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Production of $O_2$ Herzberg states in the deep UV photodissociation of ozone


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High-resolution imaging experiments combined with new electronic structure and dynamics calculations strongly indicate that the $O(^3P) + O_2$ products with very low kinetic energy release ($E_k < 0.2$ eV) formed in the deep UV (226 nm) photodissociation of ozone reflect excitation of the Herzberg states of $O_2$: $A^1 \Delta_u (v=0, 1, 2)$ and $A^3 \Sigma_u^+ (v=0, 1)$. This interpretation contradicts the earlier assignment to very high ($v \geq 26$) vibrational states of $O_2 (\Sigma_u^+)$. © 2009 American Institute of Physics. [DOI: 10.1063/1.3157236]

The strong Huggins/Hartley band system of ozone (200 nm < $\lambda$ < 350 nm) is due to excitation of the third state with $A^1 \Sigma_u^+$ symmetry ($1B_2$ in $C_{2v}$), termed $B$ in what follows. The $B$ state correlates diabatically with the singlet channel ($2, O(^1D) + O_2 (\Sigma^1 \Delta)$), with a quantum yield of the order of 90%. All distributions $P(E_k)$ have a broad maximum of around 2 eV corresponding to vibrational states of $v = 10 – 15$ of $O_2 (\Sigma_u^+)$. Classical trajectory surface hopping (TSH) calculations unambiguously attribute this main part to the $B \rightarrow R$ transition.

At all wavelengths, $P(E_k)$ extends to very low energies. For longer wavelengths the shape of the low-energy part varies gradually with $\lambda$. However, at 233 nm, there is a qualitative change: a very sharp peak near $E_k \approx 0$, several times higher than the main peak, suddenly appears out of the broad low-energy part of the distribution. At shorter wavelengths this peak broadens but remains distinct from the slowly varying component. It may be attributed to excitation of either very high vibrational states $v \geq 26$ in channel (1) or to the excited triplet (Herzberg) states (3) $O(^3P_J) + O_2 (A^1 \Delta_u)$ and (3’) $O(^3P_J) + O_2 (A^3 \Sigma_u^+)$, which open at 234.1 nm.

In this communication we present new imaging experiments and new electronic structure and dynamics calculations that provide strong evidence that the low-$E_k$ peak corresponds to the excited triplet channels (3) and (3’) rather than channel (1). The broad low-energy part of $P(E_k)$ found at all wavelengths is not part of this study and remains unexplained.

The experiments were performed using dc slice imaging in an apparatus described previously. Ozone was trapped onto silica gel in a dry ice bath. Argon at 2 bar was flowed over the silica gel into a pulsed valve. The silica gel was held in the dry ice bath during the experiment to produce a very dilute beam. The ozone beam passed through a skimmer and then crossed two laser beams each tuned 10 cm$^{-1}$ on opposite sides of the oxygen atom 2+1 ionization resonance, 44 315 cm$^{-1}$ for $O(^3P_J)$ and 44 237 cm$^{-1}$ for $O(^3P_J)$. The laser beams achieved both single photon dissociation of ozone and two-color reduced Doppler detection of the $O(^3P_J)$ product. The dissociation/ionization took place in a velocity map imaging electrode assembly optimized for high resolution slice imaging (4 electrodes, repeller of $\approx -350$ V). The resulting ions were accelerated down a 1.2 m flight tube onto a 120 mm imaging detector. Higher energy products were allowed to fly out of the detection window because our interest in this study was exclusively the very slow channel. The images were analyzed to give $P(E_k)$ using our in-house suite of programs.

Calibration of $E_k$ was performed using both $O_2$ dissociative ionization at 225 nm (Ref. 15) and OCS photodissociation to give CO($j=65$) at 230 nm. These alternative methods gave results that agreed within 5%. The resolution in the experiment was ultimately limited by the $O^+$ recoil in the resonantly enhanced multiphoton ionization step, which contributes up to 35 m/s uncertainty in the measured O atom velocity.

The measured distributions of kinetic energies for $O(^3P_1)$ and $O(^3P_2)$, shown along with the sliced images in Fig. 1 have two prominent peaks at low $E_k$ and a smaller featureless background extending to higher energies. They agree very well with the earlier observations at the same $\lambda$, although the new data have better energy resolution. The threshold energies for the possible electronic and vibrational states of $O_2$ were calculated with 1.0523 eV for the $O_2 \rightarrow O(^{3}P_2) + O_2 (\Sigma_u^+)$ dissociation energy, 158.3 cm$^{-1}$ for the $O(^3P_2) \rightarrow O(^3P_1)$ splitting, the vibrational energies ($G_v$) of
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The high-energy edges of the faster peaks coincide very close to zero that reflects the bright spot in the center of the sliced image. It may be attributed to $v=2$ and $\Omega=3$ of channel (3) and/or $v=1$ of channel (3'), which are both barely open (0.005 eV). No similar peak appears for $\Omega=3, 2,$ and 1 (in order of increasing energy). The population of the $\Omega$ states depends on both $J$ and $v$. The shift of about 0.01–0.02 eV to lower energies with respect to the threshold energies is in accordance with modest rotational excitation of the $O_2$ product. The observed peaks do not have an obvious relation to the thresholds for production of vibrationally excited $O_2(X^3 \Sigma_g^+)$.

Understanding the path leading to excitation of the Herzberg states requires knowledge on excited electronic states of ozone and the changes in electronic structure from the Franck–Condon (FC) region out to the products. Here we use a global diabatic potential energy surface (PES) for the $B$ state ($V_B$, Fig. 2), determined by multireference configuration interaction (MRCI) calculations (including Davidson correction) with the augmented correlation consistent triple-zeta basis. The lowest five states with $^1A'$ symmetry were calculated and diabatic PESs for four states $X, A, B,$ and $R$ were constructed as described by Qu et al. Five states are sufficient for describing channels (1) and (2). However, more states are required in order to explain the population of the Herzberg channels (3) and (3').

We call the $O$–$O$ bond coordinate that eventually breaks $R_1$, and the bond that remains, the $O_2$ stretching coordinate, $R_2$. Most trajectories that start on $V_B$ near the FC region very quickly dissociate into the singlet channel (2). About one tenth of the trajectories in the vicinity of the $B/R$ crossing seam “hop” to the repulsive $R$-state PES and dissociate to the triplet channel (1), yielding the highly inverted vibrational distribution of $O_2(X^3 \Sigma_g^+)$ seen in all experiments. Calculations describe the main part of this distribution very well. However, the calculated $P(E_v)$ for this pathway at 226 nm is practically zero below 0.8 eV.

Far in the exit channel ($R_1 \geq 4.5a_0$) $V_B$ has the shape of a double-well potential along the $O_2$ stretching coordinate. The deeper well at $R_2=2.3a_0$ belongs to the singlet channel (2) while the comparatively flat well around $2.8–2.9a_0$ results from the excited triplet channels (3)/(3'); a low barrier separates the two wells. To elucidate the electronic structure in this region we performed MRCI calculations as described above, but including the lowest ten $^1A'$ states. Cuts along the $O_2$ bond for two $O$–$O_2$ separations $R_1$ are depicted in Fig. 3. [Potential cuts for additional $R_1$ involving 12 states calcu-
lated at the complete active-space self-consistent field (CASSCF) level are available electronically. Many of the states are nearly degenerate for $R_1=6a_0$ and the potential curves are those of free O$_2$ with O either in the ground or the excited state. The two states with minima around $R_2=2.9a_0$ are the Herzberg states O$_2$(a$^1$A) and O$_2$(3$^3$Σ$_u^+$) in combination with O(1$^1$P). The calculations predict a separation of $\approx$660 cm$^{-1}$ in good agreement with the experimental value of 630 cm$^{-1}$. The lowest Herzberg state $c(2^1$Σ$_u^–)$ correlates, in combination with O(1$^1$P), with the triplet states of ozone and therefore does not appear in Fig. 3. All potentials approach the 3O(3$^3$P) three-body breakup limit for large $R_2$.

The representation in Fig. 3(b) is misleading in that it suggests the possibility of transitions between the states belonging to channel (2) and those representing channels (3)/(3'). However, such transitions are impossible for infinitely separated O and O$_2$ because they would require an accompanying transition in the free atom. The “crossing” exists only because of the $^3P–^1D$ excitation of the oxygen atom of 1.96 eV. The degeneracies are lifted with decreasing $R_1$ and the network of potentials becomes much more complex as illustrated for $R_1=4a_0$ in Fig. 3(a). The real crossings for infinite $R_1$ turn into a multitude of avoided crossings. All states have the same $^1$A' symmetry and nonadiabatic transitions between them are, in principle, allowed.

The double-well behavior of $V_{O_2}$ at large O–O$_2$ separations is due to the combination of channels (2) and (3) and in this sense it is an adiabatic PES. Truly diabatic PESs, which asymptotically correlate with either of the channels (2), (3), or (3'), are needed for accurate dynamics calculations. However, in view of the complexity of the potential curves in Fig. 3(a), especially in the region around $R_2=2.8a_0$, the construction of such diabatic PESs is currently impracticable. We can therefore describe the possible path leading to the Herzberg channels only with words and calculations on the single $V_{O_2}$ surface rather than true multistate dynamics calculations.

A few percent of all trajectories started near the FC point on $V_a$ with energy corresponding to 226 nm lead to products with extremely low translational energies ($\approx$4% with $E_{\mu}<0.5$ eV). All these trajectories have one feature in common: they first perform an almost perfect symmetric stretch motion on the saddle and are driven to large O–O bond lengths (Fig. 2). From the first turning point they are deflected into the dissociation channel and finally perform wide-amplitude vibrational motion almost perpendicular to the dissociation path. The initial part of the trajectory is reminiscent of a classical periodic orbit. The average rotational energy of the trajectories with very small $E_{\mu}$ is of the order of 0.01–0.02 eV.

Because $V_a$ is an adiabatic PES, trajectories with large-amplitude vibration pass through the transition state (TS) region $R_1\approx2.8a_0$ and sample both product channels. In a quantitative multistate TSH calculation, however, the highly excited transient O$_2$ would decide at each crossing of the TS whether it stays in one channel or hops to another. The nonadiabatic couplings diminish with increasing $R_1$, and beyond $R_1\approx5.5a_0$ or so transitions between channels (2) and (3)/(3') become unlikely and the populations are frozen. While O and O$_2$ separate, nonadiabatic transitions to state $R$ are also possible. They could explain part of the nearly constant background in the measured $P(E_{\mu})$ but not the distinct maxima below 0.2 eV. A quantitative evaluation of the populations of channels (3) and (3') is beyond our current capabilities. An estimation of the populations of the spin states $J=0–2$ and $\Omega=1–3$ is even more demanding and requires knowledge on the triplet and quintet states of O$_3$ correlating with the Herzberg states.

The essential step in the proposed dissociation path leading to products with very small $E_{\mu}$ is the large-amplitude symmetric stretching in the very first moments of motion in the B state. It immediately and directly results in highly excited O$_2$ products.
cally modified by collisions with N2, those experiments were not performed with enough time resolution to eliminate the possibility of rapid electronic quenching of the O2 Herzberg states.23 The vibrational distribution with a minimum at \( v = 22 \) seen in the LIF experiments was probably the superposition of a distribution of nascent vibrationally excited O2\((X \: ^3\Sigma_g^-)\) that decreased with increasing \( v \), arising from the main part of channel (1), and a distribution of product states from electronic quenching of the Herzberg states.

Our assignment of the low-energy peaks to production of the Herzberg states suggests two obvious experimental tests. The first is a direct detection of the electronically excited O2 products in either a molecular beam or a static gas photodissociation experiment. The second is the observation of the kinetics of production of very high vibrational states in O2\((X)\) after photodissociation of ozone in a bath of N2 to confirm their production by quenching of the Herzberg states.

We conclude that three processes contribute to the observed distribution of O\((^1P)\) kinetic energies. The main process requires a transition from the ozone \( B \) state to the \( R \) state, yielding ground-state molecular oxygen, and produces the main peak in \( P(E_v) \). The second is dissociation to the O2 Herzberg states \( A \: ^3\Delta_u \) and \( A \: ^3\Sigma_g^+ \) producing the sharp spikes that appear at very low kinetic energies for photolysis wavelengths \( \lambda \leq 233 \) nm. Finally, there is a third process that yields the low, broad distribution of slow O atoms at all wavelengths. That process remains unknown and merits further experimental and theoretical investigation.

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22. See EPAPS supplementary material at http://dx.doi.org/10.1063/1.3157236 for additional potential cuts.