Heat waves driven by thermal radiation in tamped flows

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(Received 31 May 1991)

The heating of matter driven by thermal radiation is studied by means of a simple model, under conditions where the material cannot freely expand. Supersonic and ablative regimes are identified when the radiation field is not in local thermodynamic equilibrium with the matter and, in a later phase, when equilibrium is reached. The scaling laws characterizing the magnitudes of the heated material are found for the different regimes. Application to an ion-beam-fusion target is examined and the effect of an active tamper that pushes the ablated material is analyzed.

PACS number(s): 52.50.Gj, 47.70.Mc, 44.40.+a

I. INTRODUCTION

Heat waves driven by thermal radiation are an important issue in ion-beam-driven inertial-confinement fusion. They develop in the interaction of the ion beam with a relatively thick spherical-shell target when the thermal radiation generated in the ion-beam-heated region impinges on the cold parts, which are optically thick to the radiation, carrying energy to regions not accessible to the ions [1–12]. The heat wave may be initially supersonic or ablative depending on the time evolution of the incident radiation flux [3]. In the first case, the supersonic heat wave (SHW) develops into an ablative heat wave (AHW) as the radiation-heated material expands and drives a shock wave which, after a certain time, catches up with the heat front and overtakes it [11,13]. In addition, when the medium starts to be heated by the radiation, its temperature is relatively low and it is in nonlocal thermodynamic equilibrium (NLTE) with the radiation field (although local thermodynamic equilibrium may be assumed for the material). In this case, the characteristic length over which the radiation flux is absorbed is given by the photon mean free path [13]. In a later phase, the heated material becomes hot enough to emit photons itself and the medium reaches local thermodynamic equilibrium (LTE) with the radiation. By this time, the photon mean free path becomes shorter than the characteristic absorption distance of the radiation flux, which is now determined by the characteristic scale of the temperature gradient. In this regime, the heat-conduction approximation is valid [13]. Thus four possible regimes may be identified during the heating process: a supersonic heat wave and an ablative heat wave during the NLTE phase (SHW-NLTE and AHW-NLTE), and a supersonic and an ablative heat wave during the LTE phase (SHW-LTE and AHW-LTE). Several of these modes may happen when the radiation heats a cold medium depending on the radiation flux S, its time evolution, and the density ρ_0 and the properties of the material. Moreover, since in the ablation regime the motion of matter is an essential feature, the characteristics of such regimes will also be dependent on the boundary conditions affecting the flow behind the heat front. For example, if the radiation penetrates the material through its interface with the vacuum [Fig. 1(a)], the heated material can freely expand, and it will fill the vacuum with a velocity of the order of

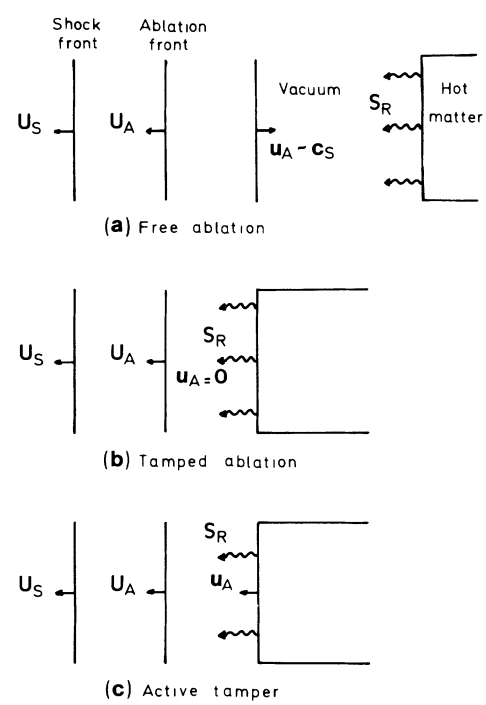

FIG. 1. Ablation driven by thermal radiation emitted by hot matter. (a) The emitter is far from the cold medium and the ablated flow is free to expand into the vacuum. (b) The emitter is in contact with the cold medium and it tamps the expansion of the ablated flow (c) The emitter is in contact with the cold medium and it expands, pushing the ablated mass towards the left.

the sound speed c_s. This is the case when the hot medium emitting the radiation is relatively far from the cold material. Such a situation has been extensively examined in the recent literature in relation to the indirect driving of a fusion capsule [6,11, 14–21]. On the other hand, if the material medium emitting the radiation is in contact with the cold medium, it may tamp the expansion of the radiation-heated matter or, furthermore, the emitter can itself expand, pushing the ablated matter [Figs. 1(b) and 1(c)]. In this case the presence of the emitter imposes boundary conditions on the ablated flow that affect the properties of the heating process.

In a typical target for ion-beam fusion, the situation shown in Figs. 1(b) and 1(c) may be present [1–3,5,10]. The basic process of the implosion of such targets lies on the fact that, during the interaction, the ion beam penetrates deeply into the thick wall of the target. The internal layer, the payload, is accelerated by the thermal pressure generated in the external layers, the absorber, where the beam energy is completely deposited (Fig. 2) [3,22,23]. The internal part of the absorber, the converter, is heated to a few hundred eV, resulting in an optically thick medium that emits part of the absorbed energy as blackbody radiation. The external part of the absorber, the tamper, is constructed of a high-z, high-density material that tamps the expansion of the converter and prevents the radiation loss. The radiation is absorbed and reemitted by the tamper and the payload and the converter becomes a closed cavity filled with equilibrium radiation [3,4]. The thermal pressure in the converter tamps the material ablated from the payload and, at later times, the converter expands, pushing the payload, so that the ablation process takes place under tamped conditions such as in the situations described in Figs. 1(b) and 1(c).

Although radiation-driven ablation in tamped flows is clearly of interest to ion-beam fusion, not much research work has been so far devoted to this subject. Indeed, theoretical work was performed in connection with experimental studies in which heat waves were generated using a thermal shock tube [24–26]. In addition, the scaling laws for the ablation driven by thermal radiation

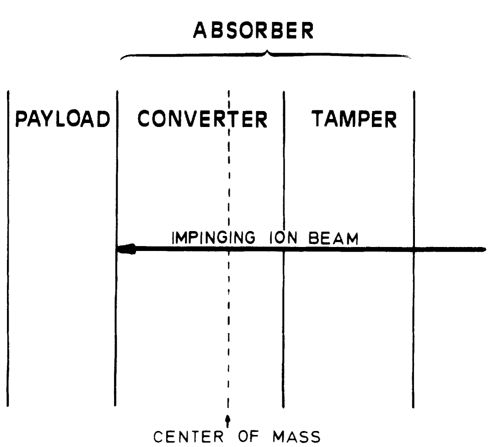

FIG. 2. Schematic diagram of the beam-foil geometry for a tampered target for ion-beam fusion.

with perfectly tamped flow ($u_A = 0$) were recently obtained by means of simple analytical models for the case of a totally ionized medium that was not in equilibrium with the radiation [27].

In this paper the previous scaling laws are generalized to LTE conditions. The proton mean free path l in a multiply ionized heated medium is in general a complicated function of both temperature and density. In particular, l as a function of the temperature T reaches a minimum at relatively low temperatures [9,13]. This minimum is in the region where the double ionization begins. For example, for carbon it occurs at $T = T_m \simeq 1-2$ eV [9]. For higher temperatures ($T > T_m$) the mean free path increases with temperature [13], and it can be considered in the usual way by assuming l in the form of a power law [1,8,9,11–18]:

$$l = l_0 \varepsilon^n / \rho^m , \tag{1}$$

where ε is the specific internal energy, ρ is the mass density, and the positive constants n and m ($n, m > 0$), and the physical parameter l_0 describe a given material. The equation of state relating the specific internal energy ε with the temperature T is assumed to be [1,8–18]:

$$\varepsilon = \varepsilon_s T^\alpha , \tag{2}$$

where the positive constant α and the physical parameter ε_s are characteristic of the material. As is well known, if the pressure p is given by the relationship

$$p = (\gamma - 1)\rho\varepsilon , \tag{3}$$

where γ is the enthalpy coefficient, thermodynamic consistency requires a dependence on density in Eq. (2): $\varepsilon \sim T^\alpha \rho^{-\beta}$ where $\beta = (\gamma - 1)(\alpha - 1)$ [9,13]. However, in most cases of interest $\beta \ll 1$ and here it is taken $\beta = 0$.

With the previous equations and assuming a self-regulating mechanism for the radiation optical depth [11,27–33], the properties on the tamped ablation driven by thermal radiation are studied. The pressure and the velocity of the ablation front are determined only by the radiation flux S and the matter density ρ_0, regardless of the heating physical mechanism. Besides, the temperature of the ablated material is, in general, higher than that which results in free-flow ablation. This effect is increased when the expansion of the converter is considered ($u_A \neq 0$).

The supersonic regimes are reviewed in Sec. II, the scaling laws are written, and the characteristic times required to achieve the equilibrium and to develop into an AHW are calculated. The ablative regimes under conditions of perfect taming ($u_A = 0$) are examined in Sec. III and the characteristic equilibrium time is calculated. The evolution of the heating process is analyzed in Sec. IV. There exists of a critical value of the flux for which the transition from the SHW to the AHW occurs just at the equilibrium time. The effect of the expansion of the converter is analyzed in Sec. V, and the main conclusions are summarized in Sec. VI.

II. SUPERSONIC HEAT WAVES

In order to study the supersonic regime, the heating process can be considered to take place without an appreciable change in the density, so that the density of the heated material is $\rho \simeq \rho_0$. In addition to this, a high-energy flux $S(t)$ is assumed to be absorbed by cold matter of state "0" (Fig. 3), and thermalized so that a relaxed new state behind the absorption region with high temperature is obtained [25]. The fluid parameters behind and ahead of the transition zone are assumed to be uniform and the power input $S(t)$ is considered to be practically constant during the transition time of a particle through the transition zone. With these assumptions the heat wave can be described by means of the conservation equations for mass, momentum, and energy in one-dimensional, integrated form. In the heat-front frame of reference, they are

$$\rho_0 v_0 = \rho v \ , \tag{4a}$$

$$p_0 + \rho_0 v_0^2 = p + \rho v^2 \ , \tag{4b}$$

$$v_0^2/2 + \gamma_0 \varepsilon_0 + S/\rho_0 v_0 = v^2/2 + \gamma \varepsilon \ . \tag{4c}$$

Since in the supersonic regimes $\rho \simeq \rho_0$, and besides $\varepsilon \gg \varepsilon_0$,

$$S \simeq \rho v \varepsilon \simeq \rho_0 v \varepsilon \tag{4d}$$

is valid, where, for simplicity, the numerical factor dependent on γ was dropped and only the essential scaling was retained. If the medium ahead of the heat wave is at rest in the laboratory frame of reference, then $v_0 \simeq v$ is the characteristic heat-front velocity.

On the other hand, a power law is assumed for the time variation of the radiation flux:

$$S = S_0 (t/\tau)^q \ , \tag{5}$$

where S_0 is a constant, and τ is a characteristic time which is longer than the time required for a particle of fluid to go through the transition region.

For the purpose of getting the scaling laws relating v and ε with S and ρ_0, a second equation, which must take into account the physical process of heating, is necessary.

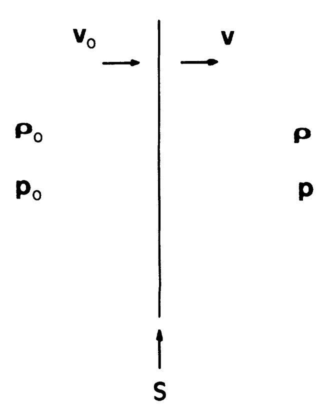

FIG. 3. Heat front induced by the absorption of a high flux S by matter of state "0."

Such a relation is found by assuming a self-regulating mechanism for the opacity of the heated region, which sets that, at any time t, the thickness of this region is given by the characteristic length l_c over which the radiation is absorbed

$$\int_0^t v(t')dt' = l_c(t) \ . \tag{6a}$$

With the assumed form for the radiation flux $S(t)$, given by Eq. (5), and considering Eqs. (1) and (4d), $v(t)$ can be written as $v(t) \sim (t/\tau)^g$. Therefore, dropping the numerical factor dependent on g, Eq. (6a) yields

$$vt \simeq l_c \ . \tag{6b}$$

The self-regulating hypothesis has proved to be true in many different situations, i.e., matter heated by a laser [28,29], by electronic thermal conduction [30] by fast electrons [31,32] and light-ion beams [33], and by thermal radiation [11,15,19,20]. In any case, the particular mechanism of energy transport provides the adequate expression for the characteristic length l_c. In the present problem l_c changes as equilibrium is approached and the radiation transport can then be described using the thermal-conduction approximation.

A. SHW-NLTE

At the beginning of the process, the heating progresses under conditions in which the matter is not in equilibrium with the radiation. Therefore the thickness of the heated region is determined by the photon mean free path l given by Eq. (1). When using Eq. (1), it should be kept in mind that in the NLTE regime the temperature T of the media is lower than the radiation temperature, and it must also be $T > T_m$, where, as it was previously mentioned, T_m is the temperature for which the mean free path l is a minimum. Actually it is not a strong limitation for the model because such a minimum occurs in the low-temperature region and it can be expected that the temperature T_m will be achieved in a time much shorter than the characteristic time for which the NLTE regime exits. Thus, from Eqs. (1), (4d), and (6b), the following scaling relations are found:

$$\varepsilon \simeq (St/l_0)^{1/(n+1)} \rho_0^{(m-1)/(n+1)} \ , \tag{7}$$

$$v \simeq (l_0/t)^{1/(n+1)} S^{n/(n+1)} \rho_0^{-(m+n)/(n+1)} \tag{8}$$

where $S = S(t)$ is given by Eq. (5). These expressions were recently reported in Ref. [11] and, of course, since the hydrodynamic motion is not relevant in the supersonic regimes, they are the same for tamped or free-flow heat waves.

The SHW-NLTE exists while the front velocity v is higher than the sound speed $c_s \simeq \varepsilon^{1/2}$. When $v \simeq \varepsilon^{1/2}$, it develops into an AHW-NLTE. Actually, such a wave can exist earlier than the SHW-NLTE depending on the time evolution of the radiation flux [13]. Assuming it in the form given by Eq. (5) the characteristic time t_{TN} for the transition from the supersonic to the ablative regime can be calculated from Eqs. (6)–(8):

$$t_{TN}=\tau(l_0/\tau)^{3/[3-q(2n-1)]}S_0^{(2n-1)/[3-q(2n-1)]}\times\rho_0^{-(3m+2n-1)/[3-q(2n-1)]} . \quad (9)$$

It is obvious that a SHW-NLTE will exist previously to the AHW-NLTE only if $q<3/(2n-1)$. Otherwise, the AHW-NLTE is generated from the very beginning and at the time t_{TN} it develops into a SHW-NLTE [13]. On the other hand, it may happen that, for $q<3/(2n-1)$, LTE conditions are reached at the time $t_{ES}<t_{TN}$ and so, the SHW-NLTE develops into a SHW-LTE. The characteristic time t_{ES} is calculated by considering that, when the equilibrium is settled, $S=\sigma T^4$ (where σ is the Stefan-Boltzmann constant). Then, from Eqs. (2), (5), and (6b) it is found that

$$t_{ES}=\tau(l_0/\tau)^{4/\{4-q[\alpha(n+1)-4]\}}\times(\varepsilon_s^{4/\alpha}/\sigma)^{\alpha(n+1)/\{4-q[\alpha(n+1)-4]\}}\times S_0^{[\alpha(n+1)-4]/\{4-q[\alpha(n+1)-4]\}}\times\rho_0^{4(m-1)/\{4-q[\alpha(n+1)-4]\}} . \quad (10)$$

B. SHW-LTE

If the equilibrium is achieved before the transition to an AHW-NLTE, the heating process is governed by a SHW-LTE. As it will be shown later, such a wave can also be reached from an AHW-LTE provided that sufficient time is available.

Under LTE conditions, the thermal-conduction approximation is valid and the radiation flux is proportional to the temperature gradient [13]:

$$S=-\tfrac{16}{3}l\sigma T^3\nabla T . \quad (11)$$

The characteristic length to be used in the self-regulating condition expressed by Eq. (6b) is given now by the scale of the temperature gradient:

$$l_T=T/|\nabla T|\simeq l\sigma T^4/S . \quad (12)$$

By combining Eq. (12) with Eqs. (2) and (4d), the following expression for l_T is obtained:

$$l_T\simeq(l_0\sigma/\varepsilon_s^{4/\alpha})\varepsilon^{[8+\alpha(2n-3)]/2\alpha}\rho_0^{-(m+1)} . \quad (13)$$

Then, introducing Eq. (13) into Eq. (6b) and using Eq. (4d), the scaling laws for the SHW-LTE result as follows:

$$\varepsilon=(t\varepsilon_s^{4/\alpha}/l_0\sigma)^{\alpha/[4+\alpha(n+1)]}\times S^{2\alpha/[4+\alpha(n+1)]}\rho_0^{(m-1)\alpha/[4+\alpha(n+1)]} , \quad (14)$$

$$v=(l_0\sigma/t\varepsilon_s^{4/\alpha})^{\alpha/[4+\alpha(n+1)]}\times S^{[4+\alpha(n-1)]/[4+\alpha(n+1)]}\times\rho_0^{-[4+\alpha(m+n)]/[4+\alpha(n+1)]} , \quad (15)$$

where $S=S(t)$ is given by Eq. (5). Equation (14) is equivalent [through Eq. (2)] to the scaling law derived in Ref. [17] by dimensional analysis.

As in the previous case, the SHW-LTE will exist if $v>\varepsilon^{1/2}$ and the characteristic times t_{TE} for the transition to an AHW-LTE can be found from Eqs. (5), (14), and (15):

$$t_{TE}=\tau(l_0\sigma/\tau\varepsilon_s^{4/\alpha})^{3\alpha/\{3\alpha-2q[4+\alpha(n-2)]\}}\times S_0^{2[4+\alpha(n-2)]/\{3\alpha-2q[4+\alpha(n-2)]\}}\times\rho_0^{-[8+\alpha(3m+2n-1)]/\{3\alpha-2q[4+\alpha(n-2)]\}} . \quad (16)$$

III. ABLATIVE REGIMES

The ablative regimes are characterized by the presence of a shock wave running ahead of the ablation front. It was previously shown that such a system can be analyzed by means of a three-phase model as the one represented in Fig. 4 [25,27,34]. The ablation front is moving towards the left with velocity U_A and it is preceded by a shock wave which is penetrating in the unperturbed phase with velocity U_s. The particle velocities in the shock wave frame of reference are v_0 and v_1 ahead and behind, respectively. In the ablation front frame of reference they are v_a and v, respectively. The velocities in the laboratory frame are u_c, u_s, and u_A for the unperturbed, shocked, and ablated phases, respectively. The following relations can be written:

$$U_s=u_c-v_0=u_s-v_1 , \quad (17)$$

$$U_A=u_s-v_a=u_A-v . \quad (18)$$

The shock wave generates the initial conditions for the AHW. If the unperturbed region is at rest, $u_c=0$ and, by considering a strong shock, the Rankine-Hugoniot equations yield

$$p_a=[2/(\gamma_a+1)]\rho_0 v_0^2 , \quad (19)$$

$$\rho_a/\rho_0=v_0/v_1=(\gamma_a+1)/(\gamma_a-1) . \quad (20)$$

$$U_s=-v_0 , \quad (21)$$

where p_a, ρ_a, and γ_a are the pressure, density, and enthalpy coefficient of the shocked phase.

For the ablation front, the conservation equations of

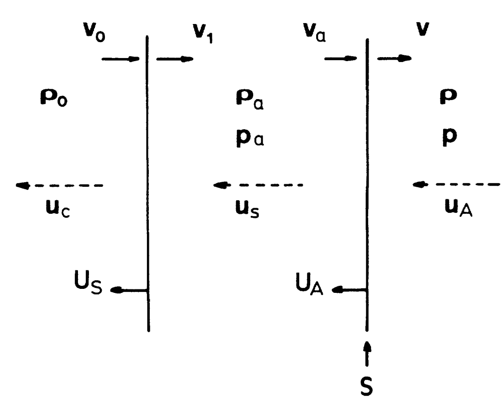

FIG. 4. Three-phase ablation model. The ablative heat wave is a weak expansion front with intense energy deposition and a preceding shock wave.

mass, momentum, and energy are

$$\rho_a v_a = \rho v \ , \tag{22}$$

$$p_a + \rho_a v_a^2 = p + \rho v^2 \ , \tag{23}$$

$$v_a^2/2 + \gamma_a \varepsilon_a + S/\rho_a v_a = v^2/2 + \gamma\varepsilon \ , \tag{24}$$

where ε_a and ε are the specific internal energies of the shocked and the ablated phases, respectively, and p, ρ, and γ are the pressure, density, and enthalpy coefficients of the ablated phase, respectively, as defined in Sec. I.

Since the AHW is an expansion front, the density will decrease across the ablation front. That is, $\rho \gg \rho_a$ ($v \gg v_a$). Besides, a large amount of energy is considered to be deposited in the ablated region, so that $S/\rho_a v_a \gg v_a^2, \varepsilon_a$. On the other hand, the pressure p is expected to remain relatively high during the process (the AHW is a weak expansion front), and so $p_a \simeq p$. With these assumptions, Eqs. (22)–(24) are easily solved. At a first step, Eqs. (22) and (23) are combined:

$$p_a = p + \rho v^2(1 - \rho/\rho_a) \ . \tag{25}$$

Because $\rho \ll \rho_a$ and $p_a \simeq p$ were assumed, Eq. (25) implies that $\rho v^2 \ll p$, or $v^2 \ll \varepsilon$. This is just the condition which characterizes a weak expansion front [25]. With the previous assumptions, Eq. (24) leads to the following form of energy conservation equation:

$$\gamma\varepsilon \simeq S/\rho_a v_a = S/\rho v \ . \tag{26}$$

From Eqs. (19) and (26) the ablation velocity v is found ($p \simeq p_a$):

$$v = [(\gamma-1)(\gamma+1)/2\gamma](S/\rho_0 v_0^2) \ . \tag{27}$$

Moreover, taking into account that $v \gg v_a$ and $u_c = 0$, another expression for v results from Eqs. (17) and (18):

$$v \simeq u_A - u_s \simeq u_A + [2/(\gamma_a+1)]v_0 \ . \tag{28}$$

Eliminating v from Eqs. (27) and (28),

$$v_0^2 + [(\gamma_a+1)/2]u_A v_0^2 = [(\gamma_a+1)(\gamma-1)/2\gamma](S/\rho_0) \ . \tag{29}$$

This equation takes into account the state of the fluid behind the ablation front and it cannot be obtained by dimensional analysis. For the particular case of free-flow ablation, $u_A \simeq \varepsilon^{1/2} \gg v_0$, u_s and Eqs. (28) and (29) lead to $v \simeq \varepsilon^{1/2}$ [Fig. 1(a)]. For tamped flow, however, $u_A = 0$ if the tamping is perfect [Fig. 1(b)], and $u_A < 0$ if the tamping is active [Fig. 1(c)]. Here, $u_A = 0$ will be considered and later, in Sec. V, the effect of active tamper resulting from the expansion of the converter region in an ion-beam-fusion target will be examined. By putting $u_A = 0$ in Eq. (29), the ablation pressure p and the ablation velocity v result:

$$p = [(\gamma-1)^2(\gamma_a+1)/2\gamma]^{1/3}\rho_0^{1/3}S^{2/3} \ , \tag{30}$$

$$v = [2(\gamma-1)/\gamma(\gamma+1)]^{1/3}(S/\rho_0)^{1/3} \ . \tag{31}$$

As can be seen, these expressions are obtained regardless of the particulars of the physical process of energy deposition in the ablated region. They are completely determined by the boundary conditions on the flow behind the ablation front ($u_A = 0$). However, in order to determine the density ρ of the ablated matter and the mass ablation rate $\dot m = \rho v$ (per unit of area), the mechanism of energy deposition must be specified through the self-regulating condition given by Eq. (6b). From Eqs. (3) and (30) it can be written

$$\rho \simeq [(\gamma_a+1)/2\gamma^2(\gamma-1)]^{1/3}\rho_0^{1/3}S^{2/3}/\varepsilon \sim \rho_0 v^2/\varepsilon \ . \tag{32}$$

A. AHW-NLTE

As was previously mentioned, the AHW-NLTE can exist from the beginning of the heating process, or it will exist after a time of the order of t_{EN} given by Eq. (9) if the process starts with a SHW-NLTE. By using the photon mean free path given by Eq. (1) in the self-regulating condition [Eq. (6b)], and combining it with Eq. (32), it is obtained that

$$\rho \simeq (l_0/t)^{1/(m+n)}S^{(2n-1)/[3(m+n)]} \times \rho_0^{(n+1)/[3(m+n)]} \ , \tag{33}$$

$$\varepsilon \simeq (t/l_0)^{1/(m+n)}S^{(2m+1)/[3(m+n)]} \times \rho_0^{(m-1)/[3(m+n)]} \ , \tag{34}$$

$$\dot m \simeq (l_0/t)^{1/(m+n)}S^{(m+3n-1)/[3(m+n)]} \times \rho_0^{-(m-1)/[3(m+n)]} \ , \tag{35}$$

where, for simplicity, the numerical factors depending on γ_a and γ in Eqs. (30)–(32) were dropped and only the essential scalings were retained.

According to Eq. (34), after a certain time the ablated material will be so hot that the flux $\sigma T^4 = \sigma(\varepsilon/\varepsilon_s)^{4/\alpha}$ emitted by the heated region becomes larger than the incident flux S. At this time, t_{EA}, the equilibrium is settled and the AHW-LTE is reached. Using Eqs. (2), (5), and (34) the time t_{EA} for which $\sigma T^4 = S$ is found:

$$\begin{gathered} t_{EA} = \tau(l_0/\tau)^{12/(12+qA)}(\varepsilon_s^{4/\alpha}/\sigma)^{3\alpha(m+n)/(12+qA)} \times S_0^{-A/(12+qA)}\rho_0^{-4(m-1)/(12+qA)} \ , \\ A = (8-3\alpha)(m+n) - 4(2n-1) \ . \end{gathered} \tag{36}$$

B. AHW-LTE

When equilibrium between the radiation and the heated matter is achieved, the characteristic length over which the radiation flux S is absorbed is given by Eq. (12). From Eq. (30), $S \simeq (\rho\varepsilon/\rho_0^{1/3})^{3/2}$, and using Eq. (2), the characteristic scale of the temperature gradient is written as

$$l_T \simeq (l_0\sigma\rho_0^{1/2}/\varepsilon_s^{4/\alpha})\varepsilon^{[8+\alpha(2n-3)]/2\alpha}\rho^{-(2m+3)/2} \ . \tag{37}$$

Then, as above, the scaling laws for ρ, ε, and $\dot m$ are obtained from Eqs. (6b), (32), and (37):

$$\rho \simeq (l_0\sigma/t\varepsilon_s^{4/\alpha})^{\alpha/[4+\alpha(m+n)]} \times S^{[8+2\alpha(n-2)]/3[4+\alpha(m+n)]} \times \rho_0^{[4+\alpha(n+1)]/3[4+\alpha(m+n)]} . \quad (38)$$

$$\varepsilon \simeq (t\varepsilon_s^{4/\alpha}/l_0\sigma)^{\alpha/[4+\alpha(m+n)]} \times S^{2\alpha(m+2)/3[4+\alpha(m+n)]} \times \rho_0^{\alpha(m-1)/3[4+\alpha(m+n)]} , \quad (39)$$

$$\dot m \simeq (l_0\sigma/t\varepsilon_s^{4/\alpha})^{\alpha/[4+\alpha(m+n)]} \times S^{[12+\alpha(m+3n-4)]/3[4+\alpha(m+n)]} \times \rho_0^{-\alpha(m-1)/3[4+\alpha(m+n)]} . \quad (40)$$

The AHW-LTE can be reached either from a SHW-LTE after a time of the order of t_{TE}, or from an AHW-NLTE after a time of the order of t_{EA}. As will be seen in the next section, it depends, for a given material, on the intensity S of the radiation.

IV. TIME EVOLUTION OF THE HEATING PROCESS

During the heating of a tamped cold medium driven by radiation, several of the four regimes previously treated will be traversed, provided that sufficient time is available. Besides, from Eqs. (9), (10), (16), and (36), which give the times required to reach either the equilibrium or the transition from a SHW to AHW, it is noted that there exists a critical value S_c of the radiation flux for which $t_{TN}=t_{ES}=t_{TE}=t_{EA}=\tau_c$. At that point, the transition from a SHW to an AHW, or vice versa, occurs simultaneously with the achievement of equilibrium:

$$S_c=(\tau/l_0)^q(\varepsilon_s^{4/\alpha}/\sigma)^{\alpha[3-q(2n-1)]/(8-3\alpha)} \times \rho_0^{8[3-q(2n-1)]/3(8-3\alpha)+q(3m+2n-1)/3} , \quad (41)$$

$$\tau_c=l_0(\varepsilon_s^{4/\alpha}/\sigma)^{\alpha(2n-1)/(8-3\alpha)} \times \rho_0^{8(2n-1)/3(8-3\alpha)-(3m+2n-1)/3} . \quad (42)$$

Introducing S_c and τ_c in Eqs. (9), (10), (16), and (36), a simpler expression for the transition times is found:

$$t_{TN}=\tau_c(S/S_c)^{(2n-1)/[3-q(2n-1)]} , \quad (43)$$

$$t_{ES}=\tau_c(S/S_c)^{[\alpha(n+1)-4]/\{4-q[\alpha(n+1)-4]\}} , \quad (44)$$

$$t_{TE}=\tau_c(S/S_c)^{2[4+\alpha(n-2)]/\{3\alpha-2q[4+\alpha(n-2)]\}} , \quad (45)$$

$$t_{EA}=\tau_c(S/S_c)^{-A/(12+qA)} . \quad (46)$$

These equations show that the space (S,t) is divided into four regions which correspond to the four possible regimes of heat waves presented in the previous sections. For the purpose of illustrating the possible time evolution of the heating process, the particular case in which the heated material is aluminum is considered here as a reference case. The corresponding diagram in the space (S,t) is shown in Fig. 5. The values for the relevant parameters were taken from Ref. [12], where they were obtained from fits to equation of state and opacity tables of the SESAME library and from Ref. [35]. They are as follows:

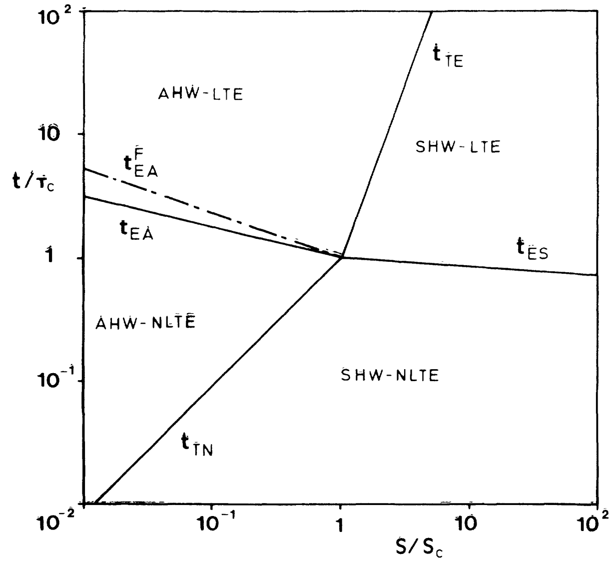

FIG. 5. Dimensionless times t_{TN}/τ_c and t_{TE}/τ_c for the transition from a SHW to an AHW for NLTE and LTE conditions, respectively, and dimensionless times t_{ES}/τ_c and t_{EA}/τ_c for the equilibrium in the supersonic and in the ablative regimes, respectively. For comparison, the dimensionless time t_{EA}^F/τ_c for the equilibrium in the ablative regime for free-flow ablation is shown.

$\alpha=\frac{6}{5}$, $m=\frac{3}{2}$, $n=\frac{25}{12}$, $\varepsilon_s=3.6\times10^{11}$ ergs/g eV$^{1.2}$, $l_0=8\times10^{-33}$ cm g$^{1.5}$/(erg/g)$^{25/12}$. Besides, $q=0$ was assumed, so that in the initial phase of the heating a SHW-NLTE is generated. If, however, $q>3/(2n-1)$ then the slope of the straight line labeled as t_{TN} becomes negative and the process starts with an AHW-NLTE. For the case shown in Fig. 5 $S_c=500\rho_0^{1.82}$ TW and $\tau_c=0.64\rho_0^{-0.64}$ ns (ρ_0 is in g/cm^3). If $S<S_c$, the SHW-LTE does not occur and the AHW-LTE is achieved after a phase where the AHW-NLTE takes place. If $S>S_c$, the intermediate phase is a SHW-LTE and the AHW-NLTE never happens. It is noted that this situation is not much different for free-flow heat waves. In fact, since the supersonic regimes are not affected by the boundary condition on the downstream flow, the transition times t_{TN}, t_{ES}, and t_{TE} are the same for both situations with free and tamped flows. The only difference is in the transition time t_{EA} from the AHW-NLTE to AHW-LTE. For comparison, the corresponding equilibrium time t_{EA}^F for free flow is shown in Fig. 5 with a dotted line. It can be seen that $t_{EA}^F>t_{EA}$. This is a general result and it is due to the fact that, in the tamped situation, a higher temperature is reached in the heated region. This is because, as a consequence of the tamping effect, less energy is drained as kinetic energy.

An expression for the critical flux S_c (with $q=0$) was previously found in relation with ion-beam energy conversion into blackbody radiation, and it was interpreted that below this critical value the heating is dominated by the hydrodynamic motion of the converter, resulting in a relatively low conversion efficiency, η_c [9,12]. The present results which apply to the radiation-heated region (not to the converter) showed a similar behavior of the flow for $S<S_c$ where the ablative regimes dominate the heating process.

The diagram shown in Fig. 5 indicates the possible evolutions of the heating process depending on the value of the radiation flux S. Actually, in a typical target for ion-beam fusion, it does not proceed to a constant S even if the incident flux S_R is kept constant (Fig. 1). This is because an increasing reemitted flux σT^4 will be circulating between the converter and the heated region. This reemitted flux is small in the NLTE phase and so $S_R=S+\sigma T^4 \simeq S$, but it becomes important as the equilibrium is approached. In an open geometry, the reemitted flux is lost and then, for inertial-confinement fusion, the working region should be limited to the NLTE phase [11]. Such a limitation may be very restrictive because the available time may be very short, as it is in the case shown in Fig. 5: $\tau_c \simeq 0.34$ ns ($\rho_0=2.7$ g/cm^3). This time can be enlarged by using a low-density foam, but it may limit the working radiation flux to a very low value. However, in a ion-beam-fusion target, the converter is actually a closed cavity where the radiation remains confined between the external tamper and the payload [1–5].

Therefore the AHW-LTE phase is the most interesting regime in which an adequately uniform pressure on the payload can be obtained. Nonetheless, the initial AHW-NLTE phase can play a role by compensating, up to a certain extent, the shortening of the ion range that takes place in the converter as it is heated by the ion beam [1–4].

The circulating energy σT^4 for a given incident flux S_R can be calculated in the AHW-LTE phase from Eqs. (2) and (39), and from the energy balance:

$$S_R=\sigma T_R^4=S(1+N)\ , \tag{47}$$

$$\begin{aligned} N=\sigma T^4/S=&(t/l_0)^{4/[4+\alpha(m+n)]}\\ &\times(\sigma/\varepsilon_s^{4/\alpha})^{\alpha(m+n)/3[4+\alpha(m+n)]}\\ &\times S^{A/3[4+\alpha(m+n)]}\\ &\times\rho_0^{-4(m-1)/3[4+\alpha(m+n)]}\ . \end{aligned} \tag{48}$$

The factor N has been previously defined in Refs. [16] and [17] and it is called quality factor for radiation confinement. For the reference case of an aluminum payload, and for a typical temperature $T_R=200$ eV in the converter, it results in $N=3$ at $t=10$ ns. That is, only 25% of the incident flux S_R is used to drive the AHW-LTE and the rest is reemitted into the converter and remains available to drive the implosion of the payload. At this time, however, it is not possible to neglect the expansion of the converter and the previous result must be modified.

V. EFFECT OF THE CONVERTER EXPANSION

In a real target for ion-beam fusion the converter expands as it is heated by the ion beam, and it pushes the payload. According to the results of previous models, the characteristic velocity of expansion u_A is [22,23,36]

$$u_A \simeq (W_i/\rho_0)^{1/3}\ . \tag{49}$$

where W_i is the part of the beam energy flux deposited in the converter, and it is around a half of the total beam energy flux [1–4,23]. A fraction η_c of W_i is converted into thermal radiation:

$$\eta_c=S_R/W_i\ . \tag{50}$$

During the NLTE phase $S_R \simeq S$ and $u_A \le (S/\rho_0)^{1/3}$. Thus the perfect tamping approximation is a reasonably good assumption. But, if $N \gg 1$ and it becomes $W_i \gg S$, then $u_A \gg (S/\rho_0)^{1/3}$ and the converter behaves as an active tamper. It is possible that, when such a situation takes place, the shock may have enough time to run through the whole payload. Then, a rarefaction wave is reflected in the internal surface of the payload and a new shock is launched when it arrives to the ablation front. This second shock will encounter material with a velocity $u_c \neq 0$. In this case, however, u_c will be a few times v_0 and it only introduces small changes in the numerical factor of the quadratic term of Eq. (30). Therefore the results are not essentially different if $u_c=0$ is taken even when $u_A \gg (S/\rho_0)^{1/3}$, and then Eqs. (28) and (30) yield

$$v_0 \simeq u_A=(W_i/\rho_0)^{1/3}\ , \tag{51}$$

$$v=(S/\rho_0)^{1/3}(S/W_i)^{2/3}\ . \tag{52}$$

The relative velocity $v=u_A-U_A$ between the ablation front and the payload-converter interface is reduced by a factor $(S/W_i)^{2/3}$ with respect to the perfect tamping case ($u_A=0$). As a consequence the temperature is increased in the heated region:

$$\begin{aligned} \varepsilon \simeq &(t\varepsilon_s^{4/\alpha}/l_0\sigma)^{\alpha/[4+\alpha(m+n)]}S^{2\alpha(m+2)/3[3+\alpha(m+n)]}\\ &\times\rho_0^{\alpha(m-1)/3[4+\alpha(m+n)]}\\ &\times(W_i/S)^{\alpha(2m+1)/3[4+\alpha(m+n)]}\ . \end{aligned} \tag{53}$$

Now, the quality factor N for the radiation confinement is found to be

$$\begin{aligned} N=&(t/l_0)^{4/[4+\alpha(m+n)]}(\sigma/\varepsilon_s^{4/\alpha})^{\alpha(m+n)/[4+\alpha(m+n)]}\\ &\times S^{A/3[4+\alpha(m+n)]}\rho_0^{-4(m-1)/3[4+\alpha(m+n)]}\\ &\times(W_i/S)^{4(2m+1)/3[4+\alpha(m+n)]}\ . \end{aligned} \tag{54}$$

By comparing this equation with Eq. (48) it is seen that N is considerably increased. Introducing Eqs. (47) and (50) into Eq. (54), an implicit expression for N can be obtained. However, for the limit $N \gg 1$, it is approximated by setting $S \simeq S_R/N$:

$$\begin{aligned} N \simeq &(t/l_0)(\sigma/\varepsilon_s^{4/\alpha})^{\alpha(m+n)/4}S_R^{A/12}\\ &\times\rho_0^{-(m-1)/3}\eta_c^{-(2m+1)/3}\ . \end{aligned} \tag{55}$$

In the regime considered, in which the converter expansion is important, the conversion efficiency is not very high. In general, the fraction of the beam energy deposited in the converter is converted into kinetic and internal energies of the converter, and kinetic energy of the payload [9,12]. Thus, as an estimate, it is assumed that

$\eta_c \leq 0.3-0.5$. For the previous example, $N \geq 25-50$, that is, one order of magnitude higher than in the case with $u_A = 0$. This means that only a very small part of the energy is drained by the AHW-LTE and most of the energy converted into radiation is being continuously absorbed and reemitted by the heated region of the payload, redistributing itself and smoothing out the pressure which drives the payload implosion. It is clear that when this stage is reached the heat front penetrates very slowly in the payload and the mass ablation does not play any role in the implosion.

VI. CONCLUSIONS

The simple model presented in the previous sections gives a qualitative picture of the main process taking place when a cold medium is heated by the thermal radiation generated in a hot material which is in contact with it and does not allow its free expansion. The obtained scaling laws may be useful in the design of ion-beam-fusion targets. In such targets an AHW-NLTE is produced at the beginning if $q > 3/(2n-1)$. Otherwise, there exists a previous stage with a SHW-NLTE. In any case, the supersonic phase is, in general, very short, as it can be seen from Fig. 5 for $S \ll S_c$. During the NLTE stage practically all the absorbed energy is used to drive the heat wave ($S_R \simeq S$). As time proceeds and S increases, the mass ablation rate and the ablation pressure increase, and the heated region becomes hotter as the equilibrium approaches. At a time of the order of t_{EA} the equilibrium is settled and the reemitted flux becomes important. At a later time, the expansion of the converter starts to be appreciable and it contributes to an increase in the reemitted flux. As a consequence, the flux which drives the AHW is reduced and the heat-front velocity inside the payload is small in comparison with the bulk velocity of the payload. By this time, the mass ablation is not very relevant, and most of the radiation energy is continuously absorbed and reemitted from the payload. Thus the energy is redistributed in the converter, contributing to the uniformity of the thermal pressure on the payload.

Mass ablation can play a role before the expansion of the converter becomes appreciable by compensating the shortening of the ion range and providing convective stabilization of Rayleigh-Taylor instabilities [3,37].

Finally, it is noted that, due to the general form assumed for the photon mean free path in Eq. (1), the present model applies to an arbitrary mechanism of energy deposition provided that the characteristic absorption length can be expressed as a power law of the density and the temperature. In particular, it could be applied to situations where the driving radiation is monochromatic x rays [11].

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