{K=3} \cdots \), \(y_{K=4} \cdots \), and \(y_{K=3} \cdots \). Note that \(y_{K=4}\) was not determined at 205 nm. The dotted line marks zero on the plots.](image)

**C. Polarization parameters**

The speed dependences of the alignment parameters with \(K=2\) and 4 are presented in Figs. 3(a)–3(c). Like the \(\beta(v)\) parameters shown in Fig. 2, the polarization parameters have been normalized to the speed distribution. For this reason the values of the parameters below 1000 m s\(^{-1}\) and in excess of 4000 m s\(^{-1}\) (not shown in the figures) should be treated with caution, since very few fragments are born with speeds outside these limits. The data obtained at 193 nm were sufficient for moments with \(K=4\) also to be determined. Although the contribution of these higher-order moments to the signal is small, the quality of the fits was improved significantly by their inclusion in the analysis, particularly in respect of the \(P^\dagger\) versus the \(F^\dagger\) transitions. In practice, only one of the fourth-order moments \(s_4\) makes a large contribution to the data, as shown more clearly once the parameters are integrated over speed (see below).

In the case of the orientation moments, data were recorded using \(P^\dagger\) and \(F^\dagger\) REMPI transitions, employing both linearly and circularly polarized photolysis radiation. Difference images, obtained after subtraction of images obtained on alternate laser shots using left and right circularly polarized probe light, are shown in Fig. 4. While the difference data for the \(P^\dagger\) data show a clear orientation signal, those for the \(F^\dagger\) transition are too small to be observed within the signal-to-noise ratio of the experiment. The line strength factors for the two transitions, i.e., the \(P_K\) term in Eq. (1), are very different, with the difference signal for the \(P^\dagger\) transition being sensitive only to the odd orientation moments with \(K=1\) and the \(F^\dagger\) transition exclusively sensitive to the \(K=3\) moments of the distribution. Thus, the difference images shown in Fig. 4 provide direct evidence that the \(K=1\) moments of the distribution are significant, while the \(K=3\) moments are considerably smaller and cannot be determined. The speed-dependent laboratory frame orientation parma-
parameters returned by the Fourier-moment analysis are shown in Fig. 5(a). A particularly interesting feature, discussed further in Sec. IV E, is that nonzero values are obtained for all three \( K=1 \) moments.

The speed-integrated polarization parameters are listed in Tables I and II. In the final two columns of Table I, which summarizes the data obtained at 193 nm, the polarization parameters are resolved into fast (\( v > 3000 \text{ m s}^{-1} \)) and slow (\( v < 3000 \text{ m s}^{-1} \)) components. This was achieved by separate integration over the two speed components, followed by renormalization according to their relative probabilities. Some of the moments change sign between the two speed regimes, while others differ in magnitude between the two components, and in a more complete analysis one might explore the polarization of the two speed components separately. However, it is evident from the data in Table I that the speed-integrated moments are generally dominated by the contribution from the major, high speed component. Thus in the present work we focus only on the speed-integrated moments, which tend to reflect the behavior of the fast \( \text{O}(^{1}D_{2}) \) photofragments. Table II includes a comparison between the polarization moments obtained in the present work at 193 and 205 nm with those derived in the previous work by Dylewski et al. at 255 nm. Bearing in mind the difficulties noted in Sec. II B concerning the analysis of the probe-only data, there is broad agreement between the data sets at 205 and 255 nm. This agreement gives strength to the argument that photodissociation at 205 nm occurs predominately via the diabatic \( B \) state. Note, on the other hand, that at 193 nm, where a different transition clearly dominates, the \( K=2 \) polarization moments are significantly reduced in magnitude compared with those observed in the previous work of Dylewski et al. at 255 nm.

### IV. DISCUSSION

#### A. Electron density and angular momentum distributions

As was noted in Sec. II B, further interpretation of the polarization parameters requires transformation into the molecular frame. Following previous work, the latter is defined with the \( z \) axis lying along the recoil direction and the \( zx \) plane containing the reference axis of the photolysis radiation, taken to be either the polarization (for linearly polarized light) or propagation (for circularly polarized light) direction of the photolysis radiation. The transformation of the polarization parameters is a generalization of that described by Bracker et al. and is given in Appendix B. In the following, we focus exclusively on the data obtained using linearly polarized photolysis and return to the issue of circularly polarized photodissociation in Sec. IV E. For the reasons discussed in Sec. III, we also concentrate on the speed-integrated behavior. The single molecular frame moment for \( K=1 \) obtained using linearly polarized photolysis at 193 nm is shown in Fig. 5(b), while the even-moment data obtained at 193 and 205 nm are shown in Fig. 6. In these figures the moments are plotted as functions of recoil angle \( \theta \), which represents the angle between the recoil velocity and the reference axis of the photolysis radiation. Figure 6 also includes the molecular frame moments derived from the polarization parameters measured by Dylewski et al. at 255 nm. The \( K=2 \) moments are dominated by the \( \rho_{20}^{\text{mol}} \) parameter, which is negative in sign, although at 205 and 255 nm this moment is reduced somewhat in magnitude for scattering angles centered around 90°. These results are consistent with the predominant population of low \( M_{L} \) states, and \( M_{L}=0 \) for \( \theta \sim 0^\circ \) and 180°, i.e., \( L \) alignment is preferentially perpendicular to the recoil direction, as noted previously by Dylewski et al. for excitation into the Hartley band. In contrast with this, the magnitude of \( \rho_{20}^{\text{mol}} \) is significantly reduced at 193 nm compared to the values obtained at longer wavelengths. The \( K=4 \) moments obtained at 193 nm are also dominated by a single term, that for \( \rho_{40}^{\text{mol}} \).

The molecular frame moments can be used to construct both the angular momentum distribution of \( L \) (Refs. 40 and 41) and the electron density (or in this case the electron hole) distribution of the atomic oxygen. Explicitly, the expression employed for the quantum mechanical angular momentum distribution is

**TABLE II.** Comparison of the speed-integrated angular momentum alignment parameters for the \( \text{O}(^{1}D_{2}) \) products of \( \text{O}_3 \) photodissociation at 193 nm with the present results at 205 nm and the previous work at 255 nm by Dylewski et al. (Ref. 22). For comparison with the present work, the latter has been averaged over O-atom speed. Numbers in parentheses give the (2\( \sigma \)) experimental error in the final significant figures. Here we employ the notation of Rakitzis and Zare (Ref. 3), as defined by Alexander (Ref. 39).

<table>
<thead>
<tr>
<th>Alignment parameter</th>
<th>This study at 193 nm</th>
<th>This study at 205 nm</th>
<th>Dylewski et al. at 255 nm</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \beta )</td>
<td>-0.21(3)</td>
<td>0.59(10)</td>
<td>1.29</td>
</tr>
<tr>
<td>( a_{20}(\parallel) )</td>
<td>-0.36(6)</td>
<td>-0.86(7)</td>
<td>-0.84</td>
</tr>
<tr>
<td>( a_{20}(\perp) )</td>
<td>-0.35(3)</td>
<td>-0.53(10)</td>
<td>-0.35</td>
</tr>
<tr>
<td>( a_{20}(\parallel, \perp) )</td>
<td>-0.17(8)</td>
<td>0.00(80)</td>
<td>-0.23</td>
</tr>
<tr>
<td>( a_{20}(\perp, \perp) )</td>
<td>-0.05(3)</td>
<td>-0.27(7)</td>
<td>-0.08</td>
</tr>
</tbody>
</table>

\[ \text{FIG. 5. (a) As for Fig. 3, but showing the laboratory frame } K=1 \text{ orientation moments obtained at 193 nm, with the notation}_1 \gamma_1-, \gamma_1-, \gamma_1- \text{. (b) the experimentally derived molecular frame moment } J_{[pn]} \text{ arising from photodissociation with linearly polarized light at 193 nm.} \]
where the $a_Q^{(K)}$ polarization parameters of Refs. 40 and 41 are equivalent to $\rho_{KQ}^{\text{mol}}/[K]^{1/2}$ used here. $(LL, K0)\mid LL\rangle$ represents a Clebsch-Gordan coefficient, and the angles $(\theta_L, \phi_L)$ define the direction of the minimum uncertainty state $|0q\rangle$ in the molecular frame. To obtain information about the electron density distribution, we make use of the equation for the case of a two-electron atom,

$$Q(\theta_L, \phi_L) = \sum_{KQ} \frac{|K|}{4\pi} a_Q^{(K)}(LL, K0)\mid LL\rangle C_{KQ}^Q(\theta_L, \phi_L),$$  

(9)

which was derived in Appendix B of Ref. 12. Because $O^+(D_2)$ has a $p^6$ configuration, as opposed to $p^2$, the distribution worked out using Eq. (10) actually represents the unfilled orbital density or the electron hole distribution. Therefore, the angles $(\theta_L, \phi_L)$ in Eq. (10) define the angular coordinates of the electron hole. The term in curly brackets in the same equation is a 6-J symbol. Note that unlike the $L$ distribution, the Clebsch-Gordan coefficient appearing in Eq. (10) ensures that the electron hole distribution only contains even terms, with $K=0$ and 2.

The two distributions are shown in Fig. 7 for a recoil angle of 45° for photolysis wavelengths of 193, 205 and 255 nm, where the latter employs the alignment parameters of Dylewski et al.22 Inspection of the molecular frame moments shown in Fig. 6 reveals that qualitatively similar results will be obtained at recoil angles other than 45°. As noted already, the derived distributions assume linearly polarized photolysis, although analogous plots can be constructed in the case of circularly polarized radiation. The $L$ distribution obtained at 193 nm includes the small contribution arising from the orientation moment $\gamma_2^L$, leading to a small preference for orientation along the $-y$ direction. In confirmation of the above qualitative discussion, the angular momentum distributions [shown in the left-hand figures labeled (a)] peak perpendicular to the recoil direction, particularly so at longer wavelengths. Furthermore, in contrast to the $L$ distributions, the corresponding electron hole distributions generated at 205 and 255 nm (Ref. 22) peak along the recoil axis $xz$. Thus, at these photolysis wavelengths, one may interpret the electron hole distribution as consisting primarily of the unfilled $p$ orbitals lying along the recoil direction, i.e., two $p_z$ holes. This is consistent with the $L$ distribution, which suggests that the orbital angular momentum $L$ together with the filled $p_{x,y}$ orbitals lie perpendicular to the recoil direction. This polarization implies that the recoiling $O^+(D_2)$ atoms are aligned in a $\Sigma$ symmetry configuration, as discussed further in Sec. IV D. Note, finally, that whereas the electron hole distributions are quite strongly aligned along the $z$ axis subsequent to excitation directly into the Hartley band, at 193 nm the distribution has more amplitude in the molecular frame $xy$ plane, a point we return to in Sec. IV D.

**B. Long range potentials**

In this section we follow the procedure employed to model polarization effects in the singlet channels of the photo-
todissociation of N₂O (Ref. 11) and NO₂.¹⁴ We have also previously employed a similar treatment to calculate the polarization moments of O(¹P) in the photofragmentation of N₂O and SO₂.¹³ We restrict attention only to the channel generating O(¹D₂)+O₂(a¹Δ₈) since this appears to be the channel that dominates and, in particular, has the most influence over the speed-integrated polarization parameters. It should be stressed at the outset that the model is an approximate one, intended to provide some insight about the interpretation of the polarization parameters measured here and elsewhere.²²,²³,²⁵ A quantitative analysis of these results would require a full quantum dynamical treatment of the photodissociation.

We start by defining a critical separation Rₖ in the exit channel and assuming the following.

- Before the critical separation Rₖ, the system evolves on some potential energy surface (the model does not say which one).
- The electronic wave function at this critical point is determined by first-order electrostatic interactions.
- After this point the fragments lose any contact and the wave function, transformed to the appropriate frame, determines the observed orbital angular momentum polarization of the oxygen atoms.
- Effects of coherent excitation of two surfaces are ignored.

In this section we take the molecular frame z axis to lie along the Jacobi coordinate R, with the diatomic molecule lying in the zx plane at some Jacobi angle γ. In Appendix C we give explicitly the γ dependence of interaction potential energy matrix Vₖₙ₉ in the long range region. This can be constructed in a basis |ALM⟩, where |A⟩ is the wave function for diatomic molecule A and |LM⟩ is the electronic orbital wave function for atom B (here we rewrite ℌₖ as ℌ). We have assumed that the long range potentials are dominated by the quadrupole-quadrupole (Q-Q) interaction between O(¹D₂) and O₂(a¹Δ₈), and that, therefore, the choice of Rₖ becomes immaterial in the model. Because of the 1/R² radial dependence of the Q-Q interaction, the latter will inevitably dominate at sufficiently large R. However, the quadrupole moment of O₂(a¹Δ₈), like that of ground state molecular oxygen, is quite small, and it is possible that at the critical distance Rₖ, higher-order electrostatic terms or the dispersion interaction may play a more important role.⁴⁹,⁵⁰ Note that only the anisotropic part of the dispersion interaction will have an effect on the atomic polarization predicted by the model. It is also worth mentioning that the atomic orbital polarization may be influenced by shorter range (anisotropic) exchange interactions. We emphasize that the present qualitative treatment does not take these additional interactions into account: These are points that would merit further consideration in future studies.⁵¹

Slices through the long range potential energy surfaces as a function of Jacobi angle are shown in Fig. 8. In the D₉ₖ point group, the ¹D₂ state of atomic oxygen transforms as Σ(M = 0), Π(M = ±1), and Δ(M = ±2). Therefore, as noted in Appendix C, combination of O(¹D₂) with O₂(¹Δ₈) at linearity generates singlet molecular states of Σ⁺, Σ⁻, Π, Δ, Φ, and Γ symmetries. Note that at linearity the ¹Δ molecular state of ozone arises from combination of O₂(¹Δ₈) with O(¹D₂) atoms aligned with M = 0. Within the present model, each of the five molecular ozone states shown in Fig. 8 is twofold degenerate (leading to ten molecular states in total), with pairs of A′ and A″ levels reflecting the Λ-doublet levels of the O₂(¹Δ₈) coproduct, which in the model are treated as degenerate. Assignment of the ozone molecular states (see Table III) is achieved partly by inspection of the predicted orbital polarization (discussed below) and partly by reference to previous ab initio work.⁴⁴,⁴⁵,⁵²,⁵⁵ Note that for the A′ states we use the diabatic notation employed by Hay et al.⁵² A key feature for the discussion below is the assignment of the lowest energy ozone state to the Δ state at linearity, as opposed to the Σ/Γ pair of states. In the model this is determined by the signs of the quadrupole moments of O₂(¹Δ₈) and O(¹D₂), and we have assumed that both of these are negative, consistent with theory.¹⁴,⁴⁶ Furthermore, the assignment also appears to be consistent with the ab initio calculations of Banichievich⁵³ and Peyerimhoff and Balotin and Balint-Kurti.³⁴ It is worth adding that the ordering of the states also coincides with that predicted on simple chemical grounds, in that one might expect the lowest molecular state of ozone to be formed when the unfilled p orbitals on O are aligned along the molecular axis of O₂, since this will involve minimum occupancy of the O₂σ₉orbitals that will lie higher in energy than the corresponding πₚ orbitals. While

![FIG. 8. Slices through the long range quadrupole-quadrupole potentials as a function of Jacobi angle γ. The notation employed is given in Table III. See text for details.](image)

### Table III. Assignment of the five doubly degenerate surfaces shown in Fig. 8. The Ωₙ, assignments are the adiabatic correlations, which are probably secure for the Λ′ states, which have been characterized recently using high level ab initio theory (Refs. 42–45), but are less clear for the Λ″ states. The “?” symbol indicates that assignment is unclear.

<table>
<thead>
<tr>
<th>State</th>
<th>Symbol</th>
<th>D₉ₖ assignment</th>
<th>Ωₙ assignment</th>
<th>C₂₉ diabatic correlation</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td></td>
<td>Δ</td>
<td>3A′(Σ⁺) and 2A″</td>
<td>1B₁ and 1A₂</td>
</tr>
<tr>
<td>2</td>
<td>-○○○○</td>
<td>Π/Φ</td>
<td>4A′(Π) and 3A″</td>
<td>2A₁ and 2A₂</td>
</tr>
<tr>
<td>3</td>
<td>⋯⋯⋯⋯</td>
<td>Φ/Π</td>
<td>5A'(Ω) and 4A″</td>
<td>2B₂ and 2B₁</td>
</tr>
<tr>
<td>4</td>
<td>⋯⋯⋯⋯</td>
<td>Σ/Γ</td>
<td>6A'(Ω) and A&quot;</td>
<td>3A₁/4A₁ and ?</td>
</tr>
<tr>
<td>5</td>
<td>⋯⋯⋯⋯</td>
<td>Γ/Σ</td>
<td>A' and A&quot;</td>
<td></td>
</tr>
</tbody>
</table>

⁴Labels of Hay et al. (Ref. 52) for the diabatic A′/A″ surfaces (A', B', C', ... ) in parentheses.
The calculated alignment moments can similarly be compared with the experimentally derived moments either obtained in the present work at 205 nm or with those measured previously by Dylewski et al. at 255 nm, both of which are shown in Fig. 6. We start by considering this K = 2 polarization data obtained subsequent to excitation into the Hartley band, at 205 and 255 nm. The experimental polarization moments, and the electron hole and angular momentum distributions derived from them, reveal that O(3P) is formed in a state with preferential population of low M_L states (see Fig. 7), corresponding to an O-atom term of Σ symmetry at linearity. Therefore, as noted above, combination with the O_2(1Δ_g) term would generate an ozone molecular state of 1Δ character at linearity. Recall that the experimental molecular frame moments are dominated by a large negative value for ρ_{20}^{mol}, particularly for recoil angles close to 0° and 180°. If we compare this result with the predicted moments shown in Fig. 10, it is clear that only two of the five pairs of states can account for such large negative ρ_{20}^{mol} moments, namely, the lowest Δ state [labeled (---) in the figure] at low Jacobi angles γ or the highest Σ=1 state [labeled (-----)] at Jacobi angles centered around 90°. The next significant feature of the experimental moments is that the ρ_{22}^{mol}, although small, is negative in sign, in agreement with the sign of the predicted ρ_{22}^{mol} moment for the state labeled Δ at linearity (-----), but opposite in sign to that predicted for the

C. Polarization moments subsequent to excitation into the Hartley band

The eigenfunctions obtained from the above procedure can be used to obtain estimates for the multipole moments of the recoiling O atom via the equation,

\[ \rho_{KQ}(L) = \sum_{MM'} (-1)^{L-M[L, M, L-M']|KQ|} \rho_{MM'}(L), \]

with density matrix elements

\[ \rho_{MM'}(L) = c_M^* c_{M'}. \]

Here \( c_M \) are the expansion coefficients of the wave function for the electronic state of interest. As for \( \rho_{KQ}^{mol} \) (see Appendix B) we have renormalized the \( \rho_{KQ}(L) \) moments such that \( \rho_{00}(L) \) is unity. In the following we have assumed that the critical point \( \lambda_c \) is at sufficiently large \( R \) that it can safely be taken to lie along the recoil direction. Hence we have not transformed the \( \rho_{MM'}(L) \) into the molecular frame used in the experiments, which is defined with \( z \) along the recoil direction. The calculated polarization moments to be presented below depend on the Jacobi angle \( \gamma \). It should be born in mind in the following discussion that a rigorous comparison between the calculated and experimental moments would require a full dynamical calculation, with the averaging over Jacobi angle taken into account explicitly.

The calculated moments are presented in Figs. 9 (for \( K = 1 \)) and 10 (for \( K = 2 \) and 4). The orientation data shown in Fig. 9 may be compared with the experimentally derived molecular frame moment obtained at 193 nm displayed in Fig. 5(b). The calculated alignment moments can similarly be compared with the experimentally derived moments either obtained in the present work at 205 nm or with those measured previously by Dylewski et al. at 255 nm, both of which are shown in Fig. 6. We start by considering this K = 2 polarization data obtained subsequent to excitation into the Hartley band, at 205 and 255 nm. The experimental polarization moments, and the electron hole and angular momentum distributions derived from them, reveal that O(3P) is formed in a state with preferential population of low M_L states (see Fig. 7), corresponding to an O-atom term of Σ symmetry at linearity. Therefore, as noted above, combination with the O_2(1Δ_g) term would generate an ozone molecular state of 1Δ character at linearity. Recall that the experimental molecular frame moments are dominated by a large negative value for ρ_{20}^{mol}, particularly for recoil angles close to 0° and 180°. If we compare this result with the predicted moments shown in Fig. 10, it is clear that only two of the five pairs of states can account for such large negative ρ_{20}^{mol} moments, namely, the lowest Δ state [labeled (---) in the figure] at low Jacobi angles γ or the highest Σ=1 state [labeled (-----)] at Jacobi angles centered around 90°. The next significant feature of the experimental moments is that the ρ_{22}^{mol}, although small, is negative in sign, in agreement with the sign of the predicted ρ_{22}^{mol} moment for the state labeled Δ at linearity (-----), but opposite in sign to that predicted for the
highest state labeled $\Sigma_{1}/\Gamma_{1}$ (---) at linearity. Thus, comparison between the experimental and calculated polarization moments seems to point to dissociation on the lowest ozone molecular state, labeled $\Delta$ at linearity, and the preferential sampling of Jacobi angles $\gamma$ less than $\sim 60^\circ$. Although it is tempting to observe that the $\gamma$-angle dependence of the calculated moments appears to mimic the $\theta$-angle dependence seen experimentally, it should be reemphasized that the dissociation will sample an undetermined range of $\gamma$ angles that could only be quantified by means of a dynamical calculation. However, Jacobi angles $\leq 60^\circ$ could be sampled preferentially, as suggested by the model, if ozone dissociated promptly at bond angles close to that of the ground state equilibrium geometry with little rotational excitation in the $\mathrm{O}_2$ cofragment, which is perhaps not unreasonable for the fastest subset of $\mathrm{O}(1D_{2})$ photofragments that tend to dominate the polarization measurements. Note that the lowest $\Delta$ symmetry state at linearity is precisely the one that is believed to correlate diabatically with the $\tilde{B}1\tilde{B}_2(3^2\Sigma_{1}^+)$ state of ozone and is known to be responsible for the Hartley band$^{42-52}$ as shown in Table III. The comparison between the calculated and observed polarization moments lends some support to the assignment of the electronic states in the long range region given in Table III and to the previously proposed (predominantly diabatic) dissociation mechanism subsequent to excitation into the Hartley band.$^{42-44}$

D. 193 nm polarization data and possible dissociation mechanism

We now turn to the polarization data obtained at 193 nm. Because of the simplicity of the model, assignment of the participating electronic states solely on the grounds of the predicted angular momentum polarization moments would be unwise. In particular, the evidence presented above suggests that excitation at 193 nm may involve at least two electronic states, and under these circumstances coherence effects cannot, in general, be neglected. As noted in Sec. III, the negative translational anisotropy parameter obtained at this wavelength is consistent with excitation to states of either $1B_1$ or $1A_1$ symmetry. The high level $ab\ initio$ calculations of Palmer and Nelson$^{55}$ suggest that excitation to either the $2^1B_1$ or $3^1A_1$ states dominates in this wavelength region, in good agreement with the present findings. Recall that some residual excitation into the $\tilde{B}(1B_2)$ state, as at 205 nm, should also not be discounted. Added to this fact, the speed distribution obtained at 193 nm is consistent with the occurrence of dissociation into two product channels, correlating with $\mathrm{O}_2$ born in either the $a^1\Delta_g$ or $b^1\Sigma^+_g$ electronic states.$^{17,23}$ Although a number of other excited states could well generate products in the same excited electronic states, the presence of the former channel, together with the similarity in the shape of the associated component of the speed distribution to that observed at longer wavelengths, provides some supporting evidence for residual excitation into the tail of the Hartley band. The production of $\mathrm{O}_2(b^1\Sigma^+_g)$, on the other hand, could reflect excitation to a state that correlates with formation of these excited products. If these excited states were then crossed by states correlating with the lower $\mathrm{O}(1D_{2})+\mathrm{O}_2(1\Delta_g)$ asymptote, this might account for the simultaneous production of $\mathrm{O}_2$ in both its $a$ and $b$ electronic states. This behavior would be reminiscent of the crossing between the diabatic $\tilde{B}$ state, which gives rise to the Hartley band and is responsible for production of the $\mathrm{O}(1D_{2})+\mathrm{O}_2(1\Delta_g)$, and the $\tilde{R}$ state, which leads to ground state atomic and molecular photofragments.$^{42,44,46}$

The $ab\ initio$ calculations of the $1A_1'$ surfaces of Schinke and co-workers$^{44,45}$ and Balitocha and Balint-Kurti$^{42,43}$ indicate the presence of a high-lying electronic state of ozone in the correct energy regime that appears to correlate with $\mathrm{O}_2(b^1\Sigma^+_g)$ products. Unfortunately, these high-lying states have not been fully characterized by modern $ab\ initio$ methods, so definitive assignment is difficult. Although we have not attempted a quantitative comparison between the model predictions and the experimental results, it would appear that weighted combinations of excitation into the Hartley band labeled (---) in Figs. 8–10, together with the $2^1B_1$ (┉) and/or $3^1A_1$ (┄) states, might give plausible matches to the observed polarization data. It is interesting to note that the $2^1B_1$ and the $3^1A_1$ states are predicted to lead to $\mathrm{O}(1D_{2})$ atoms with a preferential population of $M=\pm 1$ and $\pm 2$ at linearity, as can be seen from the values of the calculated $\rho_{20}^{\text{mol}}$ moments at $\gamma=0^\circ$ and $180^\circ$, and this probably accounts for the increased probability in the electron hole distribution in the molecular frame $xy$ plane, as was commented on at the end of Sec. IV A.

E. The $K=1$ orientation moments

As was mentioned in Sec. I B, an experimental study of the orientation of $\mathrm{O}(1D_{2})$ subsequent to excitation into the Hartley band at wavelengths between 248 and 285 nm has recently been undertaken by Lee et al.$^{23}$ These experiments revealed a significant, speed-dependent $\gamma'$ parameter, which on average is about a factor of 10 higher than we observe at 193 nm but is of the same sign. The newer work at 298 nm by Denzer et al.$^{25}$ yields a similar $\gamma'$ parameter to those obtained by Lee et al.$^{23}$ If the latter are averaged over recoil speed. Furthermore, once the $\gamma'$ parameters are transformed into the molecular frame, they also agree well with the predicted orientation generated from the state of $1A_1$ symmetry at linearity that we have derived above from the long range model. In the present work at 193 nm, a much smaller value of $\gamma'$ is obtained than was found subsequent to excitation at longer wavelengths into the Hartley band. However, as shown in Fig. 9, with the neglect of coherence effects, simultaneous excitation into the tail of the Hartley band and to either the $2^1B_1$ or the $3^1A_1$ states would be expected to generate small orientation, because the contributions to $\gamma'$ from the two states are of opposite sign.

Interestingly, in the same experiments mentioned above, Lee et al.$^{23}$ found that the other two $K=1$ orientation parameters, $\alpha_1$ and $\gamma_1$, corresponding to orientation along the $z$ and $x$ axes, respectively, were zero. For nonchiral systems, only orientation along the $y$ axis is possible, while for chiral systems orientation along all three axes is possible.$^6$ For photofragment orientation to be observed along the $z$ or $x$ axis, the chirality in the products must be derived from the chirality of
the absorbed (circularly polarized) photon. Lee et al.\(^{23}\) explained the zero values of \(\alpha_1\) and \(\gamma_1\) by noting that, unlike the case in diatomics, for nonlinear polyatomic systems excited to a single potential energy surface, the photon helicity cannot be transferred to the products, because of the quenching of the orbital angular momentum on bending (thus \(A\) and \(\Omega\) are no longer good quantum numbers). In the present work at 193 nm, we find that although all the orientation parameters are small, the \(\gamma_1\) parameter is significantly larger than the \(\gamma_1\) parameter. The question then arises as to whether or not, in contrast to the behavior observed at longer wavelengths, where excitation is believed to be to a single excited state,\(^{23,25}\) the orientation generated at 193 nm could be a consequence of coherent excitation into two electronic states leading to the same asymptotic channel, namely, \(\text{O}^\text{(1}\text{D})_2\) \(+\text{O}_2(\text{\Delta}_\text{g})\), as was suggested to be a possibility in Sec. IV D. This is an issue that needs to be investigated further in future theoretical studies.

V. CONCLUSIONS

The photodissociation dynamics of ozone at 193 nm has been studied using polarized laser photolysis coupled with velocity-map ion imaging. A comprehensive study of the \(\text{O}^\text{(1}\text{D})_2\) electronic orbital angular momentum polarization has been undertaken at this wavelength. Images obtained from dissociation at the probe laser wavelength of around 205 nm were also recorded and analyzed. The 205 nm polarization data were found to be in good agreement with previous studies by Dylewski et al.,\(^{22}\) in which ozone was photolyzed close to the maximum in the Hartley band. Furthermore, the observed \(\text{O}^\text{(1}\text{D})_2\) speed distribution at 205 nm was found to be in excellent agreement with that calculated by Balint-Kurti et al.\(^{23}\) for short wavelength dissociation of ozone via the \(\text{B}^1\text{B}_2(3\,^1\text{A}')\) electronic state. The data reveal that at 205 nm, the dominant excitation and dissociation pathway is the same as it is for the Hartley band.

The experimental data set obtained at 193 nm shows a number of features that mark it out as different from that found at longer wavelengths. In agreement with previous work,\(^{18,26}\) we find that the translational anisotropy has a different sign to that obtained subsequent to excitation into the Hartley band. The speed distribution of the \(\text{O}^\text{(1}\text{D})_2\) atoms also shows signs of bimodality, which can be defined as

\[
\langle P \rangle = \Re(\langle c_0' \rangle),
\]

where

\[
\langle c_0' \rangle = \sum_{KQ} P_K(p_{KQ})C_{KQ}(\Theta_0',\Phi_0').
\]

The integral of the moments can be written as

\[
\langle \rho_{KQ} \rangle = \sum_{qq'} (-1)^{K+Q+q'}E_{KQ}(e) \times [K]^{1/2}(f_K(q,q';v)) \times \begin{pmatrix} 1 & 1 & K \\ q' & -q & -Q' \end{pmatrix} \times \begin{pmatrix} K & K & 0 \\ Q & 0 & -Q' \\ 0 & -Q' & Q' \end{pmatrix}.
\]

Substituting for the 3-J symbols and for \(E_{KQ}(e)\) yields

\[
\langle \rho_{KQ} \rangle = \sum_{qq'} (-1)^{K+Q+q'+s} \begin{pmatrix} 1 & 1 & K \\ s & -s & 0 \end{pmatrix} \times C_{KQ}(\Theta_0',\Phi_0')(f_K(q,q';v)) \times \begin{pmatrix} 1 & 1 & K \\ q' & -q & -Q' \end{pmatrix},
\]

which can be rewritten as

\[
\langle \rho_{KQ} \rangle = \left( \frac{[K]}{[J_A]} \right)^{1/2} (-1)^{K+s+1} \begin{pmatrix} 1 & 1 & K \\ s & -s & 0 \end{pmatrix} C_{KQ}(\Theta_0',\Phi_0')(A^K).
\]

Here \((A^K)\) represents the total orientation or alignment and can be defined as

\[
(A^K) = \sum_{qq'} (-1)^{q'+s} \begin{pmatrix} 1 & 1 & K \\ q' & -q & -Q' \end{pmatrix} (f_K(q,q';v)).
\]

Note that the 3-J symbol in Eq. (A1) restricts the total alignment to moments \(K \leq 2\) and for linearly polarized photolysis (\(s = 0\)) to even \(K\) only. Now substituting for Eq. (A1) into the expression for the image intensity and contracting the prod-
uct of spherical harmonics using the spherical harmonic addition theorem yield

\[
\langle P \rangle = \Re \left\{ \langle \mathcal{E}^0 \rangle \right\} = \sum_{K} P_K \left( \frac{3[K]}{[L]} \right)^{1/2} \langle -1 \rangle^{K+1} \\
\times \left( \frac{1}{s} \begin{array}{cc} 1 & K \\ -s & 0 \end{array} \right) \Re \left\{ \langle A^K \rangle \right\} P_K (\cos \Theta_{da}),
\]

(A3)

where \( \Theta_{da} \) is the angle between the photolysis and analysis (probe) laser reference axes. This equation agrees with previous work, for example, Dixon, \(^1\) from which it is clearly seen that the intensity depends only on the angle between the pump and probe axes. The expressions for the total polarization parameters \( \langle A^K \rangle \) are

\[
\langle A^0 \rangle = 1,
\]

\[
\langle A^1 \rangle = \sqrt{2} \langle \mathcal{E}^2_1' (1,1;v) \rangle + 2 \Re \left\{ \langle \mathcal{E}^2_1 (1,0;v) \rangle \right\} = -\sqrt{\frac{2}{3}} \sqrt{3} V_1 (J_1) (\sqrt{3} a_1 + \gamma_1),
\]

\[
\langle A^2 \rangle = 2 \times \left( \frac{2}{3} \langle \mathcal{E}^2_2 (1,1;v) \rangle - \langle \mathcal{E}^2_2 (0,0;v) \rangle \right)
\]

\[
+ 2 \sqrt{3} \Re \left\{ \langle \mathcal{E}^2_1 (1,0;v) \rangle \right\} + \sqrt{6} \langle \mathcal{E}^2_1 (1,-1;v) \rangle
\]

\[
= \sqrt{\frac{2}{5}} V_2 (J_2) (\alpha_2 + \gamma_2 + \eta_2).
\]

This last equation is in agreement with Bracker et al., \(^4\) apart from the slightly different normalization employed. In general, note that just two parameters determine the total intensity, \( \langle A^1 \rangle \) and \( \langle A^2 \rangle \). Terms involving the former parameter disappear for linearly polarized photolysis, as well as for all geometries with \( \Theta_{da} = \pi/2 \). An example of the latter is the counterpropagating geometry with linearly polarized photolysis and circularly polarized probe radiation.

**APPENDIX B: MOLECULAR FRAME MOMENTS**

The following describes how to obtain the molecular frame moments from the laboratory frame polarization parameters. The starting point is Eq. (16) of Bracker et al.: \(^4\)

\[
\rho_{KQ}(\Theta, \Phi; v) = \rho_{00}(\Theta, \Phi; v) \sum_{Q'} D_{QQ'}^K(\Phi, 0; 0) \rho_{KQ}^{\text{mol}}(\theta, \phi; v),
\]

which can be inverted by premultiplying through by \( D_{QQ'}^K(\Phi, 0; 0) \) and summing over \( Q \),

\[
\sum_{Q} D_{QQ'}^K(\Phi, 0; 0) \rho_{KQ}(\Theta, \Phi; v) = \rho_{00}(\Theta, \Phi; v) \sum_{Q'} D_{QQ'}^K(\Phi, 0; 0) \rho_{KQ}^{\text{mol}}(\theta, \phi; v),
\]

\[
\times D_{QQ'}^K(\Phi, 0; 0) \rho_{KQ}^{\text{mol}}(\theta, \phi; v).
\]

Note that in the present work we have normalized \( \rho_{00}(\Theta, \Phi; v) \) to unity. Using the sum rules for rotation matrices, this equation reduces to

\[
\rho_{KQ}^{\text{mol}}(\theta, \phi; v) = \sum_{Q'} D_{QQ'}^K(\Phi, 0; 0) \rho_{KQ}(\Theta, \Phi; v),
\]

which can be rewritten as

\[
\rho_{KQ}^{\text{mol}}(\theta, \phi; v) = \sum_{Q'} D_{QQ'}^K(\Phi, 0; 0) \frac{\rho_{KQ}(\Theta, \Phi; v)}{\rho_{00}(\Theta, \Phi; v)}.
\]

(B1)

The next step is to substitute for the laboratory frame moments using Eq. (2), using the sum rules again to contract the product of rotation matrices. We can then write

\[
\rho_{KQ}^{\text{mol}}(\theta, \phi; v) = \frac{3}{4\pi} \left( \frac{[K]}{[J_\Lambda]} \right)^{1/2} \frac{1}{\rho_{00}(\Theta, \Phi; v)}
\]

\[
\times \sum_{k_d, q_d} \sum_{q'} (-1)^{k^* q^*} E_{k_d q_d}(e)
\]

\[
\times \left[ [k_d]^2 f_k^d(q, q'; v) \right]
\]

\[
\times \left( \frac{1}{q' - s} \right) D_{q_d q_d}^d(\Phi, 0). 
\]

(B2)

We now need to turn to the equation for the dissociation light polarization tensor, \( E_{k_d q_d}(e) \). In the molecular frame, the angles \( (\theta, \phi) \) define the direction of the electric vector of the photolysis light, not the direction of the laboratory frame Z axis. To achieve this, we need to set the angles \( (\Theta_{da}, \Phi_{da}) \) appearing in \( E_{k_d q_d}(e) \) equal to zero, such that the laboratory frame Z axis and the direction of the electric vector of the photolysis light are coincident. With this identity, the laboratory angles \( (\Theta, \Phi) \) become coincident with the molecular frame angles \( (\theta, \phi) \).

\[
E_{k_d q_d}(e) = (-1)^{[k_d]^2} \left( \begin{array}{ccc} 1 & 1 & k_d \\ s & -s & 0 \end{array} \right) C_{k_d q_d}(0, 0)
\]

\[
= (-1)^{[k_d]^2} \left( \begin{array}{ccc} 1 & 1 & k_d \\ s & -s & 0 \end{array} \right) \delta_{k_d}. 
\]

On substitution into Eq. (B2) the resulting expression reduces to

\[
\rho_{KQ}^{\text{mol}}(\theta, \phi; v)
\]

\[
= \frac{3}{4\pi} \left( \frac{[K]}{[J_\Lambda]} \right)^{1/2} \frac{1}{\rho_{00}(\Theta, \Phi; v)}
\]

\[
\times \sum_{k_d, q_d} \sum_{q'} (-1)^{k^* q^*} [k_d]^2 f_k^d(q, q'; v) \left( \begin{array}{ccc} 1 & 1 & k_d \\ s & -s & 0 \end{array} \right)
\]

\[
\times \left( \frac{1}{q' - s} \right) C_{k_d q_d}(0, 0). 
\]

(B3)

Note that in this equation \( Q = q' - q \) and that

\[
\rho_{00}(\Theta, \Phi; v) = \frac{W(v)}{4\pi [J_\Lambda]^{1/2}} \left[ 1 + \beta(v) P_2(\cos \Theta) \right],
\]

with
\[ W(v) = \frac{f_0(0,0;\phi) + 2f_0(1,1;\phi)}{f_0(0,0;\phi) + 2f_0(1,1;\phi)} \]

and

\[ \beta(v) = \frac{2[f_0(0,0;\phi) - f_0(1,1;\phi)]}{f_0(0,0;\phi) + 2f_0(1,1;\phi)}. \]

**APPENDIX C: LONG RANGE POTENTIALS**

In this section we take the molecular frame \( z \) axis to lie along the Jacobi coordinate \( R \), with the diatomic molecule lying in the \( xy \) plane. In the following we give explicitly the interaction potential energy matrix \( V_{\text{MM}}^{\Lambda} \) in the long range region. This can be constructed in a basis \(|\Lambda L M\rangle\), where \(|\Lambda\rangle\) is the wave function for diatomic molecule \( \Lambda \) and \(|LM\rangle\) is the electronic orbital wave function for atom \( B \) (here we rewrite \( M_L \) as \( M \)).

The multipole expansion of the Coulomb interaction between the diatomic molecule \( \Lambda \) and the atom \( (B) \) is given by

\[
\hat{V} = \sum_{l_{A,B}} \sum_{m_A,m_B} \sum_{a,b} (-1)^{l_A} \left( \frac{2l_A + 1}{2l_A} \right)^{1/2} \hat{Q}_{l_{A,B}}^{a,b} \hat{Q}_{l_{A,B}}^{a,b} |L_M\rangle \langle L_M|,
\]

where the multipole operators \( \hat{Q}^{a,b} \) are defined in the molecular frame (MF) defined above, and from now on we will drop the superscript MF notation.

The computation of the diatomic part involves a rotation to the diatomic frame (DF),

\[
\langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda' \rangle = \sum_{m'} \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda' \rangle D_{m_{A,m'}|m_{A,m'}}^{l_{A}}(0,\gamma,0),
\]

where \( D_{m_{A,m'}|m_{A,m'}}^{l_{A}}(0,\gamma,0) \) is a Wigner rotation matrix element and \( \gamma \) is the Jacobi angle. \( \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda' \rangle \) is the expectation value of the multipole moment operator of the diatomic molecule \( \Lambda \). The matrix element is zero unless \( \Lambda = m' + \Lambda' \).

Given that \( \Lambda \) and \( \Lambda' \) only take the values \( \pm 2 \), \( m' \) is restricted to 0 or 4. Thus the quadrupole term \( l_A = 2 \) term must be diagonal in \( \Lambda \), and Eq. (C2) reduces to

\[
\langle 2 | \hat{Q}_{l_{A,B}}^{a,b} | 2 \rangle = \langle 2 | \hat{Q}_{l_{A,B}}^{a,b} | 2 \rangle = \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda \rangle D_{m_{A}}^{l_{A}}(0,\gamma,0) = \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda \rangle C_{l_{A}}^{m_{A}}(\gamma,0).
\]

The first equality in this equation arises due to the reflection symmetry of the molecular wave function. Note that the hexadecapole \( l_A = 4 \) term is the first to contain an off-diagonal term \( \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda' \rangle \), but even in this case, it is probably safe to assume that it is smaller than the diagonal \( \langle \Lambda | \hat{Q}_{l_{A,B}}^{a,b} | \Lambda \rangle \) term. In the following we will restrict ourselves to the quadrupole terms of the expansion only, and therefore off-diagonal terms in the potential matrix will be neglected.

To evaluate the atomic part of the matrix elements we use the Wigner-Eckart theorem:

\[
\langle LM | \hat{Q}_{l_{B}}^{b} | L'M' \rangle = (-1)^{l_{B} - M} \left( \begin{array}{cc} L & l_{B} \\ -M & m_{B} \end{array} \right) \langle L | \hat{Q}_{l_{B}}^{b} | L \rangle \langle L | \hat{l}_{B} | L \rangle,
\]

where \( \langle L | \hat{Q}_{l_{B}}^{b} | L \rangle \) is the reduced matrix element of the multipole moment operator for atom \( B \). Coriou et al.\(^{14}\) have recently shown that for \( O(D_{2}) \), only the \( l_B = 2 \) term is nonzero and is negative in sign. Finally, the expression for the potential matrix elements is obtained by converting the Clebsch-Gordan coefficient from Eq. (C1) into a 3-J symbol through

\[
\langle l_A m_{A} | l_{B} m_{B} | (l_A + l_B) \rangle = \delta_{m_{A} - m_{B}} (-1)^{l_{A} - l_{B}}(2l_{A} + 1 + l_{B})^{1/2} \times \left( \begin{array}{ccc} l_A & l_B & l_A + l_B \\ + m_{A} & - m_{A} & 0 \end{array} \right).
\]

Thus, expressing the matrix \( V_{\text{MM}}^{\Lambda} \), in terms of a multipole expansion yields

\[
V_{\text{MM}}^{\Lambda}(\gamma) = \langle \Lambda L M | \hat{V}(\gamma) | \Lambda L M' \rangle = \sum_{l_{A,B}} \sum_{m_A,m_B} \sum_{a,b} \left( \frac{2l_A + 1}{2l_A} \right)^{1/2} \left( \begin{array}{cc} 2l_A + 2l_B + 1 \end{array} \right)^{1/2} \times \left( \begin{array}{ccc} l_A & l_B & l_A + l_B \\ - m_A & m_B & 0 \end{array} \right),
\]

As noted above, only terms diagonal in \( \Lambda \) are considered here, which is strictly only valid for the quadrupole-quadrupole term in the potential. Therefore, in the present application, the sums over \( l_A \) and \( l_B \) disappear and \( l_A = l_B = 2 \). Diagonalization of the matrix \( V \), for a given value of \( \gamma \), yields the eigenvalues (i.e., the long range potential energy surfaces) and eigenfunctions (i.e., the electronic wave functions in the \(|LM\rangle\) basis). Note that in the present case, because \( V \) is diagonal in \( \Lambda \), the problem reduces to solving two identical \( 5 \times 5 \) matrices, yielding five pairs of degenerate \( \Lambda' \) or \( \Lambda'' \) states. Group theory tells us that these states correspond to the \( \Sigma^\pm, \Pi^\pm, \Delta, \Phi, \) and \( \Gamma \) states at linearity. In this model, \( R_c \) is located at sufficiently long range that the degeneracy of these states is not lifted on bending.

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5. R. N. Zare, *Angular Momentum (World Scientific, New York, 1988).*