Polarization control in a crossed undulator without a monochromator

Tatsuo Kaneyasu¹ ² 6, Yasumasa Hikosaka¹, Masaki Fujimoto³ ⁴, Hiroshi Iwayama⁴ ⁵ and Masahiro Katoh⁴ ⁵

¹ SAGA Light Source, Tosu 841-0005, Japan
² Institute for Molecular Science, Okazaki 444-8585, Japan
³ Institute of Liberal Arts and Sciences, University of Toyama, Toyama 930-0194, Japan
⁴ Sokendai (The Graduate University for Advanced Studies), Okazaki 444-8585, Japan
⁵ Hiroshima Synchrotron Radiation Center, Hiroshima University, Higashi-Hiroshima 739-0046, Japan
⁶ Author to whom any correspondence should be addressed.
E-mail: kaneyasu@saga-ls.jp

Keywords: polarization control, crossed undulator, photoexcitation, Zeeman quantum beat

Abstract

In this study, it was demonstrated that the polarization control of light from a crossed undulator can be achieved using material response, without any prior monochromatization. Nonmonochromatized horizontally and vertically polarized radiation emitted from two undulators is irradiated to He atoms, and Rydberg states in the extreme ultraviolet region are excited resonantly. The polarization of the photoexcitation is adjustable in a controlled manner by tuning the electron beam orbit between the two undulators, which is verified by observing Zeeman quantum beats in the fluorescence decays of the He Rydberg states.

1. Introduction

The use of light polarization properties is particularly effective for studying electronic and magnetic states of matter. For example, x-ray magnetic circular dichroism spectroscopy [1] is important in studying the magnetic property of materials. The crossed undulator scheme for synchrotron radiation is a powerful method to produce light beams of various polarization states. The operation principle of the crossed undulator is based on the interference of horizontally and vertically polarized radiation fields emitted from two undulators in tandem [2]. Various polarization states, such as circular polarization and tilted linear polarization, can be obtained by tuning the phase difference between radiation fields using a phase shifter magnet that controls the electron orbit between the undulators.

Since the first proposal by Kim in 1984 [2], the crossed undulator has been adopted to achieve polarization control in synchrotron light sources [3–5] and free electron lasers [6–10]. In addition to spatially homogeneous light beams, a vector beam in which the polarization axis changes depending on the azimuthal angle in the beam cross section [11] can be produced by the interference of the radiation fields from a crossed undulator setup. Additionally, interference can be applied [12–14] in verifying the vortex properties of harmonic radiation from helical undulators [15].

Light polarization adjustment from crossed undulators requires light monochromatization. This is because the relativistic electrons passing through the two undulators emit time-separated radiation fields (radiation wave packets), and a temporal stretching of the radiation wave packets is required to achieve an interference between them [2]. Herein, we present a new concept to obtain variable polarization in the crossed undulator scheme without using a monochromator. Our proposal is based on the resonant photoabsorption of atoms exposed to nonmonochromatized radiation from a crossed undulator.
2. Polarization control

2.1. Basic concept
We first consider radiation from a single relativistic electron, for describing the basic concept of the polarization control. Figure 1 illustrates the polarization control of light from a crossed undulator by a monochromator and by the present method. The crossed undulator comprises two adjacent $N$-period undulators with a phase shifter magnet between them. A relativistic electron that passes through the crossed undulator emits a pair of horizontally and vertically polarized radiation wave packets that are characterized by an $N$-cycle oscillating field and a bandwidth of $\omega/\Delta\omega \sim N$. The two radiation wave packets are separated by a time delay $\tau$, and do not originally overlap in the time domain. In a conventional system, the radiation wave packets are temporally stretched by passing through a monochromator with a resolving power of $\omega/\Delta\omega \gg N$, which allows an interference between the two radiation fields. As the monochromator transmits a specific frequency component of the incoming radiation fields, the frequency component of the total light field of the undulator radiation can be expressed by

$$\tilde{E}_T(\tau, \omega) = \tilde{E}(\omega) \left\{ \epsilon_x e^{i\omega \tau} + \epsilon_y \right\}, \quad (1)$$

where $\epsilon_x$ and $\epsilon_y$ are polarization vectors, and $\tilde{E}(\omega)$ is the Fourier component of one of the two wave packets. Equation (1) shows that the polarization is determined by the interference of the two radiation fields, which can be controlled by tuning $\tau$. Consequently, the polarization of light after the monochromator can be controlled by adjusting $\tau$ using the phase shifter magnet.

When an atom is directly irradiated by nonmonochromatized radiation from the crossed undulator, it absorbs specific frequency components according to its discrete energy levels. The resonant width of the individual excited state is much smaller than the bandwidth of the undulator radiation. Therefore, the atomic resonant photoabsorption is expected to reflect the polarization determined by the interference of the two radiation fields at the resonant frequency. Hence, the polarization response of the atom can be interpreted as the result of polarization control at the resonant frequency. This phenomenon has been well-described in terms of the spectral interpretation of Ramsey fringes [16, 17] observed in the sequential interaction between the atom and the time-separated double pulse. While this time-domain description has been developed in laser science [18], it has been proven recently that the same description can be applied to the interaction of radiation wave packets from undulators [19, 20]. It is noteworthy that, to control the polarization in this scheme, the excited state has a lifetime that is longer than the temporal duration of the pair of radiation wave packets.

2.2. Radiation from a bunch of electrons
The basic concept described in the previous section is valid even for radiation from a bunch of relativistic electrons, as long as each atom is considered to interact with one of the identical wave packet pairs which are randomly populated in a radiation pulse. Using this simple picture, one can understand the main feature of the coherent control of atoms using the crossed undulator setup [19, 20]. However, an atom
exposed to radiation from a bunch of electrons can successively interact with numerous radiation wave packet pairs, and the successive interaction has to be considered for a more correct description.

On the basis of the time-domain description of atomic excitation [18], we discuss the successive interaction of a He atom with numerous radiation wave packet pairs. The light–atom interaction is simplified in a one-dimensional model as in figure 2. The wave packet pairs are randomly distributed in the radiation pulse. The length of the pulse is typically around 100 ps in synchrotron light sources, corresponding to the electron bunch length in a storage ring. Here, the waveform shapes and the time delay \( \tau \) are common for all wave packet pairs included in the pulse. As a result of the interaction with a single wave packet pair among them, the He atom is excited into the 1\( snp \) coherent superposition state and its wave function can be written as

\[
|\psi(t)\rangle \approx e^{-i\omega_0 t} \left\{ e^{i\omega_\tau} |x\rangle + |y\rangle \right\},
\]

where

\[
a = \tilde{E}(\omega_n) \mu.
\]

Here \( \omega_0 \) is the transition frequency, \( \mu \) is the dipole matrix element associated with the transitions into the \( |x\rangle \) and \( |y\rangle \) states characterized by the 1\( snp \) configurations, respectively.

We consider a target state

\[
|\psi_0\rangle = e^{i\omega_0 \tau_0} |x\rangle + |y\rangle
\]

which can be prepared by light with a certain polarization \( \varepsilon_x e^{i\omega_0 \tau_0} + \varepsilon_y \). The population of the target state after the interaction is expressed by

\[
n_0 = |\langle \psi_0 | \psi(t) \rangle|^2 = 2|a|^2 \left( 1 + \cos \omega_\mu \Delta \tau \right)
\]

where \( \Delta \tau = \tau - \tau_0 \). The population oscillates by varying the time delay with a time period corresponding to the transition frequency. This oscillation is due to Ramsey interference in the atomic excitation by time-separated radiation fields. On the other hand, this interference can be interpreted as the spectral modulation of radiation

\[
n_0 = 2|\mu|^2 S(\omega_n)
\]

where

\[
S(\omega) = \left| \left\{ \varepsilon_x e^{-i\omega_0 \tau_0} + \varepsilon_y \right\} \cdot \tilde{E}_T(\tau, \omega) \right|^2 = |\tilde{E}(\omega)|^2 \left( 1 + \cos \omega \Delta \tau \right),
\]

represents the spectral intensity of the total light field for a certain polarization \( \varepsilon_x e^{i\omega_0 \tau_0} + \varepsilon_y \). Therefore, the atomic resonance can be understood in view of the polarization of the interacting total light field which is determined by the time delay.

The interaction of the He atom and a single wave packet pair occurs in series. The wave function of the 1\( snp \) superposition state, resulting from the successive interaction with numerous radiation wave packet pairs, can be written as the sum of the excitation by each wave packet pair:

\[
|\psi(t)\rangle \approx \sum_j e^{-i\omega_0 \tau_j} a \left\{ e^{i\omega_\tau} |x\rangle + |y\rangle \right\} e^{i\omega_n T_j},
\]

where \( T_j \) is the time delay of the \( j \)th wave packet pair with respect to the first one. Using the wave function in (7), the population of target state is expressed by

\[
n_0 = \left| \langle \psi_0 | \sum_j e^{-i\omega_0 \tau_j} a \left\{ e^{i\omega_\tau} |x\rangle + |y\rangle \right\} e^{i\omega_n T_j} \rangle \right|^2 = 2|a|^2 \left( 1 + \cos \omega_\mu \Delta \tau \right) \sum_j e^{i\omega_n(T_j-T_0)} \\
\cong 2M|a|^2 \left( 1 + \cos \omega_\mu \Delta \tau \right),
\]

where \( M \) is the number of wave packet pairs among them.
where $M$ is the number of wave packet pairs, i.e., the number of electrons in the bunch. The population also shows the oscillation corresponding to Ramsey interference, when the time delay varies as in the single electron case. The cross terms related to different electrons in the bunch are cancelled out in (8) because the wave packet pairs are randomly distributed within the pulse. The only difference between the two cases is the maximum transition probability which is proportional to the intensity of radiation. Therefore we conclude that the atomic resonant photoabsorption reflects the polarization determined by the interference between the radiation fields in the case of radiation from a bunch of electrons. Although we have neglected the decay of the excited state in the above discussion, this conclusion holds even for an excited state with a finite lifetime. In such a case, the number of interacting wave packets may be insufficient to cancel out the cross terms related to different electrons. However we always observe an ensemble average of events for different atoms over many radiation pulses. The ensemble-averaged population of the finite-lifetime state is simply given by an expression similar to equation (8) with the corresponding decay term.

2.3. Experiment

A proof-of-concept experiment was performed at the BL1U undulator beamline of the 750 MeV UVSOR-III storage ring [21]. The UVSOR-III storage ring was operated in a single bunch mode to measure the fluorescence decay of a Rydberg state in He. A single bunch operation provides light pulses with an interval of 178 ns and a width of 300 ps (FWHM). The beam current was around 10 mA during the experiment. Therefore approximately $10^{10}$ wave packet pairs were randomly distributed within the 300 ps light pulse. The light source of the BL1U beamline comprises a twin APPLE-II-type variable polarization undulator installed in tandem. The number of magnetic periods and the period length of the undulators are 10 and 88 mm, respectively. The upstream and downstream undulators were tuned to produce the radiation wave packets of vertical and horizontal linear polarizations, respectively. The phase shifter magnet between the undulators controls the orbit length of the electrons with 0.5 nm accuracy, which corresponds to a time delay tuning with 2 as accuracy. The accuracy in controlling the time delay is limited by the resolution of the power supply system for the phase shifter magnet.

The peak photon energy of the fundamental radiation of the undulators was adjusted to the excitation energy (24.1 eV) of the 1s6p 1P state. The lifetime of the 1s6p 1P state is 13.1 ns [22], which was much longer than the few femtoseconds duration of the radiation wave packet pair. The energy spread of the undulator radiation was approximately 10%; therefore, several 1snp states were excited simultaneously. The energy width of each excited state was in the order of $\omega/\Delta\omega$ to 10$^6$, which was dominated by the Doppler spread of gas atoms at room temperature.

The polarization property of the light that resulted in the photoabsorption in He can be investigated by observing the Zeeman quantum beat in the fluorescence decay of the 1s6p 1P excited state [23]. Figure 3 shows the schematics of the experimental setup for the Zeeman quantum beat measurement. The spatially central part of the undulator radiation was selected using a 0.4 mm diameter pinhole. To eliminate visible light emitted from the undulators and bending magnets, a 75 nm-thick Al filter was inserted in front of the interaction region, to which He atoms were introduced in the form of an effusive beam. It is noteworthy that, unlike typical synchrotron radiation experiments, the undulator radiation was not monochromatized. A pair of solenoid coils (not shown in figure 3) was attached to the interaction region, and a magnetic field of approximately 20 Gauss was applied parallel to the quantization axis (z-axis).
Figure 4. Zeeman quantum beat measured for nonmonochromatized radiation from the crossed undulator. The phase difference $\Delta \phi$ between the horizontally and vertically polarized radiation fields was tuned from $\pi$ to $9\pi/2$ with a step of $\pi/2$. The gray circles represent measured data and the red curves calculated results. The polarization ellipses and vectors assumed in the calculation are shown.

Fluorescence photons with a wavelength of 345 nm, which were emitted by the decay from the 1s6p to 1s2s states, were selected and detected using a photomultiplier tube equipped with a bandpass filter and a linear polarizer. Linearly polarized photons were detected, of which the polarization axis was parallel to the $x$-axis. Although the 1s6p state could be populated via cascade decays from the highly excited 1snp ($n > 8$) states, the cascade contribution was negligible. For instance, the branching ratio was estimated to be $1 \times 10^{-4}$ for a cascade decay from the 1s8p state to the 1s6p state, according to the transition probabilities reported in [24].

3. Results and discussion

Figure 4 shows the fluorescence decay curves from the He 1s6p state, measured at eight different settings of $\tau$. The corresponding phase differences between the horizontally and vertically polarized fields are $2m\pi + \pi$ to $2m\pi + 9\pi/2$, with a step of $\pi/2$. The value of $m$ should be 13, considering the minimum delay time (approximately 2100 as) resulting from the slippage in the undulator section and the electron travel in the drift space. For simplicity, we hereinafter refer to modulo $2m\pi$ as phase difference $\Delta \phi$.

The beat structures shown in figure 3 change according to $\Delta \phi$, as follows. While a clear beat structure is observed at $\Delta \phi = \pi$, it almost disappears at $\Delta \phi = 3\pi/2$. Although a beat structure appears again at $\Delta \phi = 2\pi$, the beat structure is in antiphase compared with that at $\Delta \phi = \pi$. After the beat structure almost vanishes at $\Delta \phi = 5\pi/2$, it appears again at $3\pi$. The same trend in beat structure evolution is observed from $\Delta \phi = 3\pi$ to $\Delta \phi = 9\pi/2$.

Figure 5 shows the scheme for the observed Zeeman quantum beat with an arbitrary elliptical polarization. A magnetic field applied in the interaction region split the energies of the magnetic sublevels in the 1s6p state, and photoexcitation resulted in a coherent superposition of the $M_j = \pm 1$ magnetic sublevels. As a result of the quantum interference between the two decay paths, a beat structure of frequency $\Delta \omega$ was superimposed on the exponential decay curve of the fluorescence. Assuming an elliptical polarization tilted by angle $\chi$ with respect to the horizontal axis, the fluorescence yield is expressed as

$$I(t) \propto e^{-\gamma t} \left\{ 1 + P_L \cos(\Delta \omega t + 2\chi) \right\},$$

where $P_L$ is the degree of linear polarization along the tilt angle. Equation (9) shows that the phase and visibility of the quantum beat are directly related to the polarization (for details, see appendix A).

The quantum beat structure for the polarization expected at each phase difference in figure 4 was calculated using equation (9). In the calculation, the degree of polarization was assumed to be 0.50. The amplitudes of the horizontally and vertically polarized light fields were set as 1:0.7, as determined by their intensity measurements. This intensity difference can be explained by the different distances between the source positions and the pinhole for two undulators. Therefore, the linear polarization (expected to be produced at $\Delta \phi = n\pi + \pi/2$) should be tilted with $\pm 35^\circ$, and the elliptical polarization (expected for $\Delta \phi = n\pi$) should exhibit an ellipticity of 1.4. The obtained curves were convoluted with the time resolution.
Figure 5. Scheme of the Zeeman quantum beat measurement with the fluorescence decay of the 1s6p excited state in He. The excited state is expressed as a coherent superposition of the magnetic sublevels of $M_j=±1$, which results in the beat structure of frequency $\Delta \omega$ in the fluorescence decay curve.

($\sigma = 1.6 \text{ ns}$) of the measurement system and compared with the observations presented in figure 4. The observed beat structures were reproduced by the calculated curves. This result verified that polarization control in the crossed undulator could be achieved using atomic response. Note that the small discrepancy between the measurement and calculation was most likely due to the deviation in the phase tuning.

The degree of polarization was assumed to be 0.50 in the calculation. The instrumental effect due to a finite detection angle was removed from the value (see appendix A), whereby the actual degrees of linear and circular polarizations were estimated to be 0.71 and 0.67, respectively. The estimated degree of polarization was considerably low compared with the light source limit of 0.94, which was determined by the electron beam properties, such as the emittance and energy spread [2]. This result could be due to two reasons. The first one is the misalignment of the pinhole. The degree of polarization decreases when the pinhole position deviates from the central part of the undulator radiation. The other is the misalignment of the radiation from the two undulators, which reduces the visibility of the quantum beat. The nonoverlapping components of the horizontally and vertically polarized radiation generate quantum beats that overlap incoherently in opposite phases.

4. Conclusions and perspectives

We have demonstrated a new concept to employ variable polarizations in a crossed undulator. The proposed method is based on the resonant photoabsorption of atoms irradiated by nonmonochromatized radiation from a crossed undulator. While in this paper we showed the achievement of polarization control in the extreme ultraviolet range, this method is available in other wavelength ranges. We here discuss the advantages of this method in the soft x-ray regime where polarization control by undulators is particularly useful in various experiments. First, this method enables substantial increase of the photon flux at the sample position. In general, soft x-ray monochromatization with several mirrors and a grating reduces the photon flux by an order of magnitude or more. Second, this concept is applicable for any combination of polarization states. For example, the tilt angle of the linear polarization vector is controlled by the time delay between a cross-circularly polarized radiation from undulators [20]. In addition, the ability to combine arbitrarily polarized radiation enables the control of polarization in structured light such as vector beam, using pairs of cross-circularly polarized vortex beams from a helical undulator [11].

In particular, studies in the field of the atomic and molecular physics may benefit from these advantages. Atomic interaction with structured light has been attracting much attention recently in this field, and this method enables one to perform the investigations more simply. On the other hand, this method may be applicable to generate free-electron wave packets with a tailored momentum distribution [25, 26]. An interesting challenge is to study the coherence between the free-electron wave packets emitted from a molecular inner-shell, where different behavior against the atomic sites may be exhibited. Furthermore, applying this method to the photoelectron circular dichroism in chiral molecules [27, 28] is challenging and interesting from a fundamental point of view.
New J. Phys. 22 (2020) 083062

Figure A1. Zeeman quantum beat in the fluorescence decay curves measured for (a) horizontal and (b) vertical linear polarization. The gray circles represent measured data and the red curves show the calculated results. The polarization vectors assumed in the calculation are shown. A 1.0 mm diameter pinhole was used in this measurement.

Acknowledgments

This study was partly supported by KAKENHI Grants-in-Aid (Nos. 17H01075, 18K03489 and 18K11945) from the Japan Society for the Promotion of Science. The construction of BL1U at UVSOR was supported by the Quantum Beam Technology Program of the Ministry of Education, Culture, Sports, Science, and Technology and the Japan Science and Technology Agency.

Appendix A. Zeeman quantum beat with an arbitrary elliptical polarization

The total light field from the crossed undulator at frequency $\omega$ and time delay $\tau$ is described by equation (1). This light field can also be described by a coherent sum of the circularly polarized radiation fields.

$$\tilde{E}_T(\tau, \omega) = |\tilde{E}_R| e^{+i\chi} \varepsilon_R + |\tilde{E}_L| e^{-i\chi} \varepsilon_L,$$

(A.1)

where $\tilde{E}_{R,L}$ are Fourier components, $\varepsilon_{R,L}$ are polarization vectors, and the phase $\chi$ corresponds to the tilt angle of the polarization ellipse. In this case, the time evolution of the excited superposition state is described by

$$|\psi(t)\rangle = c_+ e^{-i\omega_+ t} e^{-\chi t} |+1\rangle + c_- e^{-i\omega_- t} e^{-\chi t} |-1\rangle,$$

(A.2)

where

$$c_{\pm} = \frac{\pm 1}{|\tilde{E}_R|} e^{i\chi} \mu_\pm.$$

Here, $\omega_{\pm}$ are transition frequencies of the $|\pm 1\rangle$ states with respect to the ground state, $\gamma$ is the decay constant, and $\mu_\pm$ are dipole matrix elements associated with $\Delta M_j = \pm 1$ transitions. The coefficients $c_{\pm}$ of the superposition state are related to the polarization of the excitation light; therefore, the fluorescence yield reflects the polarization. The fluorescence yield decays from the $|\psi(t)\rangle$ state into the $|1s2s\rangle$ state and is expressed by

$$I(t) \propto |\langle 1s2s | \psi(t)\rangle|^2$$

$$\propto e^{-\gamma t} \left\{ 1 + 2 |\tilde{E}_L| |\tilde{E}_R| \cos (\Delta \omega t + 2\chi) \right\}$$

$$= e^{-\gamma t} \left\{ 1 + P_L \cos (\Delta \omega t + 2\chi) \right\}.$$

(A.3)

Here, $P_L$ is the degree of linear polarization along the tilt angle, and $\mu$ is the dipole operator. Equation (A.3) shows that the phase of the quantum beat is directly related to the tilt angle of the polarization ellipse, and that the visibility of the beat structure is dependent on the degree of linear polarization.

Figure A1 shows the fluorescence decay curves measured for the horizontally and vertically linear polarization using one of the two undulators. The decay curves are modulated by the quantum beat, and the phase of the beat structure is dependent on the polarization direction. The beat starts from the top of the oscillation in the case of the horizontally linear polarization, while it starts from the bottom for the vertically linear polarization, indicating that information regarding the polarization of the excitation light can be derived from the quantum beat. The experimental results are well reproduced by the theoretical curves calculated using equation (A.3). The theoretical curves were convoluted with the time resolution ($\sigma = 1.6$ ns) of the measurement system. To reproduce the measured beat structure, the degree of linear polarization was assumed to be 0.70, which was much lower than that typically expected value for linearly polarized undulator radiation ($>0.99$). This disagreement was caused by the finite detection angle of the fluorescence detector, which blurred the beat structure. Therefore, we consider the effective polarization degree of 0.70 as an instrumental effect for the fully polarized light.
ORCID iDs

Tatsuo Kaneyasu https://orcid.org/0000-0002-8052-0422
Yasumasa Hikosaka https://orcid.org/0000-0001-5898-4426
Hiroshi Iwayama https://orcid.org/0000-0002-5992-5281
Masahiro Katoh https://orcid.org/0000-0001-9149-1500

References

[1] Nakamura T and Suzuki M 2013 J. Phys. Soc. Jpn. 82 021006
[2] Kim K J 1984 Nucl. Instrum. Methods 219 425
[3] Bahrdt J, Gaupp A, Gudat W, Mast M, Molter K, Peatman W B, Scheer M, Schroeter T and Wang C 1992 Rev. Sci. Instrum. 63 339
[4] Yamamoto S et al 2017 J. Synchrotron Radiation 24 352
[5] Chung T Y, Yang C S, Chu Y L, Lin F Y, Jan J C and Hwang C S 2017 Nucl. Instrum. Methods A 850 72
[6] Kim K J 2000 Nucl. Instrum. Methods A 445 329
[7] Wu Y K, Vinokurov N A, Mikhailov S, Li J and Popov V 2006 Phys. Rev. Lett. 96 224801
[8] Ding Y and Huang Z 2008 Phys. Rev. Spec. Top. Accel. Beams 11 030702
[9] Deng H et al 2014 Phys. Rev. Spec. Top. Accel. Beams 17 020704
[10] Ferrari E et al 2015 Sci. Rep. 5 13531
[11] Matsuba S, Kawase K, Miyamoto A, Sasaki S, Fujimoto M, Konomi T, Yamamoto N, Hosaka M and Katoh M 2018 Appl. Phys. Lett. 113 021106
[12] Bahrdt J, Holldack K, Kuske P, Müller B, Scheer M and Schmid P 2013 Phys. Rev. Lett. 111 034801
[13] Katoh M et al 2017 Sci. Rep. 7 61303
[14] Kaneyasu T, Hikosaka Y, Fujimoto M, Iwayama H, Hosaka M, Shigemasa E and Katoh M 2017 J. Synchrotron Radiat. 24 934
[15] Sasaki S and McNulty I 2008 Phys. Rev. Lett. 100 124801
[16] Noordam L D, Duncan D I and Gallagher T F 1992 Phys. Rev. A 45 4734
[17] Christian J F and Broers B 1995 Phys. Rev. A 52 3655
[18] Kullière C 2005 Femtosecond Laser Pulses (Berlin: Springer)
[19] Hikosaka Y, Kaneyasu T, Fujimoto M, Iwayama H and Katoh M 2019 Nat. Commun. 10 4988
[20] Kaneyasu T, Hikosaka Y, Fujimoto M, Iwayama H and Katoh M 2019 Phys. Rev. Lett. 123 233401
[21] Adachi M, Zen H, Konomi T, Yamazaki K, Hayashi K and Katoh M 2013 J. Phys.: Conf. Series 425 042013
[22] Theodouri C E 1984 Phys. Rev. A 30 2910
[23] Hikosaka Y, Iwayama H and Kaneyasu T 2020 J. Synchrotron Radiat. 27 675
[24] NIST Atomic Spectra Database 2019 version 5.7 available at http://physics.nist.gov/asd3
[25] Hockett P, Wollenhaupt M and Baumert T 2015 J. Phys. B: At. Mol. Opt. Phys. 48 214004
[26] Wollenhaupt M, Lux C, Krug M and Baumert T 2013 ChemPhysChem 14 1341
[27] Hergenhahn U, Rennie E E, Kugeler O, Marburger S, Lischke T, Powis I and Garcia G 2004 J. Chem. Phys. 120 4553
[28] Powis I, Harding C J, Garcia G A and Nahon L 2008 PhysChemPhys 9 475