Two-Temperature Model of the Corona of Laser-Irradiated Pellets

J.R. Sanmartín, J.L. Montañés, R. Ramis and J. Sanz
E.T.S.I.Aeronáuticos, Universidad Politécnica de Madrid
Ciudad Universitaria, Madrid-3, Spain

Abstract: A two electron-temperature, quasi-steady model of the corona of a laser-ablated pellet is considered. Ablation pressure, critical radius and mass flow rate are determined. Results are close to those obtained with heat flux saturation well below the free-streaming limit.

In recent, quasi-steady analysis[1-4] of the spherical, quasineutral corona outside the ablation surface of a laser-irradiated pellet, absorption was assumed to occur around the critical density n_{cr}. Using the continuity equations, $nvr^2 = \text{const} \equiv \mu$, the momentum and energy equations for the ion-electron fluid read

$$\frac{m_i}{Z_i} v \frac{dv}{dr} = \frac{kT}{r} \frac{2 - d\ln T/d\ln r}{1 - Z_i kT/m_i v^2} \qquad (1)$$

$$4\pi\mu\left(\frac{5}{2}kT + \frac{1}{2}\frac{m_i}{Z_i}v^2\right) + 4\pi r^2 q = \begin{matrix} 0, & r < r_{cr} \\ W, & r > r_{cr} \end{matrix} \qquad (2)$$

where W is the absorbed laser power; the electron heat flux q took either the least of classical ($-\bar{K}T^{5/2}\,dT/dr$) and saturated $[fnkT(kT/m)^{1/2}\,\text{sign}(dT/dr)]$ values at each point,[4] or just the classical one everywhere (formally equivalent to letting $f \to \infty$).[2] The ion temperature does not figure in (1), (2) because the charge number Z_i was assumed large. The solution satisfies conditions at

the ablation radius r_a: $T = 0$, $v/T = k\mu/P_a r_a^2$ (3)

the sonic radius r_s ($v_s^2 = Z_i kT_s/m_i$): $2 = d\ln T/d\ln r$ or $r_s = r_{cr}$ (4)

the critical radius r_{cr}: $r^2 v = \mu/n_{cr}$ (5)

infinity: $T \to 0$ (required for T not to be multivalued or negative) (6)

The eigenvalues r_{cr}, P_a (ablation pressure), and μ could be then determined in terms of W, n_{cr}, r_a, m_i/Z_i, $\bar{K}(Z_i)$, k, m/f^2. There are three dimensionless parameters $X_1 = n_{cr} r_a^3$, $X_2 = (W/n_{cr}^{5/6})(\bar{K}/k^{7/2})^{3/4}(m_i/Z_i)^{7/8}$, and $f(A_i/Z_i)^{1/2}$ ($A_i \equiv$ ion mass number), but actually r_{cr}/r_a, $P_a/n_{cr}(m_i/Z_i)U^2 \equiv Y$ and $\mu/r_a^2 n_{cr} U$, where $U = [W/4\pi r_a^2 n_{cr}(m_i/Z_i)]^{1/3}$, are functions of just $X_2/X_1^{11/12} \equiv X$ and $f(A_i/Z_i)^{1/2}$, the first two given in Figs. 1 and 2, from Ref. 4 (full lines). The f-dependence is weak for r_{cr} but strong for P_a [$\mu(P_a)$ being unaffected by f except for very low f, a discussion on the f-dependence of $\mu/r_a^2 n_{cr} U$ is not needed].

There is (controverted[5]) evidence, from experiments and simulations, that f should lie well below the free-streaming limit ($f \simeq 0.6$) and there have been attemps[6] to determine an appropriate f as local, kinetic effect. However, no local "f" could take an universal value, and saturation at a very low f would imply saturation in a collision-dominated plasma. We assume on the contrary that saturation will only occur when collisions are not dominant (f of order of unity), strong heat flux reduction being an overall effect due to the appearance of two electron temperatures. Suppose that absorption at the critical surface produces electrons hotter by a factor β than the local main population; the hot mean free path being β^2 times the cold one, hot electrons could remain decoupled from the main electrons for $r > r_{cr}$ and over part of the overdense region, even though the cold population would be collision-dominated at least up to r_{cr}.

Assume first that the hot contribution to electron macroscopic quantities are negligible for $r < r_{cr}$; then Eqs. (1)-(5) could be used for that range, T being the main temperature, and $q = -\bar{K}T^{5/2}\,dT/dr$. One can proceed as in Refs. 2, 4: Eqs. (1), (2) could be integrated from r_a outward, using (3), for given μ and P_a; hence, choosing μ, conditions (4), (5) yield $P_a(\mu)$ and $r_{cr}(\mu)$, and also $v_{cr} = v(r_{cr})$ and $T_c = T(r_{cr})$ (the "cold" temperature). In Refs. 2, 4, W(μ) would be then determined by condition (6) but now the underdense region requires a different analysis. Assume for simplicity that the hot electrons are monoenergetic, their energy at r_{cr} being $\frac{3}{2}kT_c\beta$, and make the ansatz that due to the electric field that maintains quasineutrality their energy decreases, between r_{cr} and infinity, to a small fraction of $\frac{3}{2}kT_c\beta$. Then very few cold electrons can reach infinity, and the cold contribution to the electron current and energy flow may be neglected for all $r > r_{cr}$. Let v_∞ be the ion velocity at infinity; the hot-electron macroscopic velocity will be v_∞ too because a) the hot-to-ion density ratio will approach unity as $r \to \infty$, and b) hot and ion currents are equal for all $r > r_{cr}$. Consequently if ψ is the electrostatic potential, we get

$$\frac{1}{2}\frac{m_i}{Z_i}(v_\infty^2 - v_{cr}^2) = e(\psi_{cr} - \psi_\infty) = \frac{3}{2}kT_c\beta - \frac{1}{2}mv_\infty^2 ,$$

hot electron velocities being entirely radial at infinity due to the conservation of angular momentum. It follows that the ansatz introduced is correct:

$$\frac{1}{2}mv_\infty^2 / \frac{3}{2}kT_c\beta = (Z_i m/m_i)(1 + m_i v_{cr}^2/3\beta Z_i kT_c) \ll 1.$$

The total energy flux at the surface $r = r_{cr}^+$ would be only due to ions and hot electrons:

$$W = 4\pi\mu\frac{1}{2}\frac{m_i}{Z_i}v_{cr}^2(\mu) + 4\pi\mu\frac{3}{2}kT_c(\mu)\beta. \qquad (7)$$

Eliminating μ in (7), $P_a(\mu)$, and $r_{cr}(\mu)$, we obtain Y and r_{cr}/r_a versus X, drawn in Figs. 1, 2 (dashed lines) for a few values of β; Fig. 2 shows that for β growing with X (≈ W) as expected, the X-dependence on Y is much weaker than that corresponding to free-streaming saturation. Assume

$$kT_h = C(m_i/Z_i)^{1/3}(W/4\pi r_{cr}^2 n_{cr})^{2/3} \qquad (8)$$

C being a dimensionless constant. Many experiments have proved (8) correct for planar geometry; further, (7) shows that C, defined by (8), is a weak function of the ion-hot electron energy ratio at r_{cr}. From (8) and kT_c, as provided by the analysis itself, we obtain $\beta(X, Y, r_{cr}/r_a)$ which, when used with the dashed curves in Figs. 1, 2, yields the dotted lines. In agreement with recent experimental data for spherical geometry,[7] our results appear to correspond to values of f midway between free-streaming and the strong flux reduction ($f \simeq 0.03$) found in planar geometry.

We performed preliminary calculations that include hot-electrom effects in the overdense region. The collisionless hot distribution-function is self-consistently determined by making use of energy and angular momentum conservation; the energy flux is the only electron quantity to which hot electrons make an appreciable contribution. Quantitative modifications seem weak to moderate, not affecting the conclusions of this paper; complete results will appear elsewhere.

The validity of our work is most clearly discussed in terms of the ratio r_{cr}/r_a. For r_{cr}/r_a decreasing, both β and the cold mean free path decrease and at some point, say $r_{cr}/r_a \simeq 2.5$, hot electrons thermalize at r_{cr}; below, classical results[2] are valid. At the other end, when r_{cr}/r_a reaches about 9, the cold electrons become collisionless at r_{cr} and the hot electrons reach the ablation surface, preheating the target.

|1| S.J. Gitomer et al., Phys. Fluids 20, 234 (1977); Yu.V. Afanas'ev et al. Zh. Ekps. Teor. Fiz. 71, 594 (1976) Sov. Phys. JETP 44, 311 (1977) .
|2| J. Sanz et al., Phys. Fluids 24, 2098 (1981).
|3| C.E. Max et al., Phys. Fluids 23, 1620 (1980).
|4| J. Sanz and J.R. Sanmartín, Phys. Fluids (to be published).
|5| N. St.J. Braithwaite et al., Plasma Physics 23, 713 (1981).
|6| D. Shvarts et al., Phys. Rev. Lett. 47, 247 (1981).
|7| T.J. Goldsack et al., 15th European Conference on Laser Interaction with Matter, Schliersee, Jannuary 1982.