Factors Controlling the X-Ray Pulse Emission from an Intense Femtosecond Laser-Heated Solid

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The evolution and radiation of strongly heated and ionized solid density material is calculated for conditions which are produced by an intense, femtosecond laser pulse. It is found that the spectrally integrated radiation emitted in the frequency range $h\nu > kT_e$, where T_e is the initial peak plasma temperature, can be as short in duration as ~ 100 fs if kT_e is in an optimum range, set by the target material chosen. For temperatures in this range, the radiation pulse duration is controlled primarily by hydrodynamic expansion. Low x-ray yields can be attributed to suppression of high-ion-stage populations by the high rate of three-body recombination in solid density plasma.

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The interaction of high-energy femtosecond laser pulses with solid material is an area of much current interest. Special attention has been directed at producing ultrashort x-ray pulses from the resulting hot, then rapidly cooled plasma [1], and on studying the novel properties of the very highly excited solid density material [2]. Experimentally, it is as yet unknown how short the emitted x-ray pulse can be, with current measurements instrument limited to about 2 ps [1]. In consideration of the possible duration and intensity of x rays emitted from such a hot, rapidly evolving source, attention has naturally focused on the relevant plasma cooling processes since, at least for a rapidly varying sequence of ionization equilibrium states, the radiation emission will be a temporally local function of the temperature. Previous calculations have indicated the importance of both thermal conduction and hydrodynamic expansion to the production of short pulses [3], while others have proposed special mixed-element targets with the assumption that hydrodynamic expansion and radiative recombination are small effects [4].

In this paper, we show that the dominant emission from a single-species slab target, heated to a depth consistent with a 100-fs heating pulse, can be nearly as short as this pulse, and that the most important process in controlling the emission duration is hydrodynamic expansion. It will also be seen that the criterion of short pulse emission imposes an optimum range for plasma heating set by the target material, and that the ionization dynamics at solid density is responsible for the low x-ray yield both calculated here and previously measured [1,3].

The evolution of a solid target is calculated given a specified temperature profile behind a step function interface at $t=0$. The profile is consistent with the heating provided by a 100-fs laser pulse. The laser energy deposition process is not modeled; our interest is in the impulse response of the system, which would indicate a minimum for the x-ray pulse width produced by the plasma. A relatively low-Z element, carbon, is the material modeled, mainly for the purpose of reducing the computation time. The conclusions, however, can be extended to higher-Z materials, as will be discussed later.

The hot plasma is modeled by the following "one-fluid" equations [5] for mass, momentum, and energy conservation, where ρ and v are mass density and mass weighted velocity, $P_e = N_e kT_e$ is the electron pressure expressed through the ideal-gas equation of state (N_e is the electron density and T_e is electron temperature), $\varepsilon = (\sum_Z N_Z \chi_Z + \frac{3}{2} N_e kT_e)/\rho$ is the internal energy per unit mass (the first term is energy stored in the plasma ionization state; the second term is electron thermal energy), and q is the heat flux:

$$\frac{\partial \xi_i}{\partial t} + \frac{\partial}{\partial x}(\xi_i v + \phi_i) = 0 . \tag{1}$$

Here, $\xi_1 = \rho$, $\phi_1 = 0$; $\xi_2 = \rho v$, $\phi_2 = P_e$; and $\xi_3 = \rho\varepsilon + \rho v^2/2$, $\phi_3 = P_e v + q$. Plasma quasineutrality is assumed through the relation $N_e = \sum_{Z=1}^{6} Z N_Z$, where N_Z is the density of carbon ions of charge $+Z$. At solid densities, this is a very good approximation, since spatial charge separations of at most a few angstroms can take place. The use of one spatial dimension for the plasma dynamics is based on the extremely small ratio of evolving plasma parameter scale lengths (a few thousand angstroms) to the typical laser spot size (≥ 10 μm), for times less than ~ 10 ps. It is assumed that the ion kinetic temperature is zero; on the time scales considered here, electron-ion energy transfer is mainly through the ionization, excitation, and recombination of the ions. A gradient heat conduction term with Spitzer thermal conductivity [6] was used for q [$q \propto (T_e^{5/2}/Z_{\text{eff}})\partial T_e/\partial x$]. The thermal conduction may actually be considerably less than this, either through the high resistivity of hot, solid density plasma [2], or through flux-limited heat flow [7], but even at the levels assumed here, the effect of thermal conduction will be seen as of limited importance to the generation of ultrashort x-ray pulses in the range $h\nu > kT_e$.

A collisional-radiative model, coupled to Eqs. (1), is used to calculate the time-dependent ionization state by means of the following rate equations, which couple the population densities of the ground states of the carbon ions and neutral carbon:

$$\frac{d}{dt}N_Z = S_{Z-1}N_{Z-1}N_e - (S_Z + \alpha_Z)N_Z N_e + \alpha_{Z+1}N_{Z+1}N_e , \tag{2}$$

where $S_Z = S_Z(N_e, T_e)$ and $\alpha_Z = \alpha_Z(N_e, T_e)$ are the collisional-radiative ionization and recombination rates [8], respectively. It is important to note that these rates are *not* the same as the ground-state rates [9] obtained by averaging $N_Z \sigma v$ over the electron distribution function, where σ is the particular cross section for recombination into or ionization out of the ground state of each ion. Rather, they are *effective* rates obtained by solving a much larger set of equations which include the effect of collisional excitation and deexcitation, spontaneous emission, ionization, and recombination (two and three body) with respect to the excited states of each ion. Through a reduction technique based on the near equilibrium behavior of the excited states, Summers [8] has expressed these processes in terms of effective rates coupling the ground states of each ion. In this manner, stepwise ionization and recombination with respect to excited states has been implicitly included in the calculation. As a check, it was confirmed that the solutions to Eqs. (2) converged to those given by the ground-state rates at high T_e and low N_e and to those given by local thermal equilibrium [10] (LTE) at low T_e and high N_e.

Expressions for bremsstrahlung and two-body recombination radiation were obtained from Ref. [11]. The radiation was calculated *after* Eqs. (1) and (2) were solved: The total emission amounts to less than 1% of the initial plasma energy and so the radiation contribution to the energy balance is small. The high-density plasma effect of continuum lowering [10] was taken into account by considering radiation from two-body recombination only into the two most strongly bound states of each ion and spontaneous emission from the two lowest excited states. Line radiation was calculated by assuming that the populations of the two lowest excited states of each ion were in LTE with their corresponding ground-state populations, a reasonable approximation at near-solid density. The χ_Z's of Eqs. (1) were also adjusted using a simple model [10] for continuum lowering. Opacity effects are ignored. For the heated thicknesses considered here, the radiating regions are optically thin to the continuum emission [12]. The self-absorption of line radiation, while not an important effect for carbon, will increase at higher Z. At high Z (and sufficient temperature), where spontaneous emission competes more effectively with collisional deexcitation, line radiation becomes more important, and opacity could increase the total radiation pulse duration.

As an initial condition at $t=0$, the peak temperature at the interface is taken to be 200 eV, and decays in a half-Gaussian profile of 500 Å width (approximately 100 J/cm^2 absorbed): An initial thin heated region (corresponding to a laser-deposition skin depth of 100–200 Å) can penetrate into the solid, via thermal conduction, by as much as 500 Å in 100 fs [12,13]. The temperature of the bulk material beyond the original temperature profile is taken to be 5 eV, in order to provide nonzero thermal conduction for the numerics. We consider radiation of frequency in the range of the carbon K shell. Since recombination of fully stripped carbon and recombination and spontaneous emission of hydrogenic and heliumlike carbon contribute to this radiation, the maximum temperature of 200 eV was chosen since it is just sufficient to fully ionize carbon near the original interface. The effect of higher and lower initial temperatures will be discussed shortly.

In an effort to understand the relative effects of the two main plasma cooling mechanisms, expansion and thermal conduction, three calculations were performed: (a) time-dependent collisional-radiative model [Eqs. (1) and (2)] with conduction "on" [thermal conduction included in Eqs. (1)], (b) steady-state collisional-radiative model [Eqs. (1) and (2) with the left side of Eqs. (2) set equal to zero] with conduction "on", and (c) the same as (a) but with conduction "off." Figure 1 shows the space-integrated recombination, bremsstrahlung, and line radiation emitted from the plasma under conditions (a), (b), and (c) integrated over the energy intervals (i) 200–1000 eV and (ii) 10–100 eV. Recombination radiation dominates line and bremsstrahlung radiation in interval (i) (the line radiation is weaker due to collisional deexcitation of excited states), while recombination and brems-

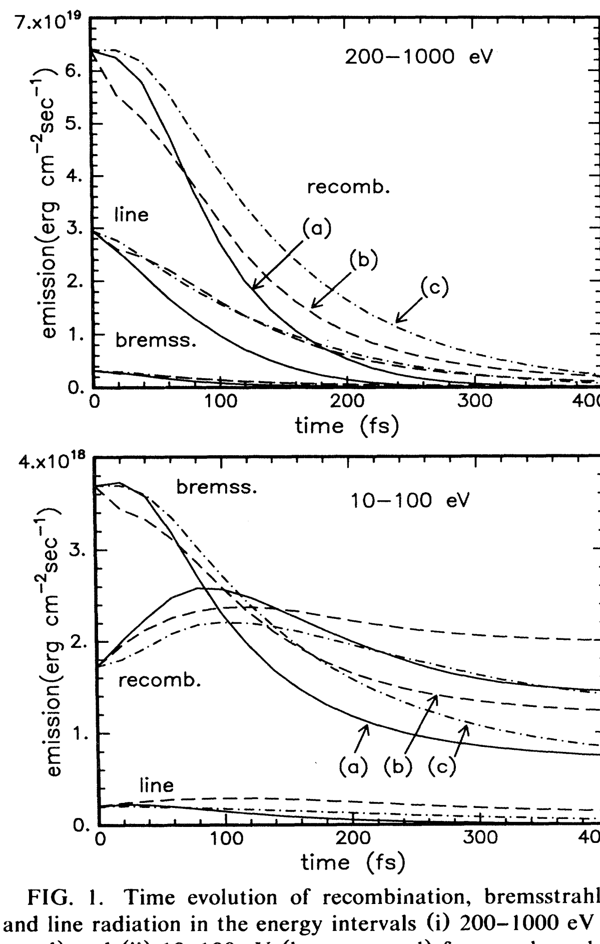

FIG. 1. Time evolution of recombination, bremsstrahlung, and line radiation in the energy intervals (i) 200–1000 eV (top panel) and (ii) 10–100 eV (bottom panel) for a carbon plasma under conditions (a), (b), and (c) (see text).

strahlung emission are comparable in region (ii), dominating line radiation. For the higher-energy interval (i), the radiation decays to half peak value in less than 150 fs for (a), (b), *and* (c). The lower-frequency radiation of (ii) is of much longer duration. The weak dependence of the 200–1000-eV emission intensity and duration on the presence of thermal conduction can be explained by an examination of Fig. 2, which shows the spatial dependence of the recombination emission at various times. The dominant emission of interval (i) clearly comes from the original heated region, which is rapidly cooled by hydrodynamic expansion and where the high ion stages are rapidly decreasing in density. The emission time scale is consistent with the hydrodynamic cooling mechanism: A plasma sound wave ($c_s \sim 2\times10^7$ cm/s) propagates across the original heated layer width in about 200 fs. In the bulk material, high ion stages are suppressed, even at the high temperatures produced there through conduction, because of the high rate of three-body recombination at solid density. This reduces the amount of high-energy recombination and line radiation from the bulk. For the 10–100-eV interval, however, thermal conduction is quite effective in distributing the emitting region well into the bulk. This emission is primarily due to recombination to lower stages of ionization and bremsstrahlung (not shown).

The effect of a time-dependent [(a)] or steady-state [(b)] analysis can be determined by the examination of Fig. 2. For (b), the state of ionization is forced to directly track the evolution of T_e, while for (a), the rapid changes in T_e cause an ionization lag resulting in less bulk penetration of both the 10–100- and 200–1000-eV emitting regions. The temperature changes sufficiently fast that the time-dependent ionization state of (a) never catches up with that of (b). While the steady-state analysis slightly overestimates the radiation emission at long times (see Fig. 1), it gives a reasonable approximation, and at considerably less computation time. In general, for a given ionization structure and temperature change time scale, the steady-state approximation will become worse at higher peak temperatures.

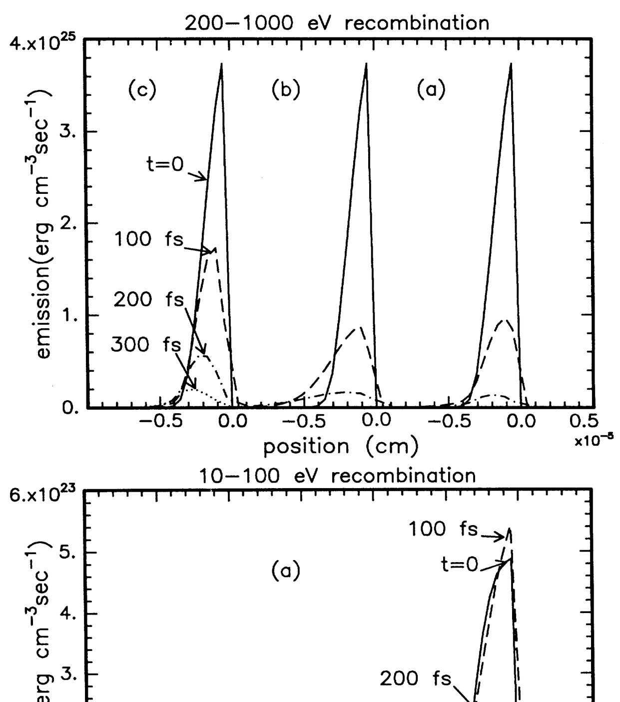

FIG. 2. Spatial distribution of recombination radiation emission at various times for energy intervals 200–1000 eV [top panel, for conditions (a), (b), and (c) (see text)] and 10–100 eV [bottom panel, for condition (a)].

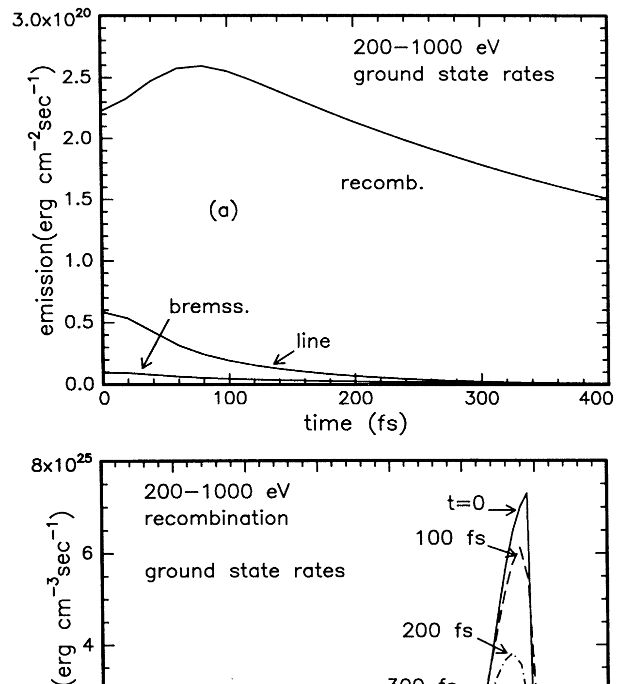

FIG. 3. (a) Time evolution of recombination, line, and bremsstrahlung radiation in the energy interval 200–1000 eV using ground-state ionization and recombination rates in Eqs. (2). (b) Spatial distribution of recombination radiation emission at various times for the energy interval 200–1000 eV, corresponding to (a).

The importance of using collisional-radiative rates in Eqs. (2) is demonstrated in Fig. 3(a), which shows the space-integrated recombination, line, and bremsstrahlung emission in the 200–1000-eV interval when ground-state rates are used. The emission is brighter and of much longer duration than that of Fig. 1. The reason is that the ground-state rates overestimate the degree of ionization in the dense plasma (by neglecting the effect of three-body recombination into excited states), and the thermal front is therefore more effective in ionizing bulk material to high stages [see Fig. 3(b)]. By integrating in time the 200–1000-eV emission curves of Fig. 1, the ratio of emitted energy to energy stored in the plasma at $t=0$ is found to be less than 1%, in reasonable agreement with recent experiments [1], while a similar ratio taken for Fig. 3(a) gives about 10%, much higher than observed. This suggests that lower-average-density targets should be more efficient x-ray emitters.

In Fig. 4, normalized plots of recombination emission in the region $h\nu > kT_e$ are shown for initial peak temperatures of 50 or 600 eV, and it is seen that the radiation decay is longer than for the 200-eV case. It is found that the radiation pulse duration increases for peak initial kT_e satisfying $kT_e \ll \alpha\chi$ or $kT_e \gg \alpha\chi$, where χ, the ionization potential of the primary radiating shell, is in the range 400–500 eV for carbon. (The factor α is in the range $\frac{1}{3} \rightarrow \frac{1}{2}$ and depends on the temperature and density.) At low temperatures, bremsstrahlung and low-ion-stage recombination contribute more strongly to the total radiation, arising from regions heated by the thermal front (see Figs. 1 and 2). At high temperatures, the thermal front can more effectively ionize bulk material to high ionization stages, thus lengthening the resulting radiation pulse. Also, at high temperatures, recombination rates simply are longer. In this regime, thermal conduction plays a more important role in plasma cooling. Clearly, for the shortest emission pulses, the driver laser needs to be adjusted to a level commensurate with the available target element and associated ionization structure.

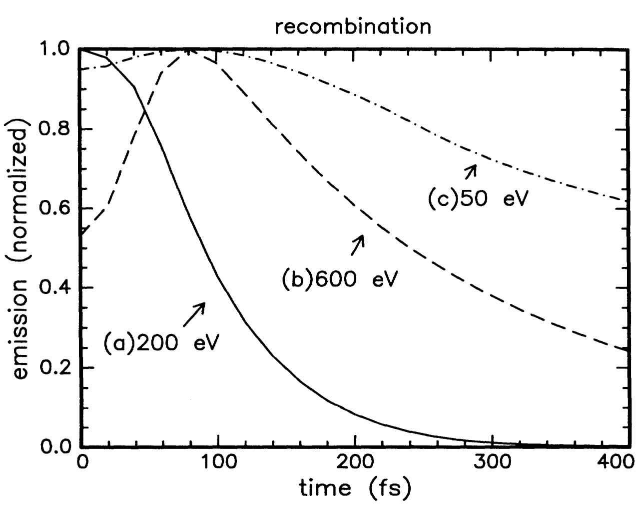

FIG. 4. Normalized plots of space- and frequency-integrated recombination emission for $h\nu > kT_e$ for peak initial temperatures kT_e of (curve *a*) 200 eV, (*b*) 600 eV, and (*c*) 50 eV. The peak emissions are (in erg cm^{-2}sec^{-1}) (curve *a*) 6.4×10^{19}, (*b*) 2.9×10^{20}, and (*c*) 1.2×10^{17}.

The results discussed thus far should also apply to higher-Z materials in the absence of strong self-absorption of line radiation. For similar characteristic emission wavelengths (and therefore temperatures), higher-Z materials would require higher average ionic charge $Z_{\rm eff}$. The energy density stored in kinetic energy of expansion, $\sim \frac{1}{2}\rho c_s^2$, scales as $Z_{\rm eff}T_e$, while Spitzer thermal conductivity scales as $T_e^{5/2}/Z_{\rm eff}$, making hydrodynamic expansion an even more important cooling channel. At the higher electron densities of these targets, three-body recombination would dominate collisional ionization to an even greater extent, with the primary radiating ion stages even more strongly suppressed in the bulk.

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