Photoionization of atmospheric gases studied by time-resolved terahertz spectroscopy

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Abstract

We investigate the laser induced ionization and plasma formation in pure oxygen and nitrogen by means of optical pump–terahertz probe spectroscopy. Focused amplified femtosecond pulses at 405 and 810 nm are used to ionize the gas molecules by nonlinear processes. The ionized gas is probed by picosecond terahertz pulses in the frequency range of $0.2 - 2\,\mathrm{THz}$ to obtain the free electron density ($10^{13} - 10^{17}\,\mathrm{cm}^{-3}$) and the electron scattering time (0.2 ps for N_2 and 0.4 ps for O_2). We demonstrate the importance of the centrifugal barrier for the photoionization process using circularly polarized light.

1 Introduction

Ionization of atoms or molecules at high laser fields has be subject of extensive research [1, 2, 3]. Femtosecond laser time-resolved studies of the gas plasma employ intense ultrashort laser pulses to ionize the examined gas [4, 5, 6]. In this regime the laser electric field amplitude in the range of $0.5-5\times10^8\,\mathrm{V/cm}$ can be easily achieved (corresponding to tens or a few hundreds of $\mathrm{TW/cm^2}$ peak pulse intensity). The forces these fields exert on the electrons are comparable to the binding forces of valence electrons in molecules leading to a number of interesting phenomena which cannot be simply described within the time-dependent perturbation theory.

Beside its contribution to laser plasmas generated in air,[7] ionized gases (oxygen in particular) play an prominent role in the upper atmosphere. For example O_2^+ is an important atmospheric source of singlet oxygen atoms [8, 9] which in turn are responsible for auroral radiation [10].

THz pulsed radiation has been shown to provide a powerful tool for probing the plasma in ionized gases [11, 12, 13]. Kampfrath et al. have been able to probe the dynamics of ionized O_2

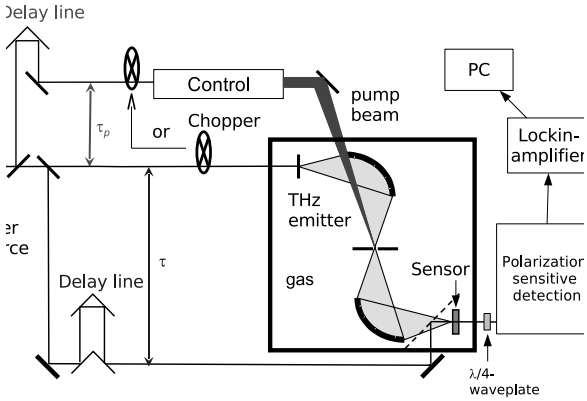

Figure 1: Experimental scheme. The pump beam branch contains a specific arrangement for a careful control of the pump beam attenuation and polarization and a second harmonic stage. The chopper can be either in the probe branch measuring the THz field transmitted through the sample, or in the pump branch to measured the photoinduced change in the transmitted THz signal.

and Ar gas by THz pulses in the frequency range 5-30 THz; from the evolution of the plasma properties the temperature of the free electron gas has been determined.[14] These successful application of THz spectroscopy for plasma characterization are related to the fact that (i) the THz pulses are highly sensitive to the presence of mobile charge carriers and that (ii) optical pump—THz probe spectroscopy is able to provide time-resolved complex THz dielectric spectra. As a result it is possible to determine both the concentration and the momentum scattering rate of the charges in the plasma.

In this paper we apply the optical pump — THz probe spectroscopy to the study of multi-photon ionization of atmospheric gases. Our experimental scheme is outlined in Sec. 2. Sec. 3 describes a model of interaction of the THz radiation with photoionized gas plasma and presents a method of extraction of the plasma characteristics from experimental data. The experimental results presented in Sec. 4 are then interpreted in Sec. 5.

2 Experimental details

We employed an experimental setup typical for time-resolved THz studies [15] which is schematically shown in Fig. 1. As a laser source, we used a Quatronix Odin Ti:sapphire multipass femtosecond pulse amplifier. It generates 1 mJ laser pulses with a 55 fs temporal length, 810 nm mean wavelength and 1 kHz repetition rate. The output laser beam is divided into three parts. Up to 80% of the laser power is used to the ionization of the examined gas (oxygen or nitrogen) either directly or using the second harmonic generation (for 405 nm pump wavelength) in an LBO crystal. We used a setup which enables controlling the pump pulse polarization and fine tuning of its peak intensity [12]. About 20% of the laser power was used for the generation and gated detection of THz pulses. For this purpose we use a pair of 1 mm thick (110)-oriented ZnTe crystals [16]. The entire THz setup was enclosed in a box, which was evacuated, and then refilled with the examined gas at atmospheric pressure. The THz radiation was focused with a pair of ellipsoidal mirrors and a 0.5 mm aperture was placed into the common focal point of the optical pump and THz probe beams. This alignment allowed us to quantify the volume and shape of the probed ionized gas column (see [12] for details).

A synchronous detection scheme locked to a mechanical chopper operating at 166 Hz was

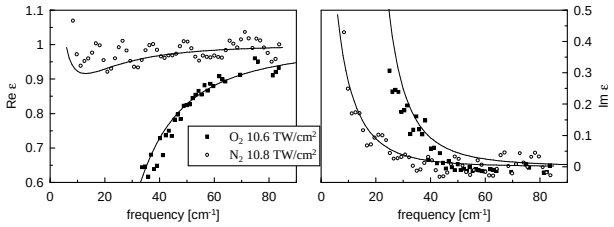

Figure 2: Dielectric spectra of oxygen and nitrogen ionized at comparable pump intensity, 405 nm and linear polarization. Symbols: experimental data; lines: fit by Eq. (4).

used. Two delay lines controlled the time delay between the THz probe and optical sampling pulses (the sampling time τ makes connection to the real time of the experiment) and between the optical pump and THz probe pulses (pump-probe delay τ_p). For the measured data we use the following notation. $E(\tau, \tau_p)$ is the THz wave form transmitted through the ionized gas at time τ_p after the photoionization event. $E(\tau)$ is the reference wave form, i.e., the THz wave form transmitted through the gas with pump beam off. $\Delta E(\tau, \tau_p) = E(\tau, \tau_p) - E(\tau)$ is the transient THz wave form; this quantity can be measured directly by placing the chopper into the pump branch. It reflects the changes in the gas induced by photoionization.

3 Interaction with probing radiation

The THz probing radiation strongly interacts with mobile delocalized carriers. Consequently, the coupling of the THz electric field to free electron plasma will dominate the interaction of the probing radiation with photo-ionized gas. This enables one to learn from transient THz dielectric spectra the free electron concentration and scattering rate at various levels of the pump intensity. In contrast, the ionized gas molecules are much heavier and their translational motion cannot be efficiently driven by the THz field.

To examine the dielectric response of the plasma, we measure the transient THz wave form $\Delta E\,(\tau, \tau_p)$ (i.e. τ is scanned while τ_p is set to $5 - 10$ ps after photoionization). From that, the transient transmission function $\Delta t(\omega)$ of the plasma can be calculated using Fourier transformation:

$$\Delta t(\omega) = \Delta E(\omega; \tau_p)/E(\omega)\,. \tag{1}$$

where $\Delta E(\omega; \tau_p)$ resp. $E(\omega)$ are $\Delta E(\tau; \tau_p)$ resp. $E(\tau)$ transformed into the frequency domain. It is then possible to evaluate the refractive index of the plasma from the equation (see Ref. [12] for a detailed description):

$$\Delta t(\omega) = \{\exp[i\omega(n + i\kappa)d_{\text{eff}}/c] - 1\}\frac{V}{d_{\text{eff}} S_{\text{a}}}\,, \tag{2}$$

where S_a is the opening of the aperture at the focus of the pump and THz beams, d_{eff} is the effective interaction length of the two beams (the effective length of the plasma column as "seen" by the THz probing pulse), V is the interaction volume (the volume of the probed part of the plasma cylinder). $\Delta N = n + i\kappa$ is the pump-induced change in the refractive index. Eq. 2 assumes that the transmission losses due to inhomogeneities (i.e., partial reflections and scattering) can be neglected compared to those related to the propagation effects described by the exponential term. Finally, we get the complex permittivity of the plasma from the equation:

$$\varepsilon(\omega) = (1 + \Delta N)^2 \tag{3}$$

To fit these spectra, we use a Drude-Lorentz model [12]:

$$\varepsilon(\omega) = 1 - \frac{\omega_p^2}{i\omega/\tau_s + \omega^2 - \omega_p^2/2} \tag{4}$$

where ω_p is the plasma frequency, which is connected with the free electron density N_e: $\omega_p^2 = N_e e^2/(\varepsilon_0 m_e)$. Here e and m_e are the electron charge and mass, respectively, and ε_0 stands for the vacuum permittivity. The term $\omega_p^2/2$ accounts for a restoring force which originates in a screening due to the spatial separation of positive ions and negative electrons [12]. This model allows us to determine the free electron density and their scattering rate for the experiments we carried out.

4 Experimental results

To examine the nature of the ionization processes we varied several parameters of the pump beam. We carried out the pump-probe experiments for two ionization wavelengths: 810 nm (the output wavelength of our laser source) and its second harmonics, i.e., 405 nm. Comparative measurements were performed with linearly and circularly polarized pump beams, where the polarization was modified using a $\lambda/4$-plate. Finally, we carefully varied the intensity of the pump beam. The peak excitation intensity was calculated based on the measured pump power, the time length of the pulse and the size of the laser spot at the point of excitation. The spot size was measured by a CCD camera in the regime where the pump intensity was strongly attenuated. We estimate that the absolute value of the pump pulse peak intensity could be determined with an uncertainty factor better than 2 (i.e. twice smaller or twice larger than the calculated value). The random error within a series of measurements was much smaller and did not exceed 10%. The applied peak pump intensities are in the range $1 - 12$ TW/cm^2 for 405 nm and $10 - 300$ TW/cm^2 for 810 nm pump wavelengths.

In Fig. 2 we plot examples of transient dielectric spectra for the linearly polarized pump beam with a wavelength of 405 nm. We show here a comparison of the spectra obtained for the two gases at about the same pump intensity. During these measurement the pump-probe delay was fixed to a value such that the THz pulse probed the plasma ~ 5 ps after the photoionization.

We measured first the pump intensity dependence of the peak values of the THz transient wave forms (i.e. $\Delta E_{max} = \Delta E(\tau, \tau_p)$, where τ was set to the maximum of the signal wave form). These measurements provide information only about the average THz response of the plasma. Nevertheless, they are fast and the temporal stability of the laser is sufficient to allow us to monitor precisely the highly nonlinear dependence of the signal versus the pump intensity. The values obtained during these measurements were then used to normalize the peaks of the entire transient THz wave forms which were measured subsequently. These latter measurements are slower but allow us to evaluate the plasma characteristics encoded into the THz signal transmitted through the photoexcited region.

The whole set of measurements allowed us to determine the electron density and scattering time versus the pump intensity [12]. The calculated electron densities obtained for the linearly polarized pump beam at 405 nm are displayed in Fig. 3. The scattering time of the free electrons was found to be (200 ± 50) fs for nitrogen and (400 ± 50) fs for oxygen. In our earlier paper [12] we reported a small pump intensity dependence of the scattering time in oxygen (150–350 fs): this was caused by a numerical error in evaluation of our data which was fixed in the present paper. Note that the values and variation of the electron plasma density presented and discussed in Ref. [12] are not influenced by this. The electron scattering time we find

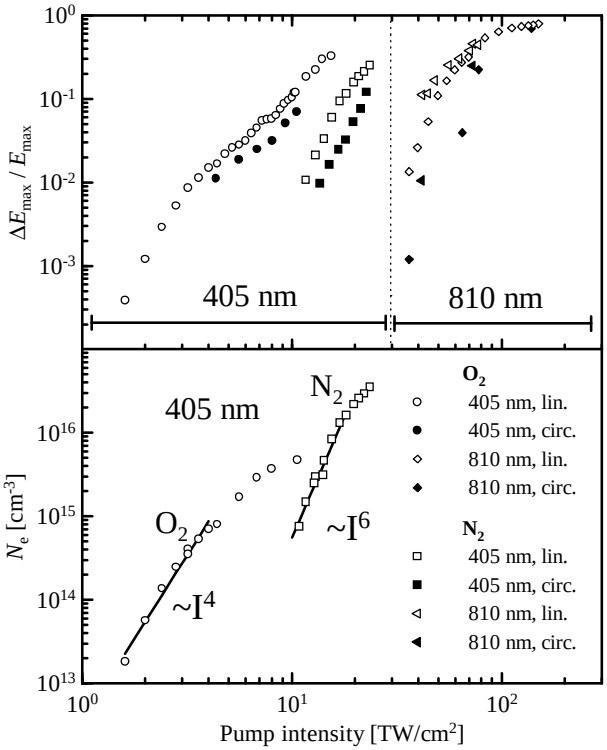

Figure 3: a) Peak transient THz signal versus the pump intensity for excitation wavelengths 405 and 810 nm for linear and circular polarization (raw data). b) Density of free electrons versus the pump intensity for 405 nm pump wavelength and linear polarization. The low-intensity part of these dependences can be fitted by simple power laws $\propto I^4$ for O_2 and $\propto I^6$ for N_2 indicating 4-photon and 6-photon processes.

is clearly sub-picosecond and it is independent of the pump intensity for both gases in our experimental conditions.

The 810 nm pump induces appreciable variation in the transmitted THz wave form only for intensities higher than about 30 TW/cm^2 as can be observed in the inset of Fig. 3. Qualitatively, the nonlinearity of these ionization processes is clearly higher than for 405 nm pump wavelength. However, for 810 nm excitation we were not able to evaluate quantitatively the plasma density. This is because the interaction volume V for such highly nonlinear processes is rather small and the approximations, under which Eq. 2 was derived, namely that of negligible Fresnel losses, is not fulfilled.

Finally, we carried out comparative measurements of the peaks of the THz signal wave form for linear (ΔE_{lin}) and circular (ΔE_{circ}) polarization of the pump beam. We measured both ΔE_{lin} and ΔE_{circ} with variable pump intensity at wavelengths of 405 and 810 nm. The ratio of the measured values $\Delta E_{circ}/\Delta E_{lin}$ for oxygen and 405 nm pump beam is 0.7 at 4 TW/cm^2 and slowly decreases to about 0.6 at 12 TW/cm^2. For nitrogen these values are found in the range of $\sim 0.4 - 0.5$ at 405 nm [Fig. 3(a)]. In contrast, at 810 nm the ionization rate with the circularly polarized beam is at relatively small pump intensities (where the probing signal does not reach saturated regime) for both gases strongly reduced compared to the experiments with linearly polarized beam; $\Delta E_{circ}/\Delta E_{lin} \approx 0.1$.

5 Discussion

The adiabatic ionization potential for oxygen is $E_{\mathrm{B}} = 12.07$ eV. For the 405 nm excitation wavelength (photon energy of 3.06 eV), at least 4 photons are required to ionize the O_2 molecule.

The ionization process can be described by the strong-field approximation proposed by Reiss [17] which consists in neglecting of the influence of the binding potential on the detached electron in comparison with the oscillating high-frequency laser field effects. In this state the electron acquires a "quiver" or ponderomotive energy given by [3]:

$$U_{\mathrm{P}} = e^2 E_L^2/(4 m_e \omega_L)\,, \tag{5}$$

where E_L is the intensity and ω_L the frequency of the laser field. It is generally accepted that multiphoton ionization channels may open and close at some specific values of the pump intensity $I_L \propto E_L^2$ [3]. This is related to the fact that the electrons, when ejected into a strong laser field, need to overcome the sum of the ionization and ponderomotive potentials. Following (5) this value increases with increasing laser intensity. Consequently, the n-photon ionization channel closes when U_{P} reaches such value that [18]:

$$n\hbar\omega_L = E_{\mathrm{B}} + U_{\mathrm{P}}\,. \tag{6}$$

It means that at lower laser intensities than those required to fulfill the condition (6) the n-photon process represents the lowest order ionization path. In the dipole approximation the relevant transition matrix element is approximately proportional to $\propto \langle \Psi_{\mathrm{f}} | r^n | \Psi_{\mathrm{i}} \rangle$ [19], where Ψ_{i} is the initial (bound) state wave function and Ψ_{f} is the final (delocalized) state one. This term usually dominates in the multiphoton ionization process. At higher intensities, the energy of n photons becomes too low and an additional photon is required to ionize the molecule; $(n+1)$-photon process becomes dominant [18]. In our case Eqs. (5) and (6) yield a value of $I_L \approx 11$ TW/cm^2 for closing the 4-photon channel. This seems to be in a good agreement with the saturation of the free electron density progressively as this irradiance value is approached (Fig. 3).

The ionization potential of nitrogen is 15.60 eV [20]. It means that $5 \times 3.06 = 15.3$ eV is slightly below the threshold. It follows, in agreement with our experimental results that the 6-photon ionization process largely dominates over the whole range of pump pulse intensities. We observe in Fig. 3 that an efficient generation of electrons starts now at higher pump intensities than for O_2 due to the higher order process involved.

The polarization of the ionizing radiation can have a dramatic effect on the nature of the ionization process. As has been shown in [19], excitation with a circularly polarized beam drives the detached electrons into free states $\Psi_{\mathrm{f}}(l)$ with a high angular momentum quantum number l creating thus a centrifugal barrier for ionization. The measure of the barrier is the overlap between the functions $r^n\Psi_{\mathrm{i}}$ and $\Psi_{\mathrm{f}}(l)$. It has been shown by photoelectron spectroscopy in xenon gas [19] that in this situation the near threshold ionization transitions can be inhibited and higher-order processes turn on and appear as higher energy peaks in the photoelectron spectrum. We think that the situation is quite analogous in our experiments, however our experiments do not give direct access to the photoelectron energy. For 405 nm pump with linear polarization the n-th order process (where $n = 4$ for O_2 and $n = 6$ for N_2) dominates while a significant contribution of the $(n + 1)$-th order process comes into play for the experiments with the circular polarization. Higher order processes involved for the 810 nm pump imply a higher centrifugal barrier and, consequently, a larger difference between the experiments with circular and linear polarizations as observed.

The scattering time of the free electrons was found to be independent on the pump intensity at 405 nm. It is 200 fs for nitrogen and 400 fs for oxygen. These numbers are compatible with those obtained for O_2 by Kampfrath and al. [14]. However, in that paper the authors worked with significantly higher plasma densities where the scattering of electrons on oxygen ions plays an important role. In our experiments the scattering of free electrons on neutral molecules dominates. As the scattering cross section of the free electrons is about the same in oxygen and nitrogen ($\sigma \sim 10^{-19}\,\mathrm{m}^2$ [21, 22]), the lower electron scattering time in nitrogen can be explained only by the higher velocity of the generated free electrons. Indeed, for the 6-photon process in nitrogen, the energy of the photons exceeds the ionization threshold by 2.8 eV, whereas for oxygen the sum of energy of the 4 photons is almost equal to the ionization potential of an O_2 molecule. The excess energy of 2.8 eV corresponds to the electron velocity of $v = 1 \times 10^6\,\mathrm{m/s}$. Taking into account that the momentum scattering time τ of free electrons is equal to

$$\tau = \frac{1}{\sigma N_0 v}\,, \tag{7}$$

where $N_0 = 2.7 \times 10^{19}\,\mathrm{cm}^{-3}$ is the concentration of gas molecules at atmospheric pressure, the value of the electron scattering time in nitrogen can be estimated: $\tau \approx 350\,\mathrm{fs}$. This is in a good semiquantitative agreement with the experiment (given the uncertainty in the value σ for example). In this approximation, the twice higher scattering time observed for oxygen implies a four times lower excess energy of electrons (0.7 eV). This value overestimates the excess energy for the four photon ionization process in O_2 suggesting the existence of a minority concentration of free electrons created by a 5-photon process with higher excess energy and lower scattering time. An effective intermediate value is then experimentally observed.

6 Summary

We have applied optical pump–THz probe spectroscopy to examine laser-induced ionization in oxygen and nitrogen at atmospheric pressure. Free electron densities in the range $10^{13} - 10^{17}\ \mathrm{cm}^{-3}$ were observed for the pump pulse peak intensities in the range $3 - 15\ \mathrm{TW/cm}^2$ at

405 nm. The dielectric properties of the generated plasma were determined and interpreted in the framework of the multiphoton ionization process.

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