Method of concentration of power in materials for x-ray amplification

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Recent experimental and theoretical results indicate that a new technique for the controlled concentration of power in materials may be feasible. The power levels that are potentially achievable are sufficient for the generation of amplification of x-ray wavelengths in the kilovolt range. The method of power concentration involves the combination of (1) a new ultrahigh brightness subpicosecond laser technology, (2) multiphoton coupling to atoms and molecules, and (3) a new channeled mode of electromagnetic propagation. The energy scaling of this approach is the most important consideration, and it is shown that the control of the propagation is the key factor that enables high levels of amplification in the kilovolt regime to be achieved with a total excitation energy of ~ 1 J.

I. Introduction

The controlled concentration of power in materials is the fundamental issue for the creation of bright sources of radiation in the x-ray range. If amplification is desired, this issue acquires a heightened prominence, since gain at these wavelengths (1–10 Å) requires truly prodigious energy deposition rates[1] spatially organized in a high-aspect-ratio volume of matter. The confluence of advances in femtosecond laser technology, the results of a range of physical measurements, and theoretical analysis all suggest that electromagnetic coupling, in the strong-field regime ($\epsilon \gg e/a_0^2$), may be capable of establishing the demanding conditions that are necessary for such amplification.[2] This paper describes research on questions pertinent to these issues, particularly with regard to certain aspects of the scaling.

Certain general aspects of the scaling relations that govern the conditions of excitation can be seen by examining the geometry of the excited volume described in Fig. 1. It is desired to organize the excitation of the volume so that (1) the maximum volume specific rate for the production of excited species (ρ^*) is achieved and (2) the maximum exponential gain G is produced per unit of deposited energy E. The first can be accomplished with nonlinear processes of absorption[2] by longitudinal excitation with the pulse at wavelength λ as shown in Fig. 1. The second requirement, which is predicated on the desire to minimize the energy of excitation, leads to a condition on the geometry of the excited volume that has direct implications concerning the physical mechanisms that govern propagation. We can write the energy as

$$E = \frac{\hbar\omega_x \rho^* \delta^2 l}{\eta_x}, \tag{1}$$

assuming that $\hbar\omega_x$ represents an x-ray quantum and η_x denotes an appropriately defined intrinsic efficiency for the production of specific excited states radiating at the wavelength λ_x. Correspondingly, gain G can be written as

$$G = \rho^* \sigma_x l, \tag{2}$$

where σ_x represents the stimulated emission cross section for the transition at wavelength λ_x. Hence, in order to minimize the energy of excitation, the

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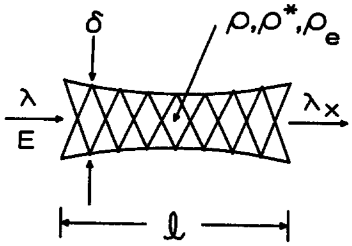

Fig. 1. Excited volume of length l and diameter δ. It is assumed that the excitation is provided by strong-field processes driven at wavelength λ with a pulse of radiation traveling along the length of the volume. The deposited energy is E, the medium density (heavy bodies) is ρ, the excited state density is ρ^*, ρ_e is the electron density, and λ_x represents the wavelength of an assumed amplified x-ray output.

ratio that we wish to maximize is given simply by

$$\frac{G}{E} = \frac{\eta_x \sigma_x}{\hbar\omega_x \delta^2}. \tag{3}$$

In Eq. (3), $\hbar\omega_x$ is given and both η_x and σ_x have upper bounds based on rigorous fundamental considerations. Consequently, the diameter δ remains as the sole variable potentially amenable to independent control. Obviously, the control of δ over a suitable range of small values is preferred. However, an optimally arranged amplifier has its length l independently set equal to the loss length, a condition that we now assume. However, if this length is now set equal to the Rayleigh range and if the propagation of the radiation at wavelength λ is governed by the laws of physical optics *in vacuo,* the condition

$$\delta \geq \sqrt{\lambda l} \tag{4}$$

applies, so that the desired small value of δ require correspondingly small lengths l. If, for example $\delta \sim \lambda$, then $l \sim \lambda$; small beam diameters, which are associated with high propagating intensities, perforce give low values of l/δ as well as small volumes of excitation. The important conclusion is that this restriction does not apply if a guided mode of propagation is used. Therefore, the existence of physical processes that can lead to confined modes of propagation in the strong-field regime basically offer the opportunity to change the scaling relationships that govern the gain G, the volume of excitation, and the excitation energy E for the generation of x-ray amplification. Importantly, there is accumulating theoretical evidence that such self-confined modes of propagation will exist under appropriate conditions.[3–5] This issue is addressed specifically in Section IV.

II. Laser Technology

Advances in femtosecond lasers are extending the exploration of multiphoton interactions well into the regime for which the external field is greater than an atomic unit (e/a_0^2). The performance projected for ultraviolet rare-gas halogen technology is currently being realized[2,6–15] and new near-infrared solid-state systems, such as $Tl{:}Al_2O_3$, are under rapid development.[16] Both technologies should reach a field strength of $\sim 100\ e/a_0^2$ with instruments that produce an output energy of ~ 1 J. With these experimental tools, intensities that approach 10^{21} W/cm^2 will be available.

The KrF* 248-nm source that is used for the studies discussed herein is schematically shown in Fig. 2. This instrument currently has an output pulse width of $\simeq 600$ fs, a typical output pulse energy in the 300–400-mJ range, and good focusability,[12] nominally within a factor of 2 of the diffraction limit. For example, with a simple $f/10$ CaF_2 lens, peak intensities of $\sim 10^{17}$ W/cm^2 are produced even though such a lens exhibits appreciable spherical aberration. This intensity corresponds to a peak electric field slightly above one atomic unit. The application of more sophisticated focusing systems can clearly generate substantially greater field strengths.[12–14] Future refinements of this system, which are currently being implemented in connection with both the source and the focusing system, are expected to lead to the limiting performance of $\sim 100\ (e/a_0^2)$, as noted above.

III. Specific Energy Deposition Rate

Studies of collision-free ionization have been a valuable source of information on multiphoton processes.[17–27] In particular, such experiments give a direct and unambiguous measure of the scale of the energy transfer that occurs between the radiation field and the free atomic or molecular target. Although most studies have been conducted with atoms, recent work has begun to examine molecules, including certain polyatomic systems.

The basic overall process under examination can be presented as the reaction

$$N\gamma + X \rightarrow Xq^+ + qe^- + \gamma, \tag{5}$$

in which multiply charged ions Xq^+, photoelectrons qe^-, and photons γ are produced. It should be noted

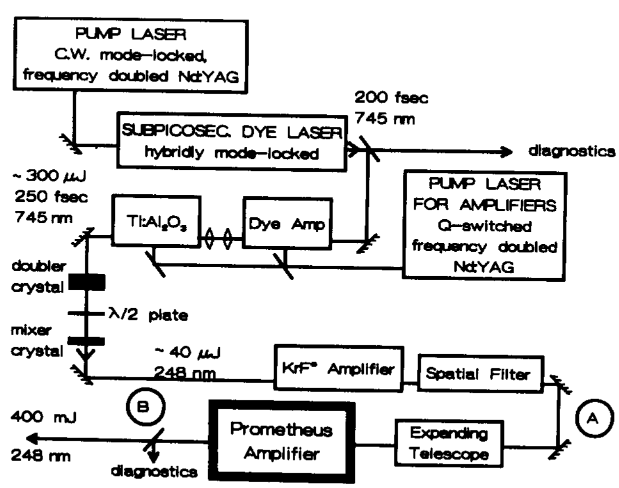

Fig. 2. Schematic showing the configuration of the ultrahigh intensity KrF* 248-nm laser system that was used in the studies discussed in the text.

that for a wavelength of 248 nm, intensities comparable with the Compton intensity, $\sim 5 \times 10^{19}$ W/cm^2, will cause strongly relativistic electronic motions to occur.

The scale of the energy transfer that occurs in subpicosecond strong-field interactions has undergone both experimental study and theoretical analysis. Recent studies[17,25] with 248-nm radiation indicate that the production of Kr-like Xe^{18+}, involving the complete removal of the $4d$ shell, will occur at an intensity of $\sim 10^{18}$ W/cm^2. This case corresponds to a total energy transfer of ~ 3.5 keV/atom. Interestingly, earlier estimates[28] of the strength of coupling, based both on available data and a simple physical model of the induced electron motions, give a total energy transfer of ~ 4 keV/atom at an intensity of 10^{18} W/cm^2, with the assumption of a pulse length of ~ 100 fs. In recent experiments Kr^{11+} has been observed with an intensity of $\sim 4 \times 10^{17}$ W/cm^2, a case that involves a total energy transfer of ~ 1350 eV. The sum of these results supports the conclusion that multikilovolt energy transfers are associated with sufficiently heavy materials for intensities above $\sim 10^{18}$ W/cm^2. Indeed, on account of the mechanisms described in Section IV governing propagation,[3–5,29–31] the range of intensities of interest for 248-nm radiation falls into the 10^{19}–10^{21}-W/cm^2 region, a zone in which the characteristic energy transfers are expected to be greater than ~ 10 keV. Thus, the strong-field multiphoton coupling, viewed in the local spatial sense, is clearly sufficient to produce copious densities of excited material whose spectroscopic structures are associated with electromagnetic transitions in the kilovolt range.

IV. Strong-Field Propagation

As discussed in Section I, the spatial organization of the deposited energy is the crucial aspect that governs the scaling relationship between the magnitude of amplification G and the total deposited energy E. This result is expressed chiefly in Eq. (3). It is obvious from the discussion that pertains to that statement that the desired spatial control of the deposition is fundamentally dependent on the nature of electromagnetic propagation in plasma under strong-field conditions. Therefore, considerable analysis of this question has been made.[3–5,29–38]

These analyses indicate that a fundamentally new regime of electromagnetic propagation is expected to arise in plasmas for subpicosecond radiation at a sufficiently high intensity. The basic formalism and procedures that describe these calculations have been described previously.[4,30] Assuming collisionless and lossless propagation[4] ($\mu = \nu = \nu_0 = 0$), the complex field amplitude $E(t, z, r)$ is governed by the nonlinear Schrödinger equation

$$\left(\frac{1}{c_1}\frac{\partial}{\partial t} + \frac{\partial}{\partial z}\right)E + \frac{i}{2k_0}\Delta_\perp E + \frac{ik_0}{2\epsilon_{R_0}}\delta\epsilon_R E = 0, \qquad (6)$$

where r is the transverse coordinate, $\Delta_\perp = \partial^2/\partial r^2 + r^{-1}\partial/\partial r$, $k_0 = 2\pi/\lambda$, $c_1 = c\epsilon_{R_0}{}^{1/2}$ is the group velocity in the plasma, and c is the vacuum speed of light. The nonlinear term $\delta\epsilon_R(r, |E|^2)$ of key significance, which describes the combined action of the relativistic and charge-displacement mechanisms, is given by

$$\delta\epsilon_R = \left(\frac{\omega_{p_0}{}^2}{\omega^2}\right)\{1 - (1 + I/I_r)^{-1/2}\max \times [0, f(r) + (c^2/\omega_{p_0}{}^2)\Delta_\perp(1 + I/I_r)^{1/2}]\}, \qquad (7)$$

where $I_r = m_e{}^2\omega^2c^3/4\pi e^2$ is the relativistic intensity,[36,39] ω_{p_0} is the plasma frequency, ω is the laser angular frequency, and $f(r)$ denotes the spatial distribution of the initially unperturbed electron density. The last term within the brackets in Eq. (7), $(c^2/\omega_{p_0}{}^2)\Delta_\perp(1 + I/I_r)^{1/2}$, describes the charge-displacement process and was neglected in the earlier studies[4] of purely relativistic propagation. The form of Eq. (7) assumes that the propagating radiation is circularly polarized, since the relativistic γ factor can be written as $(1 + I/I_r)^{1/2}$ in that case. Furthermore, the structure of the terms in Eq. (7) clearly reveals that the relativistic and charge-displacement processes are fundamentally connected and cannot be considered as truly separate and independent mechanisms.

Of primary significance, as noted in Section I, is the possible formation of quasi-stable self-channeled modes of propagation. Separately, the analyses of both the relativistic[4,35–38] and charge-displacement[3] mechanisms give indications that confined modes of propagation can exist. For radiation with a wavelength of 248 nm, the range of intensity of interest is $I > 10^{19}$ W/cm^2, a region in which the multiphoton processes that lead to energy deposition are vigorous and also one in which both relativistic[4] and charge-displacement[3] mechanisms are significant.

The relativistic influence on the electron mass and the ponderomotively driven displacement of charge both encourage the formation of channeled propagation.[3–5,30] On the basis of the dynamic picture that is revealed by the calculations, the combined effect of these two mechanisms can be, in reasonable approximation, summarized in the following simple way. The relativistic effect leads to the initial concentration of the radiation, and the resulting displacement of electronic charge reinforces this tendency and stabilizes the confinement. The strongly cooperative nature of this action appears to lead to highly stable conditions of propagation.[5,30] The effect of the charge displacement is large and cavitation of the electron distribution is a general feature of the solutions for the intensity range of interest. The results for a specific example show that approximately one half of the incident power (~ 4 TW) can be trapped in the channels that, for a wavelength of 248 nm, an electronic plasma density of $\sim 7.5 \times 10^{20}$ cm^{-3}, and an initial beam radius of ~ 3 μm, develop propagating intensities of $\sim 10^{21}$ W/cm^2. An intensity of this magnitude, which involves a field with an electric amplitude of the order of ~ 100 $(e/a_0{}^2)$, is clearly such that **very** high levels of ionization and excitation are expected to occur in any medium.[19–21,24–27]

V. Illustrative Example

The physical conditions discussed in the sections above can be used to estimate the energy of excitation required, the gain, and the output energy that could potentially be achieved. This is done simply to establish an approximate scale of magnitudes that stems from this analysis.

In order to evaluate the energy E expressed in Eq. (1), we assume $\hbar\omega_x = 1$ keV, $\delta = 3$ μm, $l = 1$ cm, and $\rho^* = 10^{20}$ cm^{-3}, parameters that are all consistent with the physical processes considered. This gives $E \simeq 1.4/\eta_x$ (mJ). For an efficiency $\eta_x \simeq 10^{-3}$, the resulting energy falls in the neighborhood of a joule.

A radiatively broadened transition at wavelength $\lambda = 1$ nm has a cross section $\sigma_x \simeq 4 \times 10^{-16}$ cm^2. If we allow for additional broadening to originate from the Doppler effect and the influence of the plasma environment, a cross section in the range of $10^{-18} \leq \sigma_x \leq 10^{-17}$ cm^2 could reasonably be expected. From Eq. (2), this would give $10^2 \leq G \leq 10^3$, a range sufficient for substantial amplification. Since the perturbation of the ionic levels by the optical Stark effect from the intense laser field can substantially lower the gain, we have assumed for this estimate that the lifetime of the upper state of the x-ray transition exceeds the pulse width that is characteristic of the excitation (~ 100 fs).

The efficiency η_x' of generation of the x-ray output E_x from such a system cannot, with the parameters defined above, exceed η_x. A reasonable possible expectation for this efficiency could be $\eta_x \geq \eta_x' \geq \eta_x/10$, a range that gives $0.1 \leq E_x \leq 1.0$ mJ. Interestingly, even the low end of this x-ray output is sufficient for the production of x-ray microholograms of biological objects.[40,41]

Although it is not within the intended scope of this work to provide an analysis of the dynamics of the population inversion, such inversions could be produced by several mechanisms. They include recombination, collisional excitation, direct multiphoton excitation of ions, and possibly the production of excited ions arising from surface crossings encountered in the dissociation of highly ionized molecules. Clearly, the specific dynamics of the electrons in the relativistic regime[34] would have to be considered in connection with the mechanisms that involve collisions and recombination. There is experimental evidence[42] that direct multiquantum molecular excitation is possible as are theoretical grounds for the selective generation of excited atomic ions caused by the dynamic action of crossings in pathways of dissociation that occur in highly charged multiphoton-produced molecular ions.[43]

VI. Conclusions

In terms of the basic question that concerns the controlled deposition of energy at high specific powers, the answer that emerges from this analysis involves three separate components. They are (1) a new ultraviolet pulsed power technology, (2) an energy deposition mechanism based on highly nonlinear coupling, and (3) a condition for channeled propagation. In can be seen that these three elements, which appear capable of producing conditions comparable to or possibly exceeding those of a thermonuclear environment, fit together in a remarkably congenial way. A principal finding is that the radiative conditions needed for the strong multiphoton coupling that governs the energy transfer rate are essentially identical to those required for the channeled propagation and that the laser technology can readily produce the necessary regime of irradiation. The compatibility of these three factors is a key feature of this method for the attainment of high energy density states[43] of matter and stimulated emission in the x-ray range.

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