

# Subsonic radiative ablation model in nonlocal thermodynamic equilibrium fusion plasmas

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In indirect-drive inertial confinement fusion, the radiative ablation process is vitally important for both building up of a uniform x-ray radiation in the high-Z hohlraum and subsequently capsule implosion ablated by the x-ray radiation. For the former circumstance, various supersonic and subsonic ablation models are developed under the local thermodynamic equilibrium (LTE) assumption. However, for the latter, since the capsule plasma is of low  $Z$ , the electron and radiation temperatures are hard to achieve equality; i.e., LTE assumption breaks. So far, how to precisely describe and simulate such a non-LTE ablation process is unknown. Here, by self-consistently including the energy exchange between radiations and electrons into the radiation hydrodynamic (RHD) equations and considering the heat wave velocity as subsonic, we derive and set up a non-LTE subsonic radiative ablation model. Different from previous models where only the Rosseland opacity  $\sigma_R$  functions, in our non-LTE model, besides  $\sigma_R$ , the Planck opacity  $\sigma_P$  plays an important role as well. Our model shows that most characteristic physical quantities of the non-LTE ablation, such as the ablation velocity, ablation pressure, and shock velocity, depend on a combination of  $\sigma_P/\sigma_R$ , while the damping scale length of the radiation temperature scales as  $L \propto (\sigma_P\sigma_R)^{1/2}$ . Therefore, we propose using a simple single multiplier  $\zeta$  on both opacities as  $\tilde{\sigma}_P = \zeta\sigma_P$  and  $\tilde{\sigma}_R = \zeta\sigma_R$  to match the RHD simulation with experimental results, where  $L$ , or more practically, the preheating radiation temperature  $T_{\text{pre}}$ , can be chosen as a proper benchmark.

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## I. INTRODUCTION

Indirect-drive inertial confinement fusion (ICF) [1–8] relies on the building up of a uniform x-ray radiation [9] within a laser-irradiated hohlraum and subsequently on the optimization of the capsule implosion [10] ablated by this x-ray radiation. In both circumstances, the radiative ablation process is vitally important. For the laser-ablated high- $Z$  hohlraum plasmas such as gold, the bremsstrahlung radiation is so strong that the electron and radiation temperatures reach equality rapidly, and therefore, the plasma can be assumed as in a local thermodynamic equilibrium (LTE) state [11–15]. Under the LTE assumption, various radiative ablation models have been derived either from the supersonic Marshak wave theory [16–18] or from a subsonic ablation theory [13–15], which all show that the ablation process depends heavily on the plasma Rosseland opacity, representing domination of the radiation transport.

In contrast, the capsule material is generally of low  $Z$ , such as typically hydrocarbon (CH) [19,20], where the bremsstrahlung radiation is weak; accordingly, the electron and radiation temperatures are unequal [3,20,21]. Therefore,

the LTE assumption breaks and the energy exchange between electrons and the radiation [13,14,18] needs to be taken into account. Furthermore, the characteristic timescale of the capsule ablation process in indirect-drive ICF is typically several nanoseconds, during which the heat wave slows down with velocities lower than the sonic speed [11,13,14]. As a result, the shock forms and the plasma motion needs to be taken into account as well. So far, how to precisely describe and simulate such a non-LTE subsonic radiative ablation process is unknown.

On the other hand, in the RHD simulations, the diffusion approximation is generally used, in which two average opacities, the Rosseland opacity and the Planck opacity, are used, which determine, respectively, the radiation energy transport process and the electron-radiation energy exchange-process [3,22]. However, the values of these opacities used in simulations are generally inaccurate so that the RHD simulation results differ significantly from the experimental results [23,24]. In practical RHD simulations, an artificial opacity multiplier needs to be introduced to match with the experimental data. In particular, to match different characteristic physical quantities measured in the experiments, frequently several different opacity multipliers are applied [23–25]. This is unphysical and difficult to operate.

In this paper, we propose a non-LTE subsonic radiative ablation model that takes into account both the electron-radiation energy exchange and the plasma fluid motion self-consistently. Different from previous models where only the Rosseland opacity  $\sigma_R$  functions, in our non-LTE model, besides  $\sigma_R$ , the Planck opacity  $\sigma_P$  plays an important role as

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well. Our model also shows that most characteristic physical quantities of the non-LTE ablation, such as the ablation velocity  $v_a$ , ablation pressure  $P_a$ , and shock velocity  $v_s$ , depend on a combination of  $\sigma_P$  and  $\sigma_R$  together as  $\sigma_P/\sigma_R$ , while the spatial damping length of the radiation temperature scales as  $L \propto (\sigma_P\sigma_R)^{1/2}$ . As a result, to make the characteristic quantities of the ablation process weakly dependent on or independent of the artificial opacity multiplier in the RHD simulation, a simple single multiplier  $\zeta$  can be used on both opacities as  $\tilde{\sigma}_P = \zeta\sigma_P$  and  $\tilde{\sigma}_R = \zeta\sigma_R$  so that  $v_a$ ,  $P_a$ , and  $v_s$  obtained in simulations can all be automatically matched with the experimental results, where  $L$  can be chosen as a proper benchmark, or more practically, choosing the preheating radiation temperature  $T_{\text{pre}}$ . The theoretical model has been verified by the RHD simulations using the code FLASH [26].

The paper is organized as follows. In Sec. II, we establish the non-LTE subsonic radiative ablation model theoretically and give analyses on the dependences of the characteristic physical quantities in the ablation process on the opacities. The RHD simulation results are presented in Sec. III to verify the theoretical model. Summary and discussion are given in Sec. IV.

## II. THEORETICAL MODEL

We consider the capsule ablation process of indirect-drive ICF in a one-dimensional (1D) planar geometry, shown schematically in Fig. 1(a), where the radiation flux is assumed to be incident from the right boundary. As well known, such a non-LTE ablation process can be described by the complete RHD equations [13,14,18] under the gray mean approximation [18] as follows:

$$\frac{\partial \rho}{\partial t} + \frac{\partial \rho u}{\partial x} = 0, \quad (1a)$$

$$\rho \left( \frac{\partial u}{\partial t} + u \frac{\partial u}{\partial x} \right) + \frac{\partial p}{\partial x} = 0, \quad (1b)$$

$$\rho \frac{\partial e}{\partial t} + p \frac{\partial u}{\partial x} = c\sigma_P a (T_r^4 - T_e^4), \quad (1c)$$

$$\frac{\partial E}{\partial t} + \frac{\partial F}{\partial x} = c\sigma_P a (T_e^4 - T_r^4), \quad (1d)$$

where  $\rho$ ,  $u$ ,  $p$ , and  $e$  are, respectively, the density, velocity, pressure, and specific internal energy of the plasma fluid.  $E$  is the radiation energy density given as  $E = aT_r^4$ , where  $a = \sigma/c$ ,  $\sigma$  being the Stefan-Boltzmann constant and  $c$  being light speed.  $F$  is the radiation flux given as  $F = -(c/3\sigma_R)\partial E/\partial x = -(c/3\sigma_R)\partial aT_r^4/\partial x$ , where  $\sigma_R = \int_0^\infty \partial B_\nu/\partial T d\nu / \int_0^\infty (1/\sigma_\nu)\partial B_\nu/\partial T d\nu$  is the average Rosseland opacity ( $B_\nu$  is the Planck distribution function and  $\sigma_\nu$  is the photon energy-dependent opacity), which determines the radiation energy transport process.

Note that here Eqs. (1a)–(1d) for the radiative ablation processes under the non-LTE condition (generally for low- $Z$  material) are inherently different from those previously under the LTE condition (generally for high- $Z$  material) [13,14]. First, since the bremsstrahlung and inverse bremsstrahlung are weak, the electron temperature  $T_e$  and the radiation temperature  $T_r$  are decoupled, i.e.,  $T_r \neq T_e$  during the ablation [see Fig. 1(a) and blue and red lines in Fig. 1(b)].

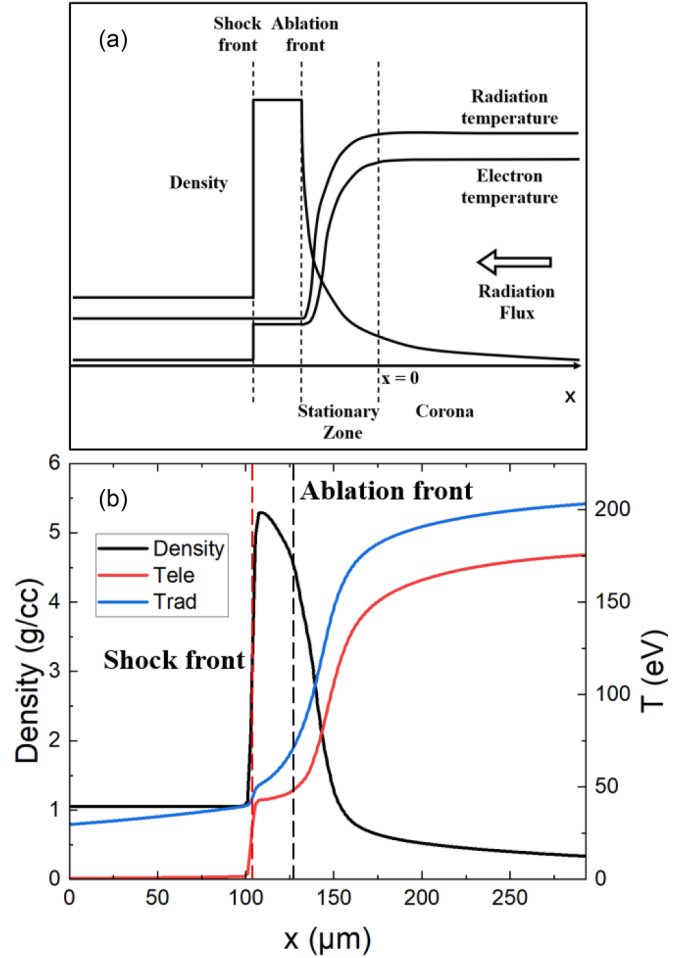


FIG. 1. Typical density and temperature profiles in non-LTE subsonic radiative ablation process: (a) schematic profiles in the theoretical model and (b) the profiles obtained from one typical RHD simulation at time  $t = 4.5$  ns, where the ablated material is CH and the drive radiation temperature is constant at 200 eV. In the figures, the radiation flux is incident from the right boundary. Tele and Trad refer to the temperatures of electrons and radiation, respectively.

As a result, when we describe the dynamics over time of the plasma internal energy, we have to take into account both the bremsstrahlung/inverse bremsstrahlung energy exchange process between the radiation and the plasma electrons and the diffusive transport process of the radiation energy [see the right-hand side of Eqs. (1c) and (1d)]. In contrast, for the previous LTE radiative ablation, due to the strong bremsstrahlung/inverse bremsstrahlung,  $T_e$  and  $T_r$  reach equality so quickly, i.e.,  $T_r = T_e$  during the ablation. Therefore, Eqs. (1c) and (1d) are simplified to only single equation as  $\partial e/\partial t + p\partial u/\partial x + \partial F/\partial x = 0$ , where only the diffusive transport process of the radiation energy needs to be taken into account [13]. Second, further because of the previous point, in contrast with the previous LTE model, where only the average Rosseland opacity  $\sigma_R$  functions (in  $F$ ), besides  $\sigma_R$ , here the average Planck opacity  $\sigma_P = \int_0^\infty \sigma_\nu B_\nu d\nu / \int_0^\infty B_\nu d\nu$  plays an important role as well [see the right-hand side of Eqs. (1c) and (1d)], which determines majorly the electron-radiation energy exchange process. Third,

also different from the previous supersonic models [16–18] where  $T_e$  and  $T_r$  are only concerned physical variables, here the fluid motion [Eqs. (1a) and (1b)] needs to be self-consistently considered [27].

Next, we try to find the solutions of Eqs. (1a)–(1d). However, as shown in the Appendix, from the perspective of mathematical principles, it is impossible to find the time-dependent self-similar analytic solution of Eqs. (1a)–(1d) as the self-similar solution [14,15,18] is unfortunately restricted to extremely stringent parameter constraints that are unable to be satisfied. To overcome this difficulty, based on the characteristics of the plasma density and temperature distribution as well as the ablation wave structure, we separate the ablated capsule plasma into three zones [see Fig. 1(a), from right to left]: the isothermal coronal zone where the plasma density and temperature vary weakly in space from vacuum till the Chapman-Jouguet point [13,28] (where the plasma Mach number  $M \approx 1$ ), the stationary zone where the density and temperature vary dramatically in space and the characteristic plasma flow can be approximately regarded as a stationary wave in the frame of reference moving with the ablation front [13,29], and the shocked region where the plasma density remains highly constant as a very narrow shelf. Typical profiles of density and temperatures in one of the FLASH simulations are also shown in Fig. 1(b), where the main characteristics are consistent with our model. However, since in the more realistic RHD simulations, the shock velocity (or the shock Mach number  $M_u$ ) changes with time during the compression process, the downstream and upstream densities  $\rho_2/\rho_1 = M_u^2(\gamma + 1)/[M_u^2(\gamma - 1) + 2]$  are different during each time of shock compression, which, after time integration, results in a nonuniform density profile in the shocked region [see Fig. 1(b)]. Correspondingly, the shock and ablation fronts do not appear as two separate discontinuity surfaces [see Fig. 1(b)], unlike the ideal schematic of Fig. 1(a).

For the stationary zone, in the frame of reference moving with the ablation front, all the terms related to the derivative over time are much smaller than those related to the derivative over space, which can thus be neglected. As a result, the Hugoniot relations can be used for Eqs. (1a)–(1d) to express the radiation flux  $F$  as a function of the local sound speed  $c_s$ . According to the literature [13], in a subsonic regime, by using the Hugoniot relations for Eqs. (1a) and (1b), we derive that the plasma flow velocity  $v_f$  depends on both the ablation velocity  $v_a$  and the local sound speed at the ablation front  $c_a$  as  $v_f = v_a c_s^2 / c_a^2$ . Furthermore, by using the Hugoniot relations for Eqs. (1c) and (1d), and considering that the low- $Z$  ablated capsule plasma can be fully ionized, as a result, ideal gas approximation with the adiabatic exponent  $\gamma = 5/3$  is used. Then, after derivations, the radiation flux can be derived as  $F = -5\dot{m}_a c_a^2 (1 - v_f/v_a)/2 = 5\dot{m}_a (c_s^2 - c_a^2)/2$  [13], where  $\dot{m}_a = \rho_a v_a$  is the mass ablation rate.

Substituting the above expression of  $F$  into, respectively, Eq. (1d) and its original definition, after derivations, we obtain

$$\frac{5}{2} \dot{m}_a \frac{dc_s^2}{dx} = c\sigma_P a (T_r^4 - T_e^4), \quad (2a)$$

$$\frac{5}{2} \dot{m}_a (c_a^2 - c_s^2) = \frac{c}{3\sigma_R} \frac{d}{dx} a T_r^4. \quad (2b)$$

Considering  $c_a^2 \ll c_s^2$  in most circumstances in the stationary zone, we ignore the high-order items involving  $c_a$  in the equations. Further substituting  $c_s^2 = Zk_B T_e / A m_p$  and  $\dot{m}_a = \rho_a v_a$  into Eqs. (2a) and (2b) and defining  $\epsilon_p = 5Zk_B \rho_a / 2A m_p$  representing the electron energy density (including both the kinetic and internal ones) per unit temperature, we get

$$\epsilon_p v_a \frac{dT_e}{dx} = c\sigma_P a (T_r^4 - T_e^4), \quad (3a)$$

$$-\epsilon_p v_a T_e = \frac{c}{3\sigma_R} \frac{d}{dx} a T_r^4, \quad (3b)$$

where  $Z$ ,  $A$ ,  $k_B$ , and  $m_p$  refer to the average charge number, average atomic weight of the materials, Boltzmann constant, and the proton mass, respectively. Equation (3a) describes how the electrons absorb the radiation energy in the frame moving with the ablation front at the stationary zone, while Eq. (3b) depicts how the radiation flux  $F$  is transferred to electrons, as the right-hand side of Eq. (3b) is actually the simplified expression of  $F$ .

In indirect-drive ICF, generally the dependence of the opacity on the electron temperature [3,13,14,18] can be expressed in the form of the power laws as  $\sigma_R = k_0 \rho T_e^{-\alpha}$  and  $\sigma_P = k'_0 \rho T_e^{-\alpha'}$  [13,14,18], where  $k_0$  and  $k'_0$  are the coefficients and  $\alpha$  and  $\alpha'$  are the relative exponents, determined by the specific plasma material properties. Then, Eqs. (3a) and (3b) become

$$\epsilon_p v_a \frac{dT_e}{dm} = \frac{ck'_0 a}{T_e^{\alpha'}} (T_r^4 - T_e^4), \quad (4a)$$

$$-\epsilon_p v_a T_e = \frac{cT_e^\alpha}{3k_0} \frac{d}{dm} a T_r^4, \quad (4b)$$

where  $m$  is the Lagrangian coordinate defined by  $dm = \rho dx$ . We see that the combination of Eqs. (4a) and (4b) constitutes a set of closed ordinary differential equations (ODEs) of the electron and radiation temperature,  $T_e$  and  $T_r$ . The boundary condition of the ODEs can be defined as  $T_r(0) = T_0$  and  $T_e(0) = T_{e0}$  at the right boundary of the stationary zone  $x = 0$ , i.e., at the Chapman-Jouguet point, and  $T_r(x_0) = T_e(x_0) = 0$  at the left boundary of the stationary zone  $x = x_0$ , i.e., at the ablation front, where  $c_a \ll c_s$  holds.

According to the mathematical principle, ensure that  $T_r$  and  $T_e$  are solvable from the system of ODEs (4a) and (4b) and also all the boundary conditions can meanwhile be satisfied, the ablation velocity  $v_a$  must be a certain value instead of a function of  $m$ , which means that  $v_a$  depends on only the parameters  $k_0$ ,  $k'_0$ ,  $\alpha$ ,  $\alpha'$ ,  $T_0$ ,  $T_{e0}$ , and  $\epsilon_p$ . Therefore, if further defining the dimensionless Lagrangian coordinate  $\xi = \sqrt{k_0 k'_0 / T_0^{\alpha+\alpha'}} m$ , the relations  $\epsilon_p v_a = ca \sqrt{k'_0/k_0} T_0^{3+(\alpha-\alpha')/2} v$ ,  $T_r(m) = T_0 f(\xi)$ , and  $T_e(m) = T_0 g(\xi)$  can be derived, where  $v$ ,  $f(\xi)$ , and  $g(\xi)$  represent the dimensionless ablation velocity, radiation temperature, and electron temperature, respectively. Substituting these dimensionless variables into Eqs. (4a) and (4b), we obtain the dimensionless ODEs as

$$-vg = \frac{g^\alpha}{3} \frac{df^4}{d\xi}, \quad (5a)$$

$$v \frac{dg}{d\xi} = \frac{g^4 - f^4}{g^{\alpha'}}, \quad (5b)$$

where the boundary conditions evolve as  $f(0) = 1$ ,  $g(0) = T_{e0}/T_0$ , and  $f(\xi_0) = g(\xi_0) = 0$ .

Combining Eqs. (5a) and (5b) and eliminating  $d\xi$ , we obtain

$$\frac{df}{dg} = \frac{3v^2g}{4f^3(f^4 - g^4)g^{\alpha-\alpha'}}, \quad (6)$$

where the boundary conditions  $f(g=0) = 0$  and  $f(g = T_{e0}/T_0) = 1$  hold. From Eq. (6), we conclude that the dimensionless ablation velocity  $v(T_{e0}/T_0)$  depends on the parameter of the opacity  $\alpha - \alpha'$ . Generally, for the capsule CH ablated plasmas, we have  $0 < \alpha - \alpha' < 1$  [28]. Figure 2(a) plots the value of  $v(T_{e0}/T_0)$  varying with  $T_{e0}/T_0$  for, respectively,  $\alpha - \alpha' = 0, 0.5$ , and  $1$ , which decreases with increasing  $T_{e0}/T_0$  at all  $\alpha - \alpha'$ .

Finally, we derive the dependences of the characteristic physical quantities on the opacities  $\sigma_P$  and  $\sigma_R$  from Eqs. (5a) and (5b). We define the typical depth that photons can penetrate into the ablated materials or say the width of the absorption zone as  $L = T_r dx/dT_r$  [30], and then its reciprocal  $k_L = (1/T_r)dT_r/dx$  representing the typical spatial damp rate of the radiation temperature. In the non-LTE subsonic radiative ablation regime, as analyzed above, we obtain  $k_L = (1/T_r)dT_r/dx = (\rho/T_r)dT_r/dm = (\rho/f)(df/d\xi)(d\xi/dm) \propto d\xi/dm = \sqrt{k_0 k'_0/T_0^{\alpha+\alpha'}} = \sqrt{\sigma_P \sigma_R}$  and  $L \propto 1/\sqrt{\sigma_P \sigma_R}$ , where  $(1/f)(df/d\xi)$  is determined by Eqs. (5a) and (5b) for certain  $\alpha, \alpha'$ . Further, from the specific dimensionless variable definition, we can obtain the ablation velocity  $v_a = ca/\epsilon_p \sqrt{k'_0/k_0} T_0^{3+(\alpha-\alpha')/2} v(T_{e0}/T_0) \propto \sqrt{k'_0/k_0} T_0^{3+(\alpha-\alpha')/2} v(T_{e0}/T_0) = \sqrt{\sigma_P/\sigma_R} T_0^3 v(T_{e0}/T_0)$ . For the shock velocity, we have  $v_s = \sqrt{(\rho_a/\rho_0)/(1-\rho_0/\rho_a)} c_a \propto c_a \propto \sqrt{c_0 v_a} \propto (\sigma_P/\sigma_R)^{1/4} T_0^{7/4} v^{1/2}(T_{e0}/T_0)$ , where  $\rho_0$  refers to the density in front of the shock and  $\rho_a/\rho_0 = 4$  is the largest limit for the strong shock situation [22];  $c_0 \propto T_0^{1/2}$  here refers to the sound speed in the corona [13]. Similarly, the ablation pressure  $p_a = \dot{m}_a c_0 = \rho_a v_a c_0 \propto \sqrt{\sigma_P/\sigma_R} v(T_{e0}/T_0) T_0^{7/2}$  can also be derived from Refs. [13,28]. As a summary, the dependences of the typical spatial damp rate of radiation temperature  $k_L$ , scale length  $L$ , ablation velocity  $v_a$ , ablation pressure  $p_a$ , and shock velocity  $v_s$  on opacities are given in the following:

$$k_L = \frac{1}{L} \propto \sqrt{\sigma_P \sigma_R}, \quad (7a)$$

$$v_a \propto \sqrt{\frac{\sigma_P}{\sigma_R}} v\left(\frac{T_{e0}}{T_0}\right) T_0^3, \quad (7b)$$

$$p_a \propto \sqrt{\frac{\sigma_P}{\sigma_R}} v\left(\frac{T_{e0}}{T_0}\right) T_0^{7/2}, \quad (7c)$$

$$v_s \propto \left(\frac{\sigma_P}{\sigma_R}\right)^{1/4} T_0^{7/4} v^{1/2}\left(\frac{T_{e0}}{T_0}\right), \quad (7d)$$

where  $\alpha, \alpha'$  are considered to be fixed parameters for certain materials in the opacities. Different from previous models where the characteristic physical quantities depend heavily on only the Rosseland opacity  $\sigma_R$ , we see from Eqs. (7a)–(7d) that most of the characteristic physical quantities of the non-LTE ablation, the ablation velocity  $v_a$ , ablation pressure  $p_a$ , and shock velocity  $v_s$ , depend on a combination of  $\sigma_P$

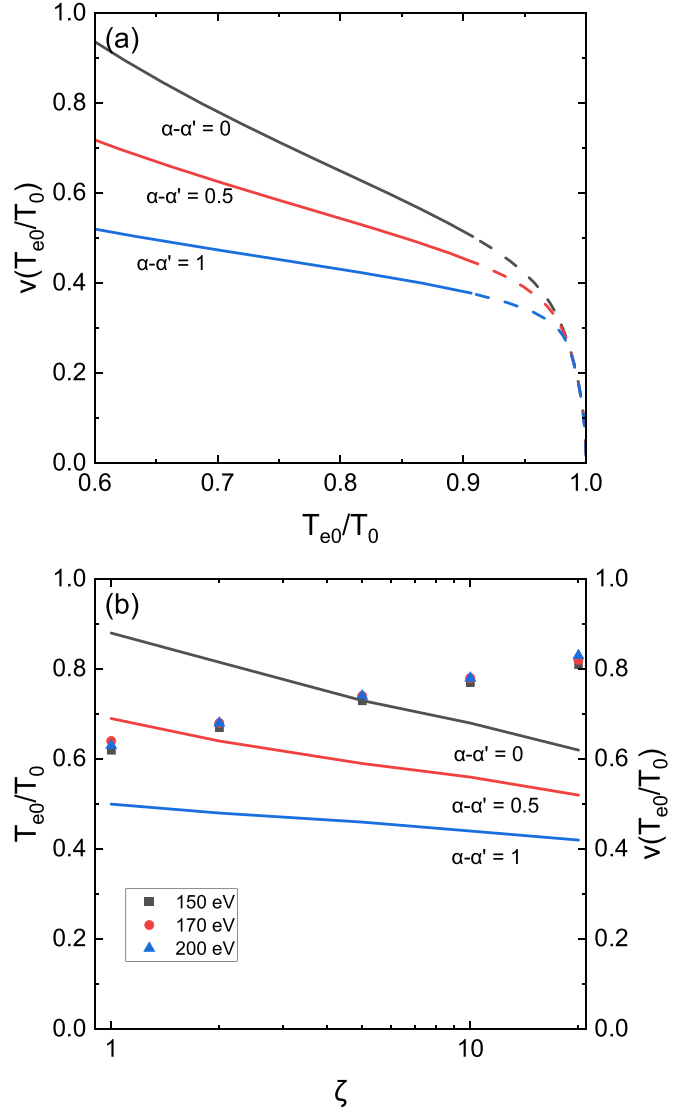


FIG. 2. (a) Dimensionless ablation velocity  $v(T_{e0}/T_0)$  varying with  $T_{e0}/T_0$  for, respectively,  $\alpha - \alpha' = 0, 0.5$ , and  $1$ , which are numerical solutions of Eqs. (5a) and (5b). (b)  $T_{e0}/T_0$  (marked with the symbols of “black square,” “red circle dots,” and “blue triangle” for the drive radiation temperature of, respectively, 150, 170, and 200 eV) obtained from the RHD simulations at different opacity multipliers  $\zeta$ . The black, red, and blue solid lines are the corresponding  $v(T_{e0}/T_0)$  for the cases of  $\alpha - \alpha' = 0, 0.5$ , and  $1$ , respectively. Detailed simulation setup is introduced in Sec. III. The results indicate that the dependence of  $v$  on opacity is rather weak.

and  $\sigma_R$  together as  $\sigma_P/\sigma_R$ , while the spatial scale length of the radiation temperature scales as  $L \propto (\sigma_P \sigma_R)^{1/2}$ . Note that for previously the LTE case, the plasma fluid motion is determined only by  $\sigma_R$  as  $v_a \propto (1/\sigma_R)^{1/2}$ ,  $p_a \propto (1/\sigma_R)^{1/2}$ , and  $v_s \propto (1/\sigma_R)^{1/4}$  [13,14]. Furthermore, in previous LTE models, an opacity multiplier  $\alpha\sigma_R$  is artificially used to match the RHD simulation results with the experimental measured data; however, generally a single fixed multiplier fails to match all the experimental observational data. Here, under the non-LTE condition, to make the characteristic quantities of the ablation process weakly dependent on or independent of the artificial

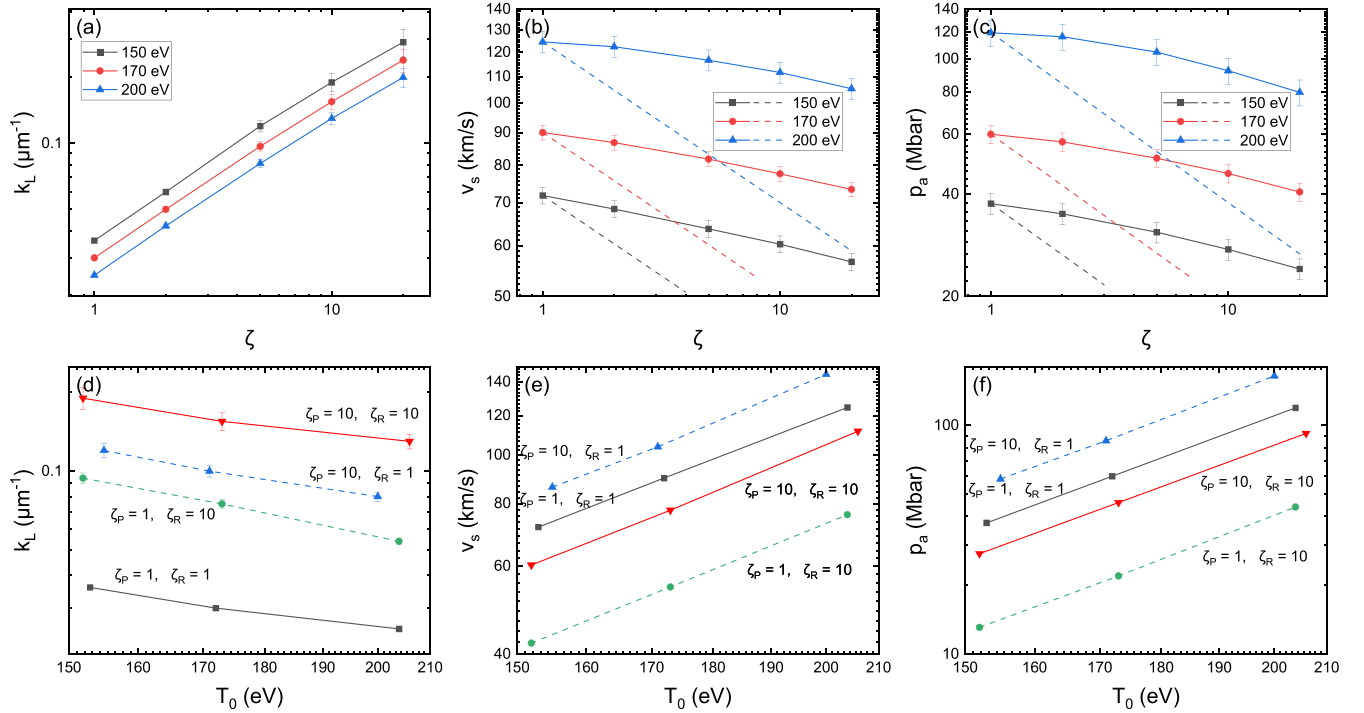


FIG. 3. RHD simulation results by using the non-LTE model: (a) damp spatial scale length of the radiation temperature  $k_L$ , (b) shock velocity  $v_s$ , and (c) ablation pressure  $p_a$  varying with different opacity multipliers  $\zeta$  ( $1 \leq \zeta \leq 20$ ) at drive radiation temperatures of 150, 170, and 200 eV, respectively. The corresponding results calculated from the LTE model [13] are also shown by the dashed lines in panels (b) and (c) for comparison. (d)–(f) Dependences of  $k_L$ ,  $v_s$ , and  $p_a$  on different drive radiation temperatures  $T_0$  for the cases of changing the average Planck and Rosseland opacities independently by using different opacity multipliers  $\zeta_P$  and  $\zeta_R$ : the red solid line for  $\zeta_P = \zeta_R = 10$ , the black solid line for  $\zeta_P = \zeta_R = 1$ , the blue dashed line for  $\zeta_P = 10$  and  $\zeta_R = 1$ , and the green dashed line for  $\zeta_P = 1$  and  $\zeta_R = 10$ .

opacity multiplier in the RHD simulation, we propose to use a simple single multiplier  $\zeta$  on both opacities as  $\tilde{\sigma}_P = \zeta \sigma_P$  and  $\tilde{\sigma}_R = \zeta \sigma_R$  so that  $v_a$ ,  $P_a$ , and  $v_s$  obtained in simulations can all be automatically matched with the experimental results, where  $\tilde{L} \propto (\tilde{\sigma}_P \tilde{\sigma}_R)^{1/2} \propto \zeta$  is chosen as a proper benchmark. The method is physically consistent with the definition of  $\sigma_P$  and  $\sigma_R$  as they both depend on  $\sigma_v$ .

### III. RHD SIMULATION RESULTS

To verify the above theoretical model, we carry out the 1D RHD simulations of the complete radiative ablation process of a CH target through the code FLASH [26]. The simulations are run in 1D Cartesian geometry. Radiation transfer is calculated through the multigroup diffusion approximation [28]. Sixty groups with photon energy ranging from  $10^{-1}$  to  $10^5$  eV are used in these simulations. The left boundary is considered to be the vacuum and the right boundary uses constant drive radiation temperature varying from 150 to 200 eV in different simulation cases. The opacity of CH is calculated through the code IONMIX [31]. The average Planck and Rosseland opacities can be changed independently. Different opacity multipliers are used in different simulation cases for comparisons. The equation of state of CH is calculated through the code BADGER [32].

One of the typical simulation results at time  $t = 4.5$  ns with the drive radiation temperature 200 eV and the opacity multiplier  $\zeta = 2$  is shown in Fig. 1(b), which clearly verifies that the profiles of the plasma density, electron temperature,

and radiation temperature are in good consistence with our theoretical model, in particular for the stationary assumptions. Figure 2(b) shows  $T_{e0}/T_0$  of the other typical simulation results when  $\zeta$  varies from 1 to 20 and the drive radiation temperatures of 150, 170, and 200 eV are chosen. We see clearly that, as expected,  $T_{e0}/T_0$  has a weak dependence on the opacity multiplier  $\zeta$ . When it varies from 1 to 20,  $T_{e0}/T_0$  increases only slightly from 0.62 to 0.82 [see the “black square,” “red circle dots,” and “blue triangle” symbols in Fig. 2(b)]. At this parameter range,  $v(T_{e0}/T_0)$  obtained either analytically from theory [see Fig. 2(a)] or from RHD simulations [see the solid lines in Fig. 2(b)] has a very small change; therefore, the corresponding ablation velocity  $v_a$  also depends weakly on  $\zeta$ . Note that if we extend the parameter  $T_{e0}/T_0$  to be close to 1.0 [dashed lines in Fig. 2(a)], i.e., in transition to the LTE case, we see from Fig. 2(a) that the dimensionless ablation velocity decreases heavily to about 0.1. This implies that the ablation velocity in the non-LTE case is much larger than that in the LTE case.

Figures 3(a)–3(c) plot the damp spatial scale length of the radiation temperature  $k_L$ , the shock velocity  $v_s$ , and the ablation pressure  $p_a$  varying with different opacity multipliers  $\zeta$  ( $1 \leq \zeta \leq 20$ ) at drive radiation temperatures of 150, 170, and 200 eV, respectively. From Fig. 3(a), we see that  $k_L$  increases almost linearly with  $\zeta$ , which agrees with the above theory [Eq. (7a)] as  $k_L \propto \zeta$ . Also, as shown in Figs. 3(b) and 3(c), both  $v_s$  and  $p_a$  decrease much slowly with increasing  $\zeta$ , which can be fitted, respectively, as  $v_s \propto \zeta^{-0.06}$  and  $p_a \propto \zeta^{-0.12}$ , with our theoretical expectations from Eqs. (7c) and

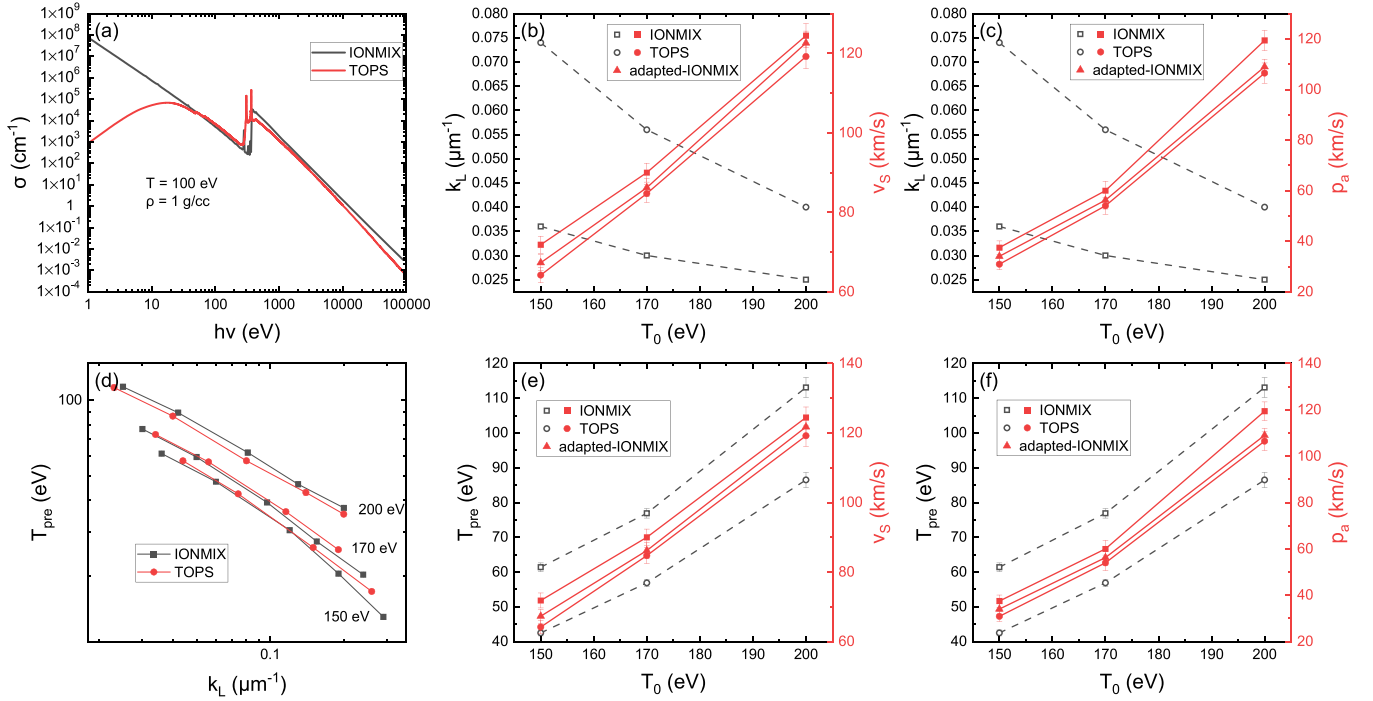


FIG. 4. (a) Spectrum of opacities varying with different photon energies calculated by IONMIX code and from TOPS database at density  $\rho = 1 \text{ g/cm}^3$  and electron temperature  $T = 100 \text{ eV}$ . (b)  $k_L$  and shock velocity  $v_s$ , and (c)  $k_L$  and ablation pressure  $p_a$  at different drive radiation temperatures. (d) Preheating temperatures  $T_{\text{pre}}$  at different  $k_L$ . (e)  $T_{\text{pre}}$  and shock velocity  $v_s$ , and (f)  $T_{\text{pre}}$  and ablation pressure  $p_a$  at different drive radiation temperatures. IONMIX and TOPS in the legends in panels (b), (c), (e), and (f) refer to the RHD simulation results from the original IONMIX and TOPS opacities, respectively. Adapted-IONMIX in the legends refers to the result when using opacity multiplier  $\zeta$  to adapt the IONMIX opacity to get the same  $k_L$  or  $T_{\text{pre}}$  as those in simulations with original TOPS opacity.

(7d) as well. For comparison, the dashed lines in Figs. 3(b) and 3(c) plot the respective results by using the LTE model [13], where only the Rosseland opacity multiplier is used, which shows that both  $v_s$  and  $p_a$  depend heavily and unphysically on the artificial opacity multiplier as  $v_s \propto \zeta^{-0.25}$  and  $p_a \propto \zeta^{-0.5}$  [13].

For further comparison and verification of our proposal of using a simple single multiplier  $\zeta$ , Figs. 3(d)–3(f) plot the RHD simulation results varying with different drive radiation temperatures  $T_0$  by using separate opacity multipliers  $\zeta_P$  and  $\zeta_R$ , i.e.,  $\tilde{\sigma}_P = \zeta_P \sigma_P$  and  $\tilde{\sigma}_R = \zeta_R \sigma_R$  [31,33], where the solid lines represent those for the cases  $\zeta_P = \zeta_R$ , while the dashed lines for the cases  $\zeta_P \neq \zeta_R$ . On one hand, we see that when the drive radiation temperature  $T_0$  increases, the radiation temperature damp scale length  $k_L$  decreases [see Fig. 3(d)] because both  $\sigma_P$  and  $\sigma_R$  decrease with increasing  $T_0$  and  $k_L \propto \sqrt{\sigma_P \sigma_R}$  [Eq. (7a)]. The shock velocity  $v_s$  [see Fig. 3(e)] and the ablation pressure  $p_a$  [see Fig. 3(f)] increase, in consistence with the theoretical equations (7b) and (7c). On the other hand, as expected, the values of  $v_s$  and  $p_a$  vary very slightly between the cases of  $\zeta_P = \zeta_R = 1$  and  $\zeta_P = \zeta_R = 10$  [compare the black and red solid lines in Figs. 3(e) and 3(f)]; however, they show big differences between the cases of  $(\zeta_P = 10, \zeta_R = 1)$  and  $(\zeta_P = 1, \zeta_R = 10)$  [compare the blue and green dashed lines in Figs. 3(e) and 3(f)], actually depending on  $\zeta_P/\zeta_R$  [agreeing with Eqs. (7b) and (7c)].

To demonstrate the application of our non-LTE model on how to apply a proper opacity multiplier by using  $k_L$  as a benchmark, we compare the RHD simulation results by using

two different sets of opacities from, respectively, the IONMIX code [31] and the TOPS database [33,34]. All other simulation setup and parameters are the same as above. A single multiplier  $\zeta$  is used on both opacities as  $\tilde{\sigma}_P = \zeta \sigma_P$  and  $\tilde{\sigma}_R = \zeta \sigma_R$ . Figure 4(a) plots the spectrum of opacities varying with different photon energies from, respectively, the IONMIX code and the TOPS database, where the plasma density  $\rho = 1 \text{ g/cm}^3$  and electron temperature  $T = 100 \text{ eV}$  are chosen as typical parameters in the radiative ablation process of ICF capsule. We see that the significant difference between them occurs at photon energies in the range less than 40 eV and the location of K edge also shifts. Such difference at the low-photon energy range comes from the difficulty of calculating accurate atomic energy level [31,33,34]. As a result, obviously the typical physical quantities,  $k_L$ ,  $v_s$ ,  $p_a$ , and so on, obtained from the RHD simulations should also be much different for two cases, which can be seen by comparing the dashed black lines between symbols of “hollow square” and “hollow circle dot” in either Figs. 4(b) or 4(c) for  $k_L$ , the solid red lines with between symbols of “solid square” and “solid circle dot” in Fig. 4(b) for  $v_s$ , and the solid red lines between symbols of “solid square” and “solid circle dot” in Fig. 4(c).

As suggested by our theoretical model, to match the two simulation results, for benchmark, we apply a multiplier  $\zeta$  for the opacities in the simulation with the IONMIX calculation so that the values of  $k_L$  become the same as those obtained in the simulation with the TOPS database. The red solid lines with the “triangle” symbols in both Figs. 4(b) and 4(c) clearly show that, no matter what the exact value of the opacity multiplier

$\zeta$  is, the adapted  $v_s$  and  $p_a$  obtained in the simulation with the IONMIX code become almost the same as those obtained in the simulations with the TOPS database. The difference is within 5% for  $v_s$  and 10% for  $p_a$  at the same  $k_L$  between the two opacities. This proves that  $k_L$  is indeed a proper benchmark to adapt the opacities in RHD simulations. Besides, we can also see from Figs. 4(b) and 4(c) that when  $k_L$  increases, the width of the absorption zone decreases, and resultantly both  $v_s$  and  $p_a$  decrease, which further verifies the validity of our model.

Finally, note that  $k_L$  is actually hard to be directly measured in the experiments, while the radiation temperature behind the target, namely, the preheating radiation temperature  $T_{\text{pre}}$  [35–38], can be taken as an effective benchmark to match the simulation results with experimental measurements [19,39,40]. The reason for this is simple as  $k_L$  characterizes the ability of the target to block radiation flux and obviously  $T_{\text{pre}}$  is determined directly by  $k_L$ , which can be seen clearly in Fig. 4(d), that are obtained from the above simulations. To further verify the validity of  $T_{\text{pre}}$  as a proper benchmark, Figs. 4(e) and 4(f) plot  $v_s$ ,  $p_a$ , and  $T_{\text{pre}}$  at time  $t = 4$  ns obtained from RHD simulations in different situations for the drive radiation temperatures of 150, 170, and 200 eV, respectively, where various  $\zeta$  are applied to match the obtained  $T_{\text{pre}}$  almost the same. Similarly, the difference is within 5% for  $v_s$  and 10% for  $p_a$  at the same  $T_{\text{pre}}$  between the two opacities.

#### IV. SUMMARY AND DISCUSSION

In summary, we have proposed a non-LTE subsonic radiative ablation model where both the electron-radiation energy exchange and the plasma fluid motion are self-consistently included. After derivations, we find that different from previous models where only the Rosseland opacity  $\sigma_R$  functions, in our non-LTE model, besides  $\sigma_R$ , the Planck opacity  $\sigma_P$  plays an important role as well. Further analyses show that most of the characteristic physical quantities of the non-LTE ablation, such as the ablation velocity  $v_a$ , ablation pressure  $P_a$ , and shock velocity  $v_s$ , depend on a combination of  $\sigma_P$  and  $\sigma_R$  together as  $\sigma_P/\sigma_R$ , while the spatial damping scale length of the radiation temperature scales as  $L \propto (\sigma_P\sigma_R)^{1/2}$ . Therefore, to match the RHD simulation results with the experimental measurements for the characteristic quantities  $v_a$ ,  $P_a$ ,  $v_s$ , etc., we can use a simple multiplier  $\zeta$  to adjust both the Planck and Rosseland opacities as  $\tilde{\sigma}_P = \zeta\sigma_P$  and  $\tilde{\sigma}_R = \zeta\sigma_R$ , where  $L$ , or more practically, the preheating radiation temperature  $T_{\text{pre}}$ , is chosen as a benchmark to remain the same between the simulation and experimental results. The 1D RHD simulations have been carried out, verifying all the theories well.

Some limitations still exist with our non-LTE subsonic radiative ablation model. First, besides the stationary zone, a very subsonic approximation of  $M \ll 1$  [13] is applied, which may not be fully satisfied, and the power law of electron temperature in the function of opacity may be quantitatively inaccurate. Second, it is actually difficult to get an accurate value of  $k_L$  in both RHD simulations and experiments when taking it as a benchmark for comparing between them. Third, some ignored physical processes in the model may also affect the simulation results such as multigroup radiation diffusion [41,42], electron thermal transport [30], and dependence of opacities on time and space [23,24].

Despite the above limitations, the proposed non-LTE radiative ablation model has vital potential applications in not only inertial confinement fusion but also other high-energy density physics fields such as astrophysics. First, in indirect-drive ICF, the proposed non-LTE radiative ablation model can be used to explain the dynamics of the capsule implosion ablated by the X-ray radiation, where the capsule material is generally of low  $Z$ , such as typically CH. Specifically, as mentioned in the manuscript, our proposed model can be used to match the RHD simulation and experimental results. Besides, several experimental diagnostic methods in indirect-drive ICF, such as the wall loss [9,43] and the reemitted flux [40] measurements that are originally based on the LTE model, need to be carefully modified for low- $Z$  material ablation according to our non-LTE model. Second, it has been proved that the non-LTE radiation transfer makes significant influence in Z-pinch [44] because the high-density, high-temperature plasmas are not in thermal equilibrium, and radiation strongly influences their energy balance, ionization state, and spectral properties. Our proposed non-LTE model can be used for studies of plasma states in Z-pinch. Finally, non-LTE radiation transfer can also be applicable in astrophysics. Radiation transfer is an essential process in multiple astronomical situations, in which Planck and Rosseland opacities both have effects. For example, different sets of opacities significantly affect the calculated hydrodynamical structure of accretion disk [45], and different sets of opacities influence the stellar model including stellar structure and evolution [46], especially the solar model [47]. Furthermore, in recent studies, non-LTE radiation transfer has been proven to be necessary under certain conditions [48,49], for example, modeling the accretion shock structures [49]. In this case, considering the non-LTE effects is important in certain astronomical situations and our non-LTE model can also be extended to these systems.

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#### DATA AVAILABILITY

The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

#### APPENDIX

In this Appendix, we provide the analysis of the requirements for parameters to get self-similar solution of Eq. (1) using the method in the literature [18]. Requirements for parameters under supersonic condition are derived in the literature [18], and the requirements under subsonic condition are more strict according to the following results. Substitute  $E = aT_r^4$ ,  $F = -(c/3\sigma_R)\partial E/\partial x$ ,  $\sigma_P = k'_0\rho T_e^{-\alpha'}$ , and  $\sigma_R = k_0\rho T_e^{-\alpha}$

into Eq. (1). Considering  $e = c_v T_e^{4\beta}$ ,  $p = 2\rho e/3$ , and setting  $U = aT_e^4$ , we have

$$\frac{\partial \rho}{\partial t} + \frac{\partial \rho u}{\partial x} = 0, \quad (\text{A1a})$$

$$\rho \left( \frac{\partial u}{\partial t} + u \frac{\partial u}{\partial x} \right) + \frac{2}{3} \frac{c_v}{a^\beta} \frac{\partial \rho U^\beta}{\partial x} = 0, \quad (\text{A1b})$$

$$\rho \frac{c_v}{a^\beta} \left( \frac{\partial U^\beta}{\partial t} + U^\beta \frac{\partial u}{\partial x} \right) = ck'_0 \rho a^{\frac{\alpha'}{4}} U^{-\frac{\alpha'}{4}} (E - U), \quad (\text{A1c})$$

$$\frac{\partial E}{\partial t} - \frac{\partial}{\partial x} \frac{c}{3k_0 \rho a^{\frac{\alpha}{4}}} U^{\frac{\alpha}{4}} \frac{\partial E}{\partial x} = ck'_0 \rho a^{\frac{\alpha'}{4}} U^{-\frac{\alpha'}{4}} (U - E). \quad (\text{A1d})$$

To get self-similar solution of the equations, we consider the boundary condition  $E(\infty) = E_0 t^{4\tau}$ . Three basic dimensions  $[L]$ ,  $[t]$ ,  $[E]$  exist in this problem, denoting the dimensions of length, time, and energy per unit volume, respectively. Dimensional variables in this problem include space and time  $x, t$ , four unknown variables  $\rho, u, E, U$ , and four parameters  $A = c_v/a^\beta$ ,  $B = c/(3k_0 a^{\alpha/4})$ ,  $C = ck'_0 a^{\alpha'/4}$ , and  $E_0$ . According to Pi theorem of dimensional analysis [50], the solution can be expressed by  $2 + 4 + 4 - 3 = 7$  dimensionless variables. If self-similar solution exists, one of them is self-similar variable and four of them are dimensionless functions used for expressing  $\rho, u, E, U$ . So, there are  $7 - 5 = 2$  dimensionless variables expressed by other parameters; i.e.,  $\Pi_i = A^{x_1} B^{x_2} C^{x_3} E_0^{x_4}$  has two independent solutions  $i = 1, 2$ . Dimensions of these parameters are listed in Table I.

TABLE I. Parameters and their dimensions.  $[L]$ ,  $[t]$ , and  $[E]$  denote the dimensions of length, time, and energy per unit volume, respectively.

| A                               | B                                      | C                                               | $E_0$                     |
|---------------------------------|----------------------------------------|-------------------------------------------------|---------------------------|
| $\frac{[L]^2}{[E]^\beta [t]^2}$ | $\frac{[t]}{[E]^{\frac{\alpha}{4}-1}}$ | $\frac{[E]^{\frac{\alpha'}{4}-1} [L]^2}{[t]^3}$ | $\frac{[E]}{[t]^{4\tau}}$ |

The equation turns into

$$2x_1 + 2x_3 = 0, \quad (\text{A2a})$$

$$-2x_1 + x_2 - 3x_3 - 4\tau x_4 = 0, \quad (\text{A2b})$$

$$-\beta x_1 + \left(1 - \frac{\alpha}{4}\right)x_2 + \left(\frac{\alpha'}{4} - 1\right)x_3 + x_4 = 0. \quad (\text{A2c})$$

Augmented matrix of the system of linear equations is

$$\begin{pmatrix} 2 & 0 & 2 & 0 \\ -2 & 1 & -3 & -4\tau \\ -\beta & 1 - \frac{\alpha}{4} & \frac{\alpha'}{4} - 1 & 1 \end{pmatrix}.$$

It can be reduced to

$$\begin{pmatrix} 1 & 0 & 1 & 0 \\ 0 & 1 & -1 & -4\tau \\ 0 & 0 & \beta + \frac{\alpha'}{4} - \frac{\alpha}{4} & 1 + 4\tau(1 - \frac{\alpha}{4}) \end{pmatrix}.$$

If there are two independent solutions for the system of linear equations, the last row of reduced augmented matrix must be zero. So, relations  $\alpha - \alpha' = 4\beta$  and  $\alpha'\tau = 1 + 4(1 - \beta)\tau$  must hold. The strict requirements are not always satisfied, so self-similar solution does not always exist.

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