Interaction of high-power laser pulses with atomic media. I. Spontaneous fields in multiphoton ionization

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A theoretical study of static electromagnetic field generation due to laser radiation force in multiphoton ionization of atomic media is presented. The static fields are spontaneously induced in atomic media and have electric field components in the radial and axial directions of the incident laser beam and an azimuthal magnetic field component. Multiphoton ionization process, which is inevitably apt to occur in the interaction, significantly enhances the spontaneous fields. Numerical calculations of the fields generated in atomic sodium vapor and helium gas have been made, assuming an incident picosecond Nd-doped yttrium aluminum garnet laser pulse. The results show the characteristic natures and magnitude of the spontaneous fields.

I. INTRODUCTION

With the recent development of high power lasers, multiphoton ionization process in atomic media[1] has received increasing attention, because this process is inevitably apt to occur whenever strong radiation interacts with matter. Also, multiphoton ionization process plays an important role, directly or indirectly, in applications of high power lasers such as laser-induced fusion, gas breakdown and isotope separation. Although the laser-induced multiphoton ionization has been extensively studied in recent years, the subjects have been concerned mainly with probabilities for the ionization of atoms and molecules and with the related problems.[1] On the other hand, a few experimental studies[2-5] have been reported so far in which high-energy electrons were found to be produced in laser-induced multiphoton ionization of atomic media. A simple theory of laser radiation force which describes field-gradient force or "ponderomotive force" has been employed to explain the fast electrons observed.[4,5] However, in order to express the radiation force acting on a multiphoton-ionized medium, it may be often crucial to take into account spatial and temporal natures peculiar to the laser-induced multiphoton ionization process.

In this paper, we wish to study static electromagnetic fields spontaneously generated by laser radiation force in multiphoton ionization process in atomic media. The spontaneous fields may exert some influences on atoms and ions in the medium during the interaction. Section II is concerned with the laser radiation force working on a dispersive atomic medium, where the plane-wave approximation for the incident laser beam is employed, and multiphoton ionization process is incorporated in the dielectric constant for the medium. Propagation of a laser pulse in the dispersive medium is also considered to see the nonsteady characteristics and anisotropy of the radiation force. In Sec. III, we derive the form for the spontaneous fields. Calculations of the fields have been made on atomic sodium vapor and on helium gas, assuming an incident pulse of Nd-doped yttrium aluminum garnet (Nd:YAG) laser. The results are presented in Sec. IV.

II. RADIATION FORCE

A. Basic equation

Consider an isotropic and nonabsorptive atomic vapor or gas with dielectric constant $\epsilon(\omega)$. The atomic medium is irradiated with a short (~psec) laser pulse with certain spatial and temporal intensity distributions. The laser frequency ω is assumed to be far from any kind of resonances of the atomic transitions.

We start with the momentum conservation equation per unit volume of the system which consists of the radiation field and the medium. This equation leads to the force $\vec{f}$ working on a unit volume of the medium. In a continuous medium, which we assume here, the force density $\vec{f}$ may be given by[6,7]

$$\vec{f} = \nabla \cdot \overleftrightarrow{\sigma} - \partial \vec{g}/\partial t\,, \tag{1}$$

where $\overleftrightarrow{\sigma}$ is the Maxwell stress tensor associated with the high-frequency electromagnetic fields $\vec{E}$ and $\vec{B}$, which involves radiation pressure acting on the medium, and $\vec{g} = (\vec{E}\times\vec{B})/4\pi c$ is the field momentum density. The stress tensor $\overleftrightarrow{\sigma}$ in a nonmagnetic medium has the form[6,7]

$$\overleftrightarrow{\sigma} = \left\{p_0 - \frac{E^2}{8\pi}\left[\epsilon - \rho\left(\frac{\partial\epsilon}{\partial\rho}\right)_T\right] - \frac{B^2}{8\pi}\right\}\overleftrightarrow{I} + \frac{1}{4\pi}(\epsilon\vec{E}\vec{E} + \vec{B}\vec{B})\,, \tag{2}$$

where $p_0 = p_0(\rho, T)$ is the medium pressure which

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is a function of density ρ and temperature T in the absence of the incident field, and $\overleftrightarrow{I}$ is the unit tensor. The stress tensor given by Eq. (2) contains pressure contribution from the high-frequency motion of bound and free electrons oscillating with the incident field, which can be characterized by the dielectric constant ϵ for the medium. For atomic media, we have a relation[6]

$$\rho\left(\frac{\partial\epsilon}{\partial\rho}\right)_T=\epsilon-1\,. \tag{3}$$

It is clear that, in a steady state, the force $\vec{f}$ averaged over an oscillation period of the incident field must be zero, if no spatial field gradient exists. However, in general, the field distribution of a laser pulse is a function of position and time, so that the medium receives certain time-averaged force from the radiation field. Obviously, we have to consider two time scales for the quantities of interest:[6-9] One is a fast time scale determined by the laser frequency, and the other is a slow time scale on which the temporal history of a laser pulse changes. Since we are interested in the radiation force on the slow time scale, it is convenient to take an average of the quantities over a cycle of field oscillation. Taking into account Eq. (3), the cycle-averaged form of Eq. (1) which gives the radiation force is

$$\langle\vec{f}\rangle=\nabla\cdot\left(-\frac{1}{8\pi}\langle E^2+B^2\rangle\overleftrightarrow{I}+\frac{1}{4\pi}\langle\epsilon\vec{E}\vec{E}+\vec{B}\vec{B}\rangle\right)-\frac{\partial\langle\vec{g}\rangle}{\partial t}, \tag{4}$$

where the brackets $\langle\,\rangle$ represent "cycle averaged". In Eq. (4), we have neglected the pressure term p_0, since we are concerned with phenomena of short duration in an initially uniform atomic medium.

B. Plane-wave approximation

Figure 1 illustrates the definition of polarization (x) and propagation (z) directions of an incident laser pulse. We assume that the laser beam has a spatial intensity distribution depending on

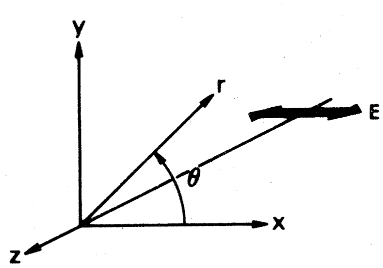

FIG. 1. The coordinates used for the incident laser beam.

only the radial position r and can be approximated by a plane wave. For example, this is the case for an ordinary laser pulse with lowest-order (TEM_{00}) Gaussian intensity distribution of field (see Appendix). Although a focused Gaussian beam differs appreciably from a plane wave, this approximation is still valid in the propagating region $|z|\lesssim b$ of the focused beam, where b is the confocal parameter of the laser beam focused at $z=0$. In this limit, we have $B^2=\epsilon E^2$ and the laser intensity $I=nc\langle E^2\rangle/4\pi$, where $n=\sqrt{\epsilon}$ is the refractive index of the medium. Using these relations and $\langle g\rangle=I/c^2$, Eq. (4) reduces to

$$\langle\vec{f}\rangle=\nabla\cdot\left(\frac{I}{2nc}\right)[(\epsilon-1)\overleftrightarrow{I}-2\epsilon\vec{k}\vec{k}]-\vec{k}\,\frac{1}{c^2}\,\frac{\partial I}{\partial t}, \tag{5}$$

where $\vec{k}$ is a unit vector along the propagation (z) direction. The same form for radiation pressure or force as Eq. (5) has been discussed in detail by Stamper[8] in his study of magnetic field generation in laser-produced plasma. It should be noted, however, that Eqs. (4) and/or (5) give a general form for the laser radiation force working on a dielectric medium such as neutral atomic vapors and fully ionized plasmas.

C. Multiphoton ionization and dielectric constant

The dielectric constant for an atomic medium in which multiphoton ionization proceeds[10] may be a function of position and time, depending on the spatial and temporal distributions of field intensity and particle density. Let us consider a laser pulse passing through a monatomic gas with local number density N at time t. The multiphoton ionization proceeds at a rate given by

$$\frac{dN}{dt}=-WN\,. \tag{6}$$

Here W is the multiphoton ionization probability defined by

$$W=\Gamma I^K(\omega)\,, \tag{7}$$

where Γ denotes the ionization cross section for a specific kind of atoms at the laser frequency ω, and K is the minimum number of laser photons required to ionize a single neutral atom. During the short interaction time concerned, we can ignore the secondary processes subsequent to the multiphoton ionization, such as avalanche ionization due to electron collisions and recombination of electrons and ions. The field depletion owing to the multiphoton absorption is negligibly small. Then, from Eq. (6), we have

$$N=N_0\exp(-F)\,, \tag{8}$$

where

$$F=\int_{-\infty}^{t} W dt\,, \tag{9}$$

and N_0 is the initial density of neutral atoms. Since the electron number density N_e created is given by

$$N_e=N_0-N\,, \tag{10}$$

the function F denotes a measure of the ionization rate N_e/N_0, i.e.,

$$N_e/N_0=1-\exp(-F)\,. \tag{11}$$

The atomic medium in which multiphoton ionization proceeds can be modeled as a mixture of neutral atoms and a cold electron gas to derive the dielectric constant $\epsilon(\omega)$. From the usual frequency-dependent susceptibilities for an atomic gas[11] and for a plasma,[12] the dielectric constant ϵ for the mixture may be expressed as

$$\epsilon(\omega)=1+\delta(\omega)\,, \tag{12a}$$

where

$$\delta(\omega)=\frac{4\pi e^2}{m}N\sum_i \frac{f_i}{\omega_i^2-\omega^2}-\frac{4\pi e^2}{m\omega^2}N_e\,. \tag{12b}$$

In Eq. (12), m and e are the mass and charge of an electron, respectively, and f_i is the oscillator strength for the atomic transition i with excitation frequency ω_i. The first and second terms on the right-hand side of Eq. (12b) represent dispersions due to atomic bound electrons and free electrons, respectively.[13] Equation (12) indicates that the dielectric constant of the medium is a function of position and time through the particle densities N and N_e.

The dispersion term $\delta(\omega)$ given by Eq. (12b) is usually much smaller than unity, i.e., $\delta \ll 1$, under the conditions of interest, but this term is essentially responsible for the radiation force, as expected from Eq. (5). In some cases, depending on the atomic species and the laser frequency ω employed, an additional dispersion term due to the ionic bound electrons should be incorporated in $\delta(\omega)$. The ionic dispersion is usually much smaller compared with the atomic one. In the following treatment, therefore, we will disregard the ionic term for simplicity.

D. Nonsteady force and anisotropy

In order to study the nonsteady contribution and resulting anisotropy of the radiation force, the propagation of a laser pulse in the dielectric medium will be considered here. We start with the derivation of propagation velocity v of a laser pulse, which is the velocity of field momentum or energy transported through the medium. The propagation velocity is generally defined by[6,11]

$$v=\langle S\rangle/\langle U\rangle\,, \tag{13}$$

where $\langle\vec{S}\rangle$ and $\langle U\rangle$ are the cycle-averaged Poynting vector and electromagnetic energy density in the medium, respectively.[14] For the nonmagnetic medium concerned, a general form for $\langle U\rangle$ may be given by[6]

$$\langle U\rangle=\frac{1}{8\pi}\left(\frac{d(\omega\epsilon)}{d\omega}\langle E^2\rangle+\langle B^2\rangle\right). \tag{14}$$

Substitution of Eq. (12) and $B^2=\epsilon E^2$ into Eq. (14) leads to

$$\langle U\rangle=\frac{\langle E^2\rangle}{4\pi}[1+\gamma(\omega)]\,, \tag{15a}$$

where

$$\gamma(\omega)=\frac{4\pi e^2}{m}N\sum_i \frac{f_i\omega_i^2}{(\omega_i^2-\omega^2)^2}\,. \tag{15b}$$

Comparing Eq. (15) with the pure field energy $\langle E^2\rangle/4\pi$ in vacuum, Eq. (15) contains the additional energy that is stored in atoms by virtue of their high-frequency polarization. Since $\langle S\rangle=I=nc\langle E^2\rangle/4\pi$ and $n=\sqrt{\epsilon}$, we have the final form of v:

$$v=\frac{c\sqrt{\epsilon(\omega)}}{1+\gamma(\omega)}\,. \tag{16}$$

In the fixed laboratory coordinates, the temporal history of a propagating laser pulse may be equivalently expressed as a function of the position z by making use of the relation

$$z=-vt\,. \tag{17}$$

Using Eqs. (16) and (17), we rewrite Eq. (5) and get

$$\langle\vec{f}\rangle=\nabla\cdot\left(\frac{I}{2nc}\right)[(\epsilon-1)\vec{I}-2\gamma\vec{k}\vec{k}]\,, \tag{18}$$

where we have neglected the higher-order terms for δ and γ such as δ^2, γ^2, and $\delta\gamma$ since $\delta,\gamma\ll 1$. The anisotropy of radiation force $\langle\vec{f}\rangle$ and its origin are clearly shown in Eq. (18). The anisotropy arises from only the atomic dispersion or the energy stored in atoms. As will be described in Sec. III, the anisotropy makes a solenoidal contribution to the spontaneous electric field and accounts for the magnetic field generation.

It is instructive to check on the energy conservation of the system concerned. The time derivative of the energy $\langle U\rangle$, given by Eq. (15), leads to

$$\frac{\partial\langle U\rangle}{\partial t}=-\nabla\cdot\langle\vec{S}\rangle\,, \tag{19}$$

with the help of Eqs. (16) and (17). Equation (19) represents the energy conservation in the medium, which can be shown as follows: When we consider a small volume V of the system, the conservation of energy flow may be given by[11]

$$\int_s \langle \vec{S} \rangle \cdot d\vec{s} + \int_V (\text{dissipation})\, dV = -\int_V \frac{\partial \langle U \rangle}{\partial t} dV. \quad (20)$$

The terms on the left of Eq. (20) are the rates of energy loss in the volume V by transfer through its surface s and by dissipation, while the term on the right gives the total rate of loss of the energy stored within V. For the nondissipative medium of interest, Eq. (20) reduces to Eq. (19).

III. SPONTANEOUS FIELDS

A static electric field due to the cycle-averaged radiation force may be derived as follows.[15] The radiation force comes from the high-frequency motion of atomic and free electrons in an incident inhomogeneous field. Therefore, it may be possible to consider that the force acts initially on the electrons in the medium concerned. Due to their small mass of electrons, the cycle-averaged force makes a static displacement of electrons with respect to the relatively immovable ions. The force is transmitted to the ions by a self-consistent static electric field $\langle \vec{E}_s \rangle$ which may be given by the relation[8,15,16]

$$\langle \vec{f} \rangle - eN_0 \langle \vec{E}_s \rangle = 0, \quad (21)$$

where N_0 is the effective number density of the static dipoles induced, i.e., $N_0 = N_e + N$. The static force density on the ion species is approximately given by $N_0 e \langle \vec{E}_s \rangle$, and then, the force density working on the entire medium concerned is $\langle \vec{f} \rangle$. Note that in this situation the quasineutrality of charge in the medium is maintained.

From Eqs. (18) and (21), we have the form for the each component of $\langle \vec{E}_s \rangle$ in the cylindrical coordinates,

$$\langle E_s \rangle_r = \frac{1}{2eN_0 c} \frac{\partial}{\partial r} (\epsilon - 1) I, \quad (22a)$$

$$\langle E_s \rangle_\theta = 0, \quad (22b)$$

$$\langle E_s \rangle_z = \frac{1}{2eN_0 c} \frac{\partial}{\partial z} (\epsilon - 1 - 2\gamma) I$$

$$= \frac{1}{2eN_0 c^2} \frac{\partial}{\partial t} (1 - \epsilon + 2\gamma) I. \quad (22c)$$

Here Eqs. (16) and (17) have been used to obtain $\langle E_s \rangle_z$ in the time derivative. From $\nabla \times \langle \vec{E}_s \rangle$, we have components of the spontaneous magnetic field $\langle \vec{B}_s \rangle$ of the form

$$\langle B_s \rangle_r = \langle B_s \rangle_z = 0, \quad (23a)$$

$$\langle B_s \rangle_\theta = \frac{1}{eN_0 c} \frac{\partial}{\partial r} (\gamma I), \quad (23b)$$

taking into account $I = 0$ at $t = -\infty$. As suggested in Sec. II, only the energy stored in atoms or the anisotropic part in the radiation force is related to the magnetic field generation.

To obtain the explicit forms for $\langle \vec{E}_s \rangle$ and $\langle \vec{B}_s \rangle$, we may separate the spatial and temporal distribution functions of the incident laser intensity I and assume the form of

$$I(\omega) = I_0(\omega) \psi(r) \phi(t), \quad (24)$$

where I_0 is the maximum field intensity, and ψ and ϕ are the dimensionless shape functions normalized with respect to I_0. With the help of Eqs. (7), (8), (9), (12), and (15b), substitutions of Eq. (24) into Eqs. (22) and (23) give the explicit forms for the nonzero components of $\langle \vec{E}_s \rangle$ and $\langle \vec{B}_s \rangle$:

$$\langle E_s \rangle_r = \frac{2\pi e}{mc} \left((1 - KF) \exp(-F) \sum_i \frac{f_i}{\omega_i^2 - \omega^2} - [1 - (1 - KF) \exp(-F)] \frac{1}{\omega^2} \right) \frac{\partial I}{\partial r}, \quad (25a)$$

$$\langle E_s \rangle_z = \frac{2\pi e}{mc^2} \left[\left(\frac{\partial I}{\partial t} - WI \right) \exp(-F) \sum_i \frac{f_i (\omega_i^2 + \omega^2)}{(\omega_i^2 - \omega^2)^2} + \left([1 - \exp(-F)] \frac{\partial I}{\partial t} + WI \exp(-F) \right) \frac{1}{\omega^2} \right], \quad (25b)$$

$$\langle B_s \rangle_\theta = \frac{4\pi e}{mc} (1 - KF) \exp(-F) \left(\sum_i \frac{f_i \omega_i^2}{(\omega_i^2 - \omega^2)^2} \right) \frac{\partial I}{\partial r}. \quad (25c)$$

It is easy to see from Eq. (25) that the spontaneous fields are independent of the particle density, and that $\langle E_s \rangle_r$ and $\langle B_s \rangle_\theta$ vanish at the beam center for a Gaussian beam since $\partial I / \partial r = 0$ at $r = 0$. By setting $F = W = 0$ in Eq. (25), we have the forms for $\langle \vec{E}_s \rangle$ and $\langle \vec{B}_s \rangle$ in a pure atomic medium without charged species:

$$\langle E_s \rangle_r = \frac{2\pi e}{mc} \left(\sum_i \frac{f_i}{\omega_i^2 - \omega^2} \right) \frac{\partial I}{\partial r}, \quad (26a)$$

$$\langle E_s \rangle_z = \frac{2\pi e}{mc^2} \left(\sum_i \frac{f_i (\omega_i^2 + \omega^2)}{(\omega_i^2 - \omega^2)^2} \right) \frac{\partial I}{\partial t}, \quad (26b)$$

$$\langle B_s \rangle_\theta = \frac{4\pi e}{mc} \left(\sum_i \frac{f_i \omega_i^2}{(\omega_i^2 - \omega^2)^2} \right) \frac{\partial I}{\partial r}. \quad (26c)$$

For $F, W \to \infty$, we have the forms for the fields in a fully ionized plasma:

$$\langle E_s \rangle_r = -\left(\frac{2\pi e}{mc\omega^2} \right) \left(\frac{\partial I}{\partial r} \right), \quad (27a)$$

$$\langle E_s \rangle_z = \left(\frac{2\pi e}{mc^2\omega^2}\right)\left(\frac{\partial I}{\partial t}\right), \tag{27b}$$

$$\langle B_s \rangle_\theta = 0. \tag{27c}$$

Equations (26) and (27) are equivalent to the forms derived from ponderomotive force in a pure atomic medium[17-20] and in a fully ionized plasma[4,8,15,21], respectively.

IV. CALCULATIONS AND DISCUSSION

To have insight into the characteristic natures and magnitude of the spontaneous fields, numerical calculations have been made for a Nd : YAG laser pulse irradiating atomic sodium (Na) vapor and helium (He) gas. The incident laser pulse is assumed to have the Gaussian spatial and temporal distributions of the field intensity (see Appendix) and the pulse duration $\tau = 28$ psec.

A. Sodium vapor

For Na vapor, we take into account only the resonance $3s$-$3p$ transition in the atomic dispersion, since the contributions from the other atomic transitions are negligibly small at the Nd : YAG laser frequency ($\omega = 1.77\times 10^{15}$ sec^{-1}). If N and N_e are comparable in the ionized Na vapor, the atomic and plasma dispersion terms in Eq. (12b) are in the same order of magnitude and have the opposite sign each other. In the calculations, the conditions and constants employed are as follows: The multiphoton ionization probability $W = 6.8\times 10^{-49} I^K$ sec^{-1}, which has been given experimentally by the present authors,[10] where I is in W/cm^2 and $K = 5$; the spot size $r_0 = 152$ μm; the incident peak intensity $I_0 = 8\times 10^{11}$ and 3×10^{12} W/cm^2.

In Fig. 2, we show the temporal changes in the ionization rate N_e/N_0. At the high input intensity [Fig. 2(b)], the Na vapor in the central region of the laser beam is fully ionized in the early stage of the incident pulse, and a steep gradient of N_e/N_0 is created in the radial distribution. Figure 3 represents the radial electric field $\langle E_s \rangle_r$ generated in the interaction region under the same conditions as in Fig. 2, and demonstrates the following. In the region in which N_e/N_0 is small, $\langle E_s \rangle_r$ arises principally from the atomic dispersion and is negative since $\omega_i > \omega$ and $\nabla I < 0$ in Eq. (26a). With the increasing value of N_e/N_0, $\langle E_s \rangle_r$ changes the sign and has the positive value. The maximum of $\langle E_s \rangle_r$ in the radial distribution appears in the region where N_e/N_0 has the steepest gradient. This implies that a large radial electric field can be produced by a high-order multiphoton ionization process or by a tight focus of the incident beam which will generate a steep gradient of N_e/N_0.

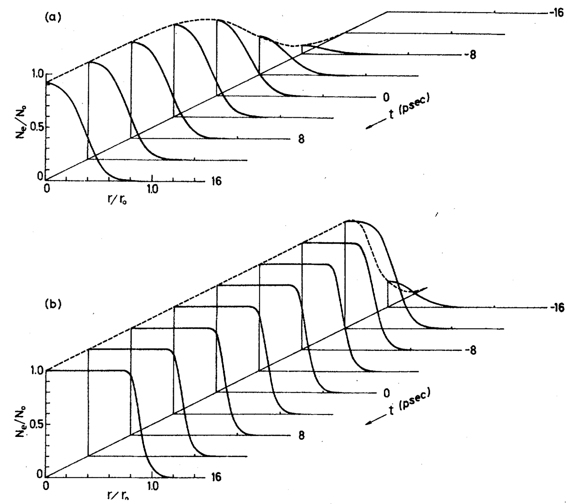

FIG. 2. Temporal changes in the radial dependence of multiphoton ionization rate in atomic Na vapor. The incident Nd:YAG laser pulse is assumed to have the duration $\tau = 28$ psec. (a) Laser intensity $I_0 = 8\times 10^{11}$ W/cm^2; (b) $I_0 = 3\times 10^{12}$ W/cm^2.

Figures 4 and 5 display the axial electric field $\langle E_s \rangle_z$ and the azimuthal magnetic field $\langle B_s \rangle_\theta$, respectively, under the same conditions as in Figs. 2 and 3. Although $\langle E_s \rangle_z$ is very weak compared with $\langle E_s \rangle_r$, it is noted that $\langle E_s \rangle_z$ can have the maximum value at the beam center ($r = 0$) where $\langle E_s \rangle_r$ always vanishes. The magnitude of $\langle B_s \rangle_\theta$ is also very small in the present case, but it may be possible to enhance $\langle B_s \rangle_\theta$ under certain resonance conditions that increase the energy stored in atoms.

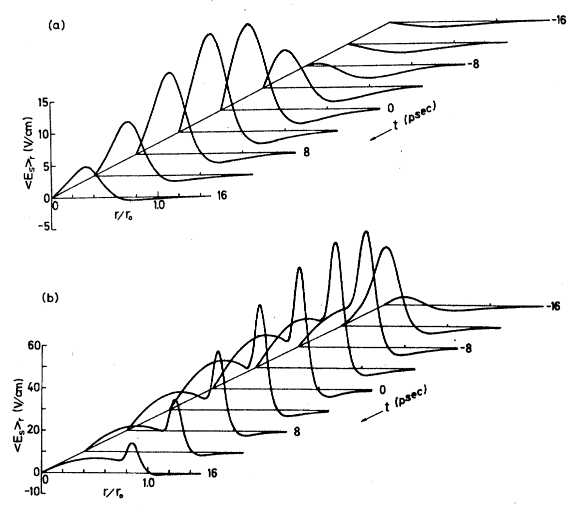

FIG. 3. Temporal changes in the radial distribution of radial electric field generated in atomic Na vapor. The incident Nd:YAG laser pulse is assumed to have the duration $\tau = 28$ psec and the spot size $r_0 = 152$ μm. (a) $I_0 = 8\times 10^{11}$ W/cm^2; (b) $I_0 = 3\times 10^{12}$ W/cm^2.

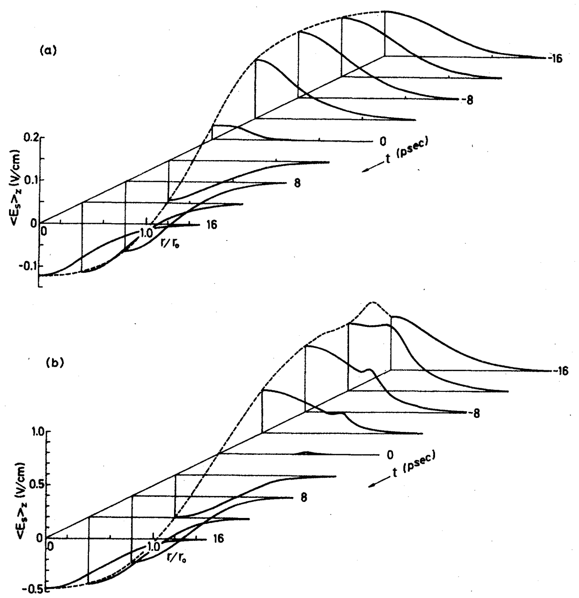

FIG. 4. Temporal changes in the radial distribution of axial electric field generated under the same conditions as in Fig. 3. (a) $I_0 = 8\times10^{11}$ W/cm^2; (b) $I_0 = 3\times10^{12}$ W/cm^2.

B. Helium gas

Multiphoton ionization of He atoms is a twenty-two photon process at the Nd:YAG laser frequency. The excitation frequency ω_i of He atoms is much larger than the laser frequency ω, i.e., $\omega_i \gg \omega$. Hence we neglect the atomic dispersion. No magnetic field is generated in this case. The

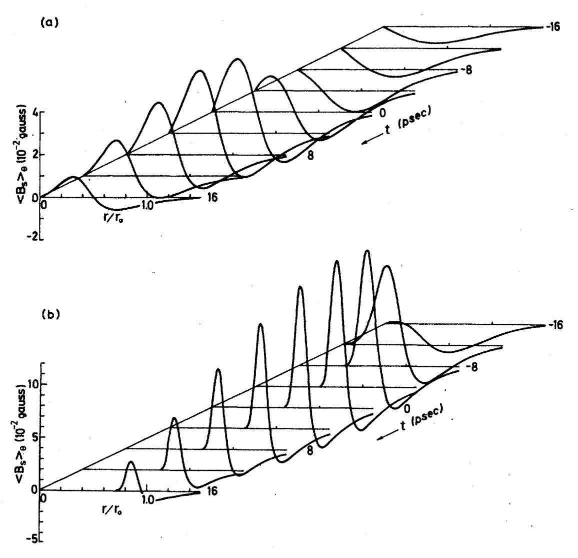

FIG. 5. Temporal changes in the radial distribution of magnetic field generated under the same conditions as in Fig. 3. (a) $I_0 = 8\times10^{11}$ W/cm^2; (b) $I_0 = 3\times10^{12}$ W/cm^2.

following constants are used in the calculations: The multiphoton ionization probability $W = 1.2 \times 10^{-321} I^K$ sec^{-1}, which is taken from Lompre *et al.*,[22] where I is in W/cm^2 and $K = 22$; the spot size $r_0 = 15$ μm; the incident peak intensity $I_0 = 1.5\times10^{15}$ W/cm^2.

The results are shown in Fig. 6. As predicted

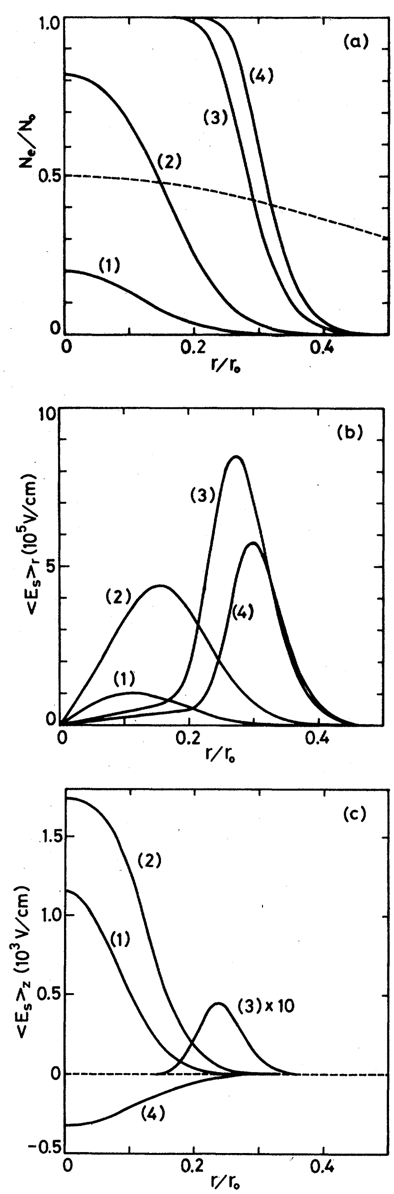

FIG. 6. Radial dependences of (a) multiphoton ionization rate, (b) radial electric field, and (c) axial electric field in He gas at the time t = (1) −5.5, (2) −3.8, (3) 0, and (4) 9.5 psec. The incident Nd:YAG laser pulse is assumed to have the duration τ = 28 psec, the spot size r_0 = 15 μm, and the intensity $I_0 = 1.5\times10^{15}$ W/cm^2. The dashed line indicated in (a) represents the intensity distribution of the incident laser beam.

above, because of the high-order process of multiphoton ionization, the small gradient of the laser intensity distribution is found to produce a steep gradient of N_e/N_0 in Fig. 6(a), and the strong fields $\langle E_s \rangle_r$ and $\langle E_s \rangle_z$ are demonstrated in Figs. 6(b) and (c), respectively.

One can easily predict some phenomena induced by the intense dc electric field in atomic media. The intense field will induce the dc Stark effect in the atoms and ions, which may be the Stark shift and mixing of atomic and ionic energy levels, the lowering of ionization potential, and so forth.

V. CONCLUSION

We have investigated the static electromagnetic fields which are spontaneously generated in the laser-induced multiphoton ionization process in atomic media. A theory is presented which gives the expressions for the nonvanishing radial and axial electric fields $\langle E_s \rangle_r$ and $\langle E_s \rangle_z$, respectively. and the azimuthal magnetic field $\langle B_s \rangle_\theta$. The characteristic natures and magnitude of the spontaneous fields have been demonstrated by calculations on atomic Na vapor and He gas, assuming an incident mode-locked Nd : YAG laser pulse.

The spontaneous fields, as well as multiphoton ionization, are inevitably induced in the interaction of an intense laser pulse with atomic media. Various effects of the fields on the interaction may be expected. We have recently observed optical second-harmonic generation induced by the spontaneous electric field in atomic Na vapor.[23] The second-harmonic generation observed is an experimental evidence for the spontaneous field generation predicted in this paper. The experiment will be described in detail in the following paper.[24]

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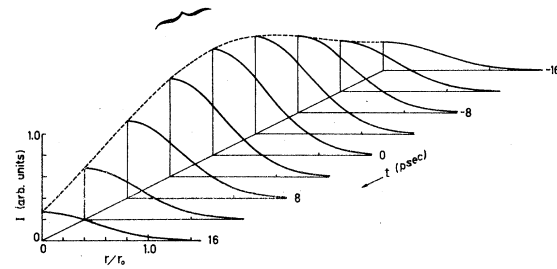

FIG. 7. Temporal change in the radial intensity distribution of the Gaussian laser beam used in the calculations. The peak intensity is normalized to unity at $r=t=0$.

APPENDIX

The explicit forms for the functions ψ and ϕ used in the calculations in Sec. IV are

$$\psi(r) = \exp(-2r^2/r_0^2) \tag{A1}$$

and

$$\phi(t) = \exp(-4t^2/\tau^2)\,. \tag{A2}$$

The temporal change in the intensity distribution given by Eq. (24) is illustrated in Fig. 7, where the peak intensity I_0 is normalized to unity at $r=t=0$.

The general form for the spatial distribution of lowest-order Gaussian beam focused at the center ($z=0$) of a medium may be given by[25]

$$\psi(x,y,z) = (1+\xi^2)^{-1}\exp[-2kr^2/b(1+\xi^2)]\,, \tag{A3}$$

where $b=2\pi r_0^2/\lambda$ is the confocal parameter, $k=2n\pi/\lambda$ is the wave number, and $\xi=2z/b$ is the normalized z coodinate. Equation (A1) which is the formula for ψ at $z=0$ can be regarded as an approximated form for the focused Gaussian beam in the near field region $|z| \lesssim b$. In addition, the refractive index $n=1$ is assumed in Eq. (A1), since the inclusion of the appropriate form of n, which can be easily derived from Eq. (12), makes only small high-order contribution of δ and γ to $\langle \vec{E}_s \rangle$ and $\langle \vec{B}_s \rangle$ under the conditions concerned.

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[13]We therefore assume that the high-frequency polarization induced in the medium is proportional to the incident field, to the first approximation.

[14]If anomalous dispersion as well as absorption is not so large, which is the case of interest, the velocity v is equal to the group velocity v_g defined in the usual way: $v_g = \partial\omega/\partial k$, where k is the wave number.

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