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Laser-induced gas-breakdown model

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Résumé. — Nous présentons un modèle unidimensionnel destiné à l'interprétation détaillée du spectre continu et discret émis par un plasma d'hélium créé par un laser à CO_2 . Ce modèle tient compte simultanément de la dynamique d'ionisation, du transfert de rayonnement, et de l'hydrodynamique (traitée par une simulation lagrangienne) du plasma. L'état d'ionisation du plasma est évalué par des modèles collisionnels radiatifs mettant en jeu de nombreux états excités pour chaque type d'ions. Le spectre émis par le plasma de claquage est calculé en fonction de la position et du temps au moyen d'un algorithme de transfert qui tient compte de la réabsorption et de l'élargissement Stark. Les résultats de la simulation qui concernent l'expansion du plasma, l'intensité absolue de l'émission lumineuse et le spectre continu et discret sont en accord favorable avec les résultats expérimentaux obtenus au laboratoire.

Abstract. — A one-dimensional model oriented toward a thorough understanding of line and continuum emission from a CO_2 laser-induced helium plasma is presented. The model combines a time-dependent calculation of the ionization and radiation dynamics with a lagrangian hydrodynamic simulation code. Detailed collisional radiative models involving many neutral and ion excited levels are used to predict the plasma ionization state. A radiation transport algorithm including self-trapping effects and Stark broadening is developed which directly gives the space- and time-resolved spectrum of the breakdown plasma. Calculated plasma expansion, time- and space-resolved absolute light intensities and line and continuum spectral shapes compare quite favorably with experimental data.

1. **Introduction.** — In this paper, we describe the details of a model for the interaction of a focused laser pulse with helium gas at moderately high pressures. Early studies of laser-induced breakdown in gases have concentrated on the power threshold for breakdown [1] and a qualitative description of the spatial structure and temporal development of the breakdown spark. Laser-induced breakdown studies have been reviewed by De Michelis [2] for papers published up to 1968 ; Ostrovskaya and Zaïdel [3] provide a thorough review of the topic up to 1971 while a brief survey by Wood [4] includes work done with CO_2 lasers at 10.6μ .

These lasers have renewed [5-8] the interest in gas-breakdown phenomena because they offer a relatively simple way of generating megawatt power light pulses at high repetition rates with good shot-to-shot reproducibility. The rapid rate of data acquisition allows precise studies on the hydrodynamic mode of propagation of the spark along the beam

through detailed space-and-time resolved measurements of plasma parameters. Partially resolved structure of CO_2 laser-induced plasmas [9-11] have been obtained using both conventional spectroscopic techniques [12, 13] and interferometric measurements [8, 10].

Our theoretical understanding of gas-breakdown phenomena has been based for a long time on the ideas developed in a very comprehensive work [14, 15] by Raïzer. The speed of the laser heated wavefront was computed on the basis of various breakdown propagation schemes, including detonation and radiative mechanisms. Comparison with experiments have shown [9, 16-17] qualitative agreement with Raïzer's predictions.

The possibility of utilizing CO_2 lasers for heating magnetically confined fusion plasmas have given rise to the development of one-dimensional models [18, 19] of laser energy addition to nonuniform plasmas. These models were designed to predict the global

magnetohydrodynamic expansion of laser-heated plasmas irrespective of their detailed ionization state and excited-level populations.

Until the expansion of laser-initiated fusion studies little attention has been paid on an accurate theoretical description of both ionization dynamics and hydrodynamic motion of laser created plasmas. Actually, laser target models like Lasnex [20] and others [21, 22] incorporate atomic physics processes among the relevant excited and ionized species. Similarly, the development of suitable diagnostics and interpretation of data in laser-induced gas-breakdown experiments points to the necessity of modeling the numerous ionization and excitation stages of the gas, radiation dynamics of the emitting plasma layers, and laser absorption wavefront motion on a well-grounded basis.

In this work, we will present a one-dimensional model oriented toward a thorough understanding of line emission from a laser-induced helium plasma. Focused laser radiation at CO_2 wavelength enters the gas and is partially reflected and absorbed at the critical electron density where the laser and plasma frequencies equalize. Two temperature hydrodynamic equations for mass, momentum and energy conservation are solved on a lagrangian grid. A collisional-radiative (CR) model based on rate-equation calculations is used to predict ground and excited states populations of neutral (He I) and ionized (He II) helium atoms.

Although this study specifically concerns helium laser-induced plasmas, the treatment of ionization dynamics and hydrodynamic motion has been designed for application to higher- Z gases by simply modifying the atomic parameters (Sect. 3.1) contained in the input data.

The frequency distribution of radiation intensity within the plasma is computed taking into account re-absorption effects. The line emission spectrum is an important test of the model with experiment : radiation intensities are evaluated at various space and time positions as they depend on the correct modeling of laser absorption, rate equations and radiation transport.

Particular attention will be directed towards the study of spatial evolution of the absorption wavefront as a function of laser power and gas pressure. With the long pulses attainable from CO_2 lasers, a transition is observed in the expansion mechanism from a radiation-driven ionization wave to a thermal expansion. Detailed comparison of the results with experimental data taken in our laboratory will be published elsewhere. However, the good agreement between the model and experiment will be illustrated through typical examples.

2. Hydrodynamic equations. — The plasma is assumed to consist of a charge-neutral mixture of electrons (e) and ground state helium atoms (0),

ground state helium ions (1), and alpha particles (2) collectively referred to as ions. Excited-ion state populations are deduced from rate equations including the main dominant collisional and radiative processes relevant to the ionic composition of the plasma. The plasma dynamics is described by 4 main states variables, namely the mass density ρ , the fluid velocity u , the electron and ion temperature T_e and T_i .

The fluid equations then write

$$\frac{\partial \rho}{\partial t} + \underline{\nabla} \cdot (\rho u) = 0 \quad (1)$$

$$\rho \left(\frac{\partial}{\partial t} + u \cdot \underline{\nabla} \right) u = - \underline{\nabla} P \quad (2)$$

$$\begin{aligned} \left(\frac{\partial}{\partial t} + u \cdot \underline{\nabla} \right) E_e + (E_e + P_e) \underline{\nabla} \cdot u = \\ = P_L - \underline{\nabla} \cdot Q_e - R_v - \frac{3}{2} \tau_{ei}^{-1} N_e k(T_e - T_i) \end{aligned} \quad (3)$$

$$\begin{aligned} \left(\frac{\partial}{\partial t} + u \cdot \underline{\nabla} \right) E_i + (E_i + P_i) \underline{\nabla} \cdot u = \\ = \frac{3}{2} \tau_{ei}^{-1} N_e k(T_e - T_i) \end{aligned} \quad (4)$$

where N is the total number density, $P = P_e + P_i$ is the total pressure, E_e is the electron thermal plus ionization energy density, E_i is the thermal ion energy density, P_L is the laser energy deposition power density and R_v is the rate of change of E_e due to radiation losses. The following additional relations exist between the previously defined variables

— electron pressure

$$P_e = N_e k T_e \quad (5)$$

— ion density

$$N = N_0 + N_1 + N_2 \quad (6)$$

— ion pressure

$$P_i = N k T_i \quad (7)$$

— charge neutrality

$$N_e = N_1 + 2 N_2 \quad (8)$$

— electron energy density

$$E_e = \frac{3}{2} N_e k T_e + E_1 N_1 + E_2 N_2 \quad (9)$$

— ion energy density

$$E_i = \frac{3}{2} N k T_i \quad (10)$$

where E_1 and E_2 are respectively the ionization potentials of HE I and He II.

The heat flux Q_e is related to temperature gradients by

$$Q_e = - K_e \underline{\nabla} T_e \quad (11)$$

where K_e is the electron thermal conductivity listed in Gauthier and Delpech [23]. Ion and neutral thermal conduction have been considered ; they were found

negligible under the present conditions. The electron-heavy particle energy transfer time writes

$$\tau_{eh} = (\nu_{e0} + \nu_{ei})^{-1} \quad (12)$$

where ν_{e0} and ν_{ei} are respectively the electron-neutral and electron-ion collision frequencies for energy transfer [24].

2.1 ABSORPTION OF LASER LIGHT. — During the initial stage of breakdown, laser energy is absorbed by inverse bremsstrahlung of electrons on neutrals and ions. The classical absorption coefficient [25] for ions of average charge Z_{eff} was used,

$$K = \frac{16 \pi Z_{\text{eff}} N_e^2 e^6 \ln A}{3 c v^2 (2 \pi m_e k T_e)^{3/2} (1 - \omega_p^2/\omega^2)^{1/2}} \quad (13)$$

where $\nu = \omega/2\pi$ is the laser radiation frequency, ω_p the plasma frequency and $\ln A$ the usual Coulomb logarithm. For neutral helium atoms, the absorption coefficient was numerically integrated from the approximate bremsstrahlung cross section of Dalgarno and Lane [26] using accurate values of the electron-helium atom momentum transfer cross section [27]. Results were in good agreement with tabulated values of Geltman [28] who used atomic model-potential calculations.

Laser energy deposition is computed by solving the 10.6μ radiation transport equation along a photon path by an analytic cell-integrated method [29]. Simulation of the laser focusing system is obtained by scaling the average energy absorbed along the laser beam to the experimentally determined axial beam profile. The integration method is energy conservative and preserves laser flux continuity. Refraction effects, which results from electron density gradients, and laser-trapping [30] processes can also be incorporated in the model. The incident wave, on reaching the critical region with electron density $N_e \approx 10^{19} \text{ cm}^{-3}$, is partially absorbed and reflected. The reflected wave further deposits energy in the underdense ionized gas.

2.2 RADIATION LOSSES. — The emitted radiation energy loss R_v in eq. (3) is made up of free-free, free-bound and line emission. The spectrum of frequencies in both continuum and line radiation is very wide and complicated in that the two stages of ionization will be passed as the temperature changes. Excited levels of each ion are capable of absorption of line and continuum radiation. In this way, radiation transport contribute to level population, ionization, and heating of the plasma.

This model is designed primarily to describe only the main features of the emitted visible and UV spectra. For this purpose, approximate values [31] of the radiative losses integrated over all wavelengths were found to be adequate. The term R_v can be written

$$R_v = R_{bb} + R_{fb} + R_{ff} \quad (14)$$

with explicitly

$$R_{bb} = \sum_{i,j,Z} A(i,j) (E(Z,i) - E(Z,j)) N(Z,i) \quad (15)$$

$$R_{fb} = N_e \sum_{Z,j} N(Z+1,0) \alpha^{rr}(Z+1,0;Z,j) \chi(Z,j) \quad (16)$$

$$R_{ff} = 1.2 \times 10^{-19} (kT_e)^{1/2} N_e \sum_{Z,j} Z^2 N(Z,j) \quad (17)$$

where $A(i,j)$ is the radiative transition probability, $\alpha^{rr}(Z+1,0;Z,j)$ the radiative recombination coefficient connecting the ground state ion stage $Z+1$ to the j -excited ion stage Z , $N(Z,i)$, $E(Z,i)$, $\chi(Z,i)$ are respectively the population, energy and ionization energy of the level labelled i in the ionization stage of charge Z .

Under the present conditions, free-free absorption is usually negligible and most of the bound-bound radiative energy losses occur in the resonance line. At the high neutral densities encountered in atmospheric pressure gas-breakdown plasmas, ground state reabsorption of the recombination continuum and resonance radiation must be considered. In such cases, the loss rate is smaller than the rate predicted by assuming the plasma to be optically thin. Accordingly, resonance radiation losses are completely negligible: absorption coefficient in excess of 10^5 cm^{-1} are typical. For free-bound radiation, eq. (16) must be modified to include an attenuation factor which incorporates bound-free reabsorption, mainly photo-ionization effects. In any case, the total radiative energy losses were compared to blackbody radiation to check the consistency of the results.

2.3 NUMERICAL TECHNIQUES. — Stability considerations influence the choice of numerical techniques for the simultaneous integration of the excitation and ionization rate equations with the hydrodynamic equations. The lagrangian form of mass, momentum and energy conservation equations is particularly well suited for one-dimensional calculations where intense compression is produced. Mass and momentum equations were solved explicitly on a lagrangian grid using time-centered finite difference approximations where the fluid velocity was computed at times half-way the energy time-grid. In addition, we introduced, as usual [32], an empirically determined pseudo-viscous pressure in eqs. (2) and (4).

Electron — and heavy particles — temperature equations were solved implicitly using a modified predictor-corrector algorithm. The heat conduction term of eq. (3) was processed by a Crank-Nicholson implicit scheme [32] with iterations in plasma regions where heat transport was important. The energy exchange term was treated by the exponential method of Christiansen and Roberts [33]. A more simple and less time-consuming solution was adopted in the low-heat conduction regime by integrating simultaneously the electron and ion energy equations

with an explicit version of temperature exchange and conduction.

The ion and excited states populations were determined from the solution of a tridiagonal system of equations (to be described in section 3) using an implicit integration scheme. We found that the ionization energy term causes a numerical instability due to the rapid variation of ionization and recombination cross-sections with electron temperature. This instability was controlled in the following manner : suppose we neglect in eq. (3) all sources, losses and inhomogeneous terms retaining only the balance between electron and ionization energy densities. We get

$$\frac{\partial}{\partial t} \left(\frac{3}{2} k N_e T_e \right) = - E_i \frac{\partial}{\partial t} N_e \quad (18)$$

which may be solved together with the electron density buildup equation

$$\frac{\partial N_e}{\partial t} = S(T_e) N_0 N_e \quad (19)$$

where $S(T_e)$ is the ionization rate. The formal implicit numerical solution of eqs. (18) and (19) is

$$T_e^+ - T_e^- = - (2 E_i / 3 k + T_e^+) S(T_e^+) N_0 \Delta t \quad (20)$$

where T_e^+ and T_e^- are the electron temperatures before and after the current time step Δt . Considering the right hand side of eq. (20), we find from studies of the ionization rate temperature variations [34] that this term is highly non linear in T_e . An approximate solution is obtained by expanding $S(T_e^+)$ to first order

$$S(T_e^+) = S(T_e^-) + (T_e^+ - T_e^-) \left[\frac{\partial S(T_e)}{\partial T_e} \right]_{T_e = T_e^-} \quad (21)$$

Thereby, we get a fully-linearized form of eq. 20, namely

$$T_e^+ = T_e^- - \left((2 E_i / 3 k + T_e^+) S(T_e^-) + (2 E_i / 3 k + T_e^-) (T_e^+ - T_e^-) \frac{\partial S(T_e)}{\partial T_e} \right) N_0 \Delta t \quad (22)$$

Total plasma energy is conserved in a time step by correcting the ion populations using a combination of eqs. (19) and (21). In that way, the plasma species populations and energies are evaluated in a fully implicit manner which is energy conservative.

The time scale for ion population changes was determined at each point of the grid for each time step together with the heat conduction time, the bremsstrahlung absorption time and the sound velocity. Hydrodynamic and rate equations integrations were performed consistently with this data. The whole procedure was checked by verifying total energy conservation.

3. Excitation and ionization dynamics. — The hydrodynamic equations (eqs. (1)-(4)) contain the various ion stage densities, the ionization energy density $E_i = E_1 N_1 + E_2 N_2$ and the effective charge $Z_{\text{eff}} = N_e / (N_0 + N_1 + N_2)$. In order to obtain these quantities, as well as the rate of change of E_e due to radiation processes discussed in 2.2, we must solve a system of rate equations which takes into account non-equilibrium collisional and radiative processes among the many excited and ionized states of the gas.

The inhomogeneous temperature and density regime characterizing laser-induced plasmas rules out the use of a local thermodynamic equilibrium (LTE) model based on the Saha equation or the coronal approximation. For many plasma regions, the atomic collision processes proceed at rates which are not fast enough to establish LTE. We have used a time dependent atomic model including most of the collisional and radiative processes which we believe to be dominant under laser gas-breakdown conditions.

3.1 ATOMIC PARAMETERS. — In this section, we give some specific rate coefficients. The accuracy of the model in deriving the actual plasma state will depend primarily on how well the rate coefficients adopted in the simulation reproduce the available measured or predicted atomic data. Fortunately, helium-like and hydrogen-like ions have been studied to a certain extent : experimental data was used for line strengths, level ionization energies, and wherever possible, for electron excitation and ionization cross-sections.

We have considered the following atomic processes connecting the ground and excited states of He I, He II ($Z = 1$) and He III ($Z = 2$).

i) Collisional ionization

ii) collisional recombination

iii) radiative recombination

iv) spontaneous emission (radiative decay)

v) collisional excitation and deexcitation

with the corresponding rate coefficients :

$$S(Z, j ; Z + 1, 0)$$

for collisional ionization of state (Z, j) to the ground state ion of next charge, $\alpha^{\text{cr}}(Z, 0 ; Z - 1, j)$ for the

inverse process of collisional (three-body) recombination, $\alpha^{rr}(Z, 0; Z-1, j)$ for radiative (two-body) recombination resulting in the emission of a photon with energy $\chi(Z, j)$, $A(i, j)$ for the radiative transition probability and $X^{-1}(i, j)$ for the collisional deexcitation rate from upper level i to lower level j of the charge Z ion ($X(i, j)$ stands for collisional excitation).

Direct double excitation, dielectronic recombination, and forbidden collisional and radiative processes were found to be negligible under the present conditions. Collisional rate coefficients were integrated by assuming the electron velocity distribution function to be maxwellian.

3.1.1 Neutral helium. — Accurate rate coefficients data are available for excited helium atoms at low electron temperature ($T_e < 0.4$ eV). Indirect [35] and direct [36, 37] experimental evidence have shown that Johnson's [38] hydrogenic reaction rates give a satisfactory description of collisional radiative recombination of He^+ ions with electrons in the low temperature range. Helium may be expected to behave in a quasi-hydrogenic fashion, as far as transfer between highly excited states is concerned. In the 1-10 eV temperature range, Johnson's cross-sections are thought to be adequate [39, 40] for principal quantum numbers $n > 2, 3$.

Accordingly we have used the Johnson's integrated cross-sections for ionization from the i -th level to evaluate $S(Z, i; Z+1, 0)$. Ionization from the ground state was taken from Lotz's compilation [34] of experimental and theoretical electron-impact ionization data. A convenient three parameter empirical formula was used for the rate coefficient, namely (in cm^3/s)

$$S(Z, 0; Z+1, 0) = 5.36 \times 10^{-6} T_e^{-3/2} [E_i(x)/x - 1.19 E_i(y)/y] \quad (28)$$

$$\text{where} \quad E_i(x) = \int_x^\infty e^{-t}/t \, dt$$

is the exponential integral, $x = \chi(Z, 0)/T_e$ (T_e in eV) and $y = x + 0.46$.

Collisional recombination rate coefficients were inferred by detailed balancing

$$N_Z \alpha^{cr}(Z, 0; Z-1, j) = N_E(Z-1, j) S(Z-1, j; Z, 0) \quad (29)$$

where N_Z is the Z ion density and $N_E(Z-1, j)$ is the equilibrium population of the j -th level of $(Z-1)$ -ion obtained from Saha equation

$$N_E(Z-1, j) = \left(\frac{2\pi\hbar^2}{mkT_e} \right)^{3/2} \frac{g(Z-1, j)}{2g(Z, 0)} N_e N_Z \times \exp\left(\frac{\chi(Z-1, j)}{kT_e} \right). \quad (30)$$

Radiative recombination into all levels of hydrogenic systems has been given by Seaton [41]

$$\alpha^{rr}(Z, 0; Z-1, j) = 5.2 \times 10^{-14} N_e Z \lambda^{1/2} (0.43 + 0.5 \ln \lambda + 0.47 \lambda^{-1/3}) \quad (31)$$

$$\text{where} \quad \lambda = \chi(Z-1, j)/T_e.$$

The transition probabilities for spontaneous emission were taken from the NBS tables [42]. Excitation rate coefficients responsible for populating the bound levels were also taken from Johnson's expressions using effective quantum numbers $n^* = n - \delta$ (where δ is the known quantum defect of helium) and real level energies.

3.1.2 Ionized helium. — With its single electron, He^+ has an H-like collisional and radiative behaviour. Collisional excitation of hydrogenic ions has been reviewed in Bates *et al.* [43] and more recently by Percival and Richards [44]. Following Tallents [45], we used the rather simple analytic expressions listed in ref. [43] for collisional excitation and ionization of He^+ ions. Again we used Lotz's expression [34] for ionization from the ground state. Radiative transition probabilities were computed using the well-known expression for hydrogen oscillators strengths [46]

$$f_{ij} = \frac{64}{3\pi\sqrt{3}} \frac{1}{2i^2} \left(\frac{1}{i^2} - \frac{1}{j^2} \right)^{-3} \frac{1}{i^3 j^3} g_{bb}(i, j) \quad (32)$$

where we used the values of the Gaunt factors g_{bb} quoted in Johnson [38].

3.2 COLLISIONAL RADIATIVE MODEL. — The collisional radiative model (CR) equations taking into account the atomic processes discussed above can be written as

$$\begin{aligned} \frac{dN(Z, j)}{dt} = & -N_e N(Z, j) S(Z, j; Z+1, 0) + N_e N_{Z+1} (\alpha^{rr}(Z+1, 0; Z, j) + N_e \alpha^{cr}(Z+1, 0; Z, j)) - \\ & - N(Z, j) \sum_{i < j} A(j, i) + \sum_{i > j} A(i, j) N(Z, i) - N(Z, j) N_e \left(\sum_{i < j} X^{-1}(j, i) + \sum_{i > j} X(j, i) \right) \\ & + N_e \left(\sum_{i < j} X(i, j) N(Z, i) + \sum_{i > j} X^{-1}(i, j) N(Z, i) \right). \end{aligned} \quad (33)$$

This model has been widely used [45, 47, 48] since the pioneering work of Bates and his coworkers [43]. Numerical integration of this infinite set of equations is greatly simplified in high density plasmas. Ionization energies of free atoms or ions must be corrected to account for Debye screening effects; this implies that bound states whose energies are between the reduced and the uncorrected ionization energies become free states. The reduction in ionization energy in a plasma writes

$$\Delta\chi(Z, 0) = \frac{(Z + 1) e^2}{\rho_D} \quad (34)$$

where ρ_D is the Debye length

$$\rho_D = \left[\frac{kT_e}{4\pi e^2 \left(N_e + \sum_Z Z^2 N_Z \right)} \right]^{1/2} \quad (35)$$

Levels with principal quantum number j greater than

$$j_{\max} = \left(\frac{Z^2 E_{II}}{\Delta\chi(Z, 0)} \right)^{1/2} \quad (36)$$

may be considered as part of the continuum.

At sufficiently large j collisional processes are high enough to bring levels in thermodynamic equilibrium. This further reduce the dimension of the system to be solved since populations for quantum numbers above j may be deduced from Saha equation (eq. (30)).

At the high densities encountered in laser-breakdown plasmas, self-absorption of line and continuum radiation may be considered in calculating the level populations. A detailed general treatment of the problem is scarcely feasible in inhomogeneous plasmas. We made the assumption that resonance radiation was fully self-absorbed while all other transitions were treated in the optically thin case. This statement is reasonable since experimental evidence based on emission and absorption data [49] indicate that only a very restricted number of lines ending on the $2^1,3S$ and $2^1,3P$ He I levels display slightly self-absorbed profiles.

A quasi steady state (QSS) solution of the rate equations may be sought by considering the large difference between ground and excited levels relaxation time constants. Conditions for the validity of the QSS solution have been discussed by Bates *et al.* [43] and Limbaugh and Mason [50] on a wide range of electron density and temperature. However, as the collisional rates used in ref. [43] differ appreciably from the rates used in the present study we have checked the results obtained by the QSS method against the numerical solution of the time dependent system of rate equations using typical breakdown-values of plasma parameters.

Following a classical technique, the level scheme of helium atom was reduced by combining levels

to form closely-spaced quasi-levels. Singlet and triplet levels with $j \geq 5$ and various orbital quantum number l were lumped together in a single equivalent hydrogenic level of ionization energy E_{II}/j^2 . Below $j = 5$, we take into account all the actual helium levels to compute light intensities over the visible spectrum. A comparison of the QSS and transient solutions of eqs. (33) is shown on figure 1 for a typical non-equilibrium situation where

$$N(0, 1) = N_0 = 1.28 \times 10^{19} \text{ cm}^{-3} \text{ (0.5 atm) ,}$$

$$N_e = 10^{18} \text{ cm}^{-3} \text{ and } T_e = 3 \text{ eV .}$$

Initial conditions for the transient solutions are $N(0, j) = 0$ for all $j \neq 1$. Transient solutions are found to approach QSS solutions on a picosecond time scale. The corresponding time variations of ion (and hence electron) densities occur on a considerably larger scale of tens of nanoseconds (which appear as constant on the time period covered in figure 1). As a result, excited state populations follow almost instantaneously any variation of ion densities to reach a new QSS equilibrium determined only by the current values of electron density and temperature. It is also clear from figure 1, that the QSS equilibrium is very far from the local thermodynamic equilibrium (LTE, represented by the line $N/N_{\text{Saha}} = 1$) under typical experimental conditions. Similar results were obtained in ionized helium, except for the fact that higher electron temperatures were necessary to significantly populate the levels by stepwise ionization.

Fig. 1. — He I excited states populations reduced to Saha equilibrium populations as a function of time for $N_0 = 1.28 \times 10^{19} \text{ cm}^{-3}$, $N_e = 10^{18} \text{ cm}^{-3}$ and $T_e = 3 \text{ eV}$. Straight lines on the right give the results of the QSS model. Equilibration time is about 5 ps. Note the large difference between QSS and LTE ($N/N_{\text{Saha}} = 1$) solutions.

The early state of ionization and recombination shows inversion as the levels fill up and the QSS population densities are established. This effect, which is most easily found in recombining plasmas, is the basis of recently proposed short wavelength lasers [51, 52]. Cooling times in the range of 10^{-13} s [51] are necessary to produce significant gain on high- Z materials UV transitions. Observations of IR laser oscillation in CO_2 laser-produced plasmas of argon, krypton and xenon have been yet reported [53].

According to the above discussion, the complete excitation state of the gas may now be described by a ground-state model where the rate of change of the ground levels of each ion state due to collisional ionization and recombination is evaluated using effective collisional-radiative rate coefficients. Such models have been studied [54, 55] in connection with laser target ionization dynamics. Each ion state of charge Z is connected to states $Z + 1$ and $Z - 1$, and the resulting ionization rate equations take a simple tridiagonal form

$$\begin{aligned} \frac{dN_{Z-1}}{dt} &= -S_{Z-1,Z} N_e N_{Z-1} + \alpha_{Z,Z-1} N_e N_Z \\ \frac{dN_Z}{dt} &= S_{Z-1,Z} N_e N_{Z-1} - (\alpha_{Z,Z-1} + S_{Z,Z+1}) N_e N_Z + \alpha_{Z+1,Z} N_e N_{Z+1} \end{aligned} \quad (37)$$

$$\frac{dN_{Z+1}}{dt} = S_{Z,Z+1} N_e N_Z - \alpha_{Z+1,Z} N_e N_{Z+1}$$

where $S_{Z-1,Z}$ and $\alpha_{Z,Z-1}$ are respectively the ionization and recombination rate coefficients of ion state Z .

These rates were deduced from a detailed comparison of the solutions of the QSS model (with excited states) and the ground-state model (no excited states), taking advantage of the linearity of eqs. (37) on ionic populations. Thereby, $S_{Z-1,Z}$ and $\alpha_{Z,Z-1}$ include the numerous collisional and radiative processes occurring between ground and excited states.

Neglecting (for the moment) hydrodynamic motion, we compared the ionization state obtained from local thermodynamic equilibrium to the one obtained from the CR model. Results for the plasma effective charge Z_{eff} as a function of electron temperature are shown on figure 2 for two typical values of the electron density. For electron temperatures below 4 eV, He^+ ions are dominant and LTE results are accurate. At higher electron temperatures, a significant fraction of He^+ ions are ionized to He^{++} ions; the higher density curves yield lower effective ionization and larger divergence from LTE results.

Combining the time dependent calculation of ground state populations (eqs. (37)), the QSS evaluation of excited state populations (eqs. (33)) and the hydrodynamic description of the plasma (eqs. (1-4)) we are in position to derive the radiative properties of the laser-breakdown plasma.

Fig. 2. — Effective ion charge as a function of electron temperature for two electron densities. Continuous curves : LTE equilibrium, dashed curves : QSS collisional radiative equilibrium.

3.3 CALCULATION OF LIGHT INTENSITIES. — The continuous and line emission spectrum emitted by the modeled plasma is an important feature to compare with experiments. An exact calculation is difficult since we have shown that reabsorbed radiation should be coupled back to electron and ion thermal energies.

Radiation emerging from one point of an optically thick plasma may reach there by many different paths. Along a photon path, the steady state equation of radiation transport, is

$$\frac{dI_v}{dx} = \rho(j_v - k_v I_v) \quad (38)$$

where j_v and k_v are the mass emission and absorption coefficients and I_v the specific radiation intensity. Even when the geometry is axisymmetric, as in the present case, the solution of eq. (38) has to be worked out at least in a 2 $D(r, \theta)$ or 2 $D(z, r)$ variable space.

In a one-dimensional problem, the question arises of the validity of radiation transport. In the experimental situation, however, the radial distribution of volume emission is obtained from the Abel inversion of side-on measurements where light intensities emitted along a chord are measured at various radial positions on a cylindrical cross-section of the plasma. By analogy, a crude estimate of the radiation distribution population perpendicular to the laser axis may be obtained by assuming the medium to be uniform along the radial position with plasma parameters derived from the 1 $D(z)$ laser plasma model. The radial extent of the plasma is supposed to be equal to the width of the focusing cone. This is strongly supported by experimental evidence during the first

500 ns of plasma evolution as shown on figure 3. The three contours represent the plasma extent as deduced from time-resolved strioscopic pictures of a one atmosphere helium-gas CO₂ laser-spark. Laser energy was one Joule in a pulse consisting of a peak ≈ 100 ns wide followed by a 1 μ s tail. The curve at 400 ns fits closely the focusing cone; at later times, the radial expansion of the plasma is significant. In addition, spectroscopic observations reported by George *et al.* [13] and verified in our laboratory [49] have shown that at times later than 400 ns, the plasma is no more homogeneous in the radial direction but rather comprises a hot core and a cold outer shell. The approximations of plasma homogeneity and negligible radial expansion may thus be thought to be correct in the first 500 ns of plasma expansion, as far as light emitting properties of the plasma are concerned.

Fig. 3. — Strioscopic picture of the time evolution of a 1 atm CO₂ laser-produced plasma of helium. Laser energy 1 J. The shape of the focusing cone is shown for comparison.

In the framework of the approximations made, the line and continuous emissions of the plasma were derived by using the well-known solution of eq. (38)

$$I_{\text{out}}(z) = \frac{j_v}{k_v} (1 - e^{-\tau}) \quad (39)$$

where $\tau = k_v L(z)$ is the plasma opacity and $L(z)$ the radial extent of the plasma. The mass emission and absorption coefficients for free-free and free-bound radiation in helium were taken from Anderson and Griem [56] taking into account the lowering of ionization energy. Line radiation was computed using tabulated transitions probabilities and statistical weights [42] together with the level populations derived from the QSS model of sect. 3.2. Doppler and pressure broadening were found to be negligible in front of Stark broadening under the present conditions. Stark widths parameters were taken from the recently revised results of Griem [57].

4. Discussion of the results. — The results that follow represent data from typical simulation runs where a TEA CO₂ laser at 10.6 μ was focused on helium gas. The influence of gas and laser characteristics on the radiation emission and plasma expansion was particularly studied. Results were obtained on the plasma dynamics both along the laser axis and in the radial direction. In the first case, the model was driven using slab geometry with the position along the laser axis as the spatial variable. In the second case, we used cylindrical geometry with the radius, measured from the laser axis, as the spatial variable.

4.1 AXIAL RESULTS. — In the model calculations we started the plasma with an uniform electron density of 10^8 cm^{-3} and an electron temperature of 1 eV. Neutrals and ions were supposed to be at room temperature. Results were found independent of this particular set of initial conditions. The total length of the simulated plasma was generally 1.5 cm. 75 cells of 200 μ each were found adequate to preserve numerical stability while keeping the computer cost to a reasonable level. Typical computing time was 4 ms per cell per iteration on a Univac 1110.

The laser pulse used in the simulation was tailored on the measured laser shape: it consists of an intense spike ≈ 100 ns wide, followed by a weaker emission about 20 % of the peak of some 1 000 ns duration. A bump in the laser tail occurs for large relative values of N_2 flow rate in the lasing mixture: this feature was occasionally introduced in the simulated laser to study the plasma heating effect of CO₂ radiation.

The focal spot area of the lens used to concentrate laser energy in the gas is an important parameter regarding laser energy deposition. In most of the runs, the focal spot diameter was set to 300 μ , close to the measured value for the lens used in the experiment. The focal length and the diameter of the lens were 60 and 23 mm, respectively.

4.1.1 Breakdown versus gas pressure. — In a first series of comparison runs, the laser energy was held constant at 270 mJ in a pulse consisting of a peak 60 ns wide (full width at half maximum) followed by a 1 μ s tail. This corresponds to a laser intensity of $\approx 2 \times 10^9 \text{ W/cm}^2$ at focus, a value close to the breakdown threshold of helium at atmospheric pressure. The choice of a low-energy pulse results in a reduction of the total number of time steps to reach the time of 500 ns after breakdown (according to the discussion in sect. 3.3) thus reducing the computational time without loss of information. The gas density was varied from $1.2 \times 10^{19} \text{ cm}^{-3}$ to $5 \times 10^{19} \text{ cm}^{-3}$ (0.45 to 1.86 atmosphere).

Results are presented in table I and figures 4-5. Table I gives the plasma energy balance as a function of time and pressure. We cover the transition from the ionization build-up regime at laser peak (45 ns) to the laser-driven expansion regime during the laser

Table I. — Plasma energy balance at 45 ns (laser peak), 200 ns (laser tail), 500 ns (low laser) for three neutral densities (cm^{-3}).

Energy (mJ)	45 ns			200 ns			500 ns		
	1.2×10^{19}	3×10^{19}	5×10^{19}	1.2×10^{19}	3×10^{19}	5×10^{19}	1.2×10^{19}	3×10^{19}	5×10^{19}
Thermal	0.5	15	23	19	28	38	27	36	44
Potential	0	21	22	29	34	31	29	35	32
Kinetic	0	< 1	1	4	5	7	5	7	10
Radiation	0	< 1	1	1	1	1	8	3	2
Absorbed	0.5	36	46	54	67	77	70	79	88

tail (200 ns) and to the thermally-driven expansion regime (500 ns). The low pressure run was prepared to favor shock wave generation and seek their effect on plasma dynamics. At laser peak, the ion density at this pressure is still too small to produce appreciable radiation and internal energy. At the two higher neutral densities, the critical electron density of 10^{19} cm^{-3} is reached. Kinetic energy is very small at that time since the plasma essentially builds up at rest in the focal volume. During the laser tail, the potential energy is almost independent on gas pressure and time. This results from the fact that the higher electron densities reached at high pressures (more laser energy is absorbed) are compensated by a smaller plasma extension. During expansion,

the total number of ions varies slowly as a result of the balance between volume growth and electron-ion recombination. On the other hand, the increase of thermal energy with pressure is not directly related to greater electron and ion temperatures, but rather to the combination of a larger number of heated neutrals at higher pressures (ions and neutrals have been supposed to be at the same temperature), a larger laser absorption and a decrease of radiation emission. The last point comes from the approximate radiation transport included in the model which states that radiation losses decrease and more energy is found in the ion states when the pressure increases.

Also, we found that kinetic energy is almost independent of pressure in the range studied. The increase of plasma front velocity as pressure decreases is balanced by the variation of total plasma mass with

Fig. 4. — Space- and time-resolved evolution of a) electron density and b) electron temperature at an initial neutral density of $5 \times 10^{19} \text{ cm}^{-3}$ (1.87 atm). Laser from left to right (energy 270 mJ). Temperature spikes on the right at 500 ns are due to shock reflections on the simulation grid end. Focus indicated by an arrow.

Fig. 5. — Same as figure 4 at an initial neutral density of $1.8 \times 10^{19} \text{ cm}^{-3}$ (0.67 atm).

neutral density since the initial volume is the same in the three cases.

Figures 4a and 5a show the electron density as a function of position along the laser axis and time for $N_0 = 5 \times 10^{19} \text{ cm}^{-3}$ and $N_0 = 1.8 \times 10^{19} \text{ cm}^{-3}$, respectively. Near the focus position, we see that the time required for the plasma to reach the cutoff density ($N_c = 10^{19} \text{ cm}^{-3}$) varies considerably with neutral gas density. This is a direct result of the ionization dynamics included in the calculation. Simple considerations show [15] that the characteristic ionization time is inversely proportional to the gas density and the time-averaged laser power. It depends only logarithmically on the initial electron density. Results shown on figures 4-5 and obtained at various laser powers reproduced quite well the above predictions.

The wave velocity in laser-created plasmas has been studied in great detail [15] (we use Raizer's terminology for the various waves). For the conditions studied here, it turns out that the *absorption wave* velocity ($v \approx 4 \times 10^6 \text{ cm/s}$) is nearly equal to the *breakdown wave* velocity at $N_0 = 3 \times 10^{19} \text{ cm}^{-3}$. However, pressure variations of these waves are opposite. This is particularly illustrated in the data shown in figures 4-5. By 500 ns, the plasma front has reached approximately the same position around 4-5 mm toward the lens in the two pressure cases. At high pressure, the large initial breakdown velocity is balanced by a smaller thermally driven expansion velocity during the laser tail, when laser power has been greatly reduced. Obviously, the contrary applies in the low pressure case.

Opposite to the plasma front moving toward the lens (the backward front), a forward front travelling at a much lower velocity is present in the simulations. The velocity of this front, which is slightly dependent on gas pressure, indicates a thermally driven expansion regime. The forward wave has been observed in some details in low [10] and high [58] pressure laser induced helium plasmas.

The electron temperature is shown in figures 4b and 5b. During the laser peak, the energy deposition within the gas takes place in two distinct successive regimes which are particularly illustrated on figure 4b. The first 30 ns are governed by the partial absorption of the laser beam going with ionization of the gas. The temperature reaches its maximum in the vicinity of the focus at $z \approx 11.6 \text{ mm}$ and, otherwise, follows closely the shape of the focusing cone. Ionization is particularly intense around the focal spot resulting in greater laser absorption in this region, where the avalanche process of breakdown starts. The laser is totally absorbed when the electron density equals the critical density. Then, a second regime arises where, first, the electron temperature falls down to the neutral temperature behind the cutoff region and, second, the laser energy is absorbed in a thin layer toward the lens resulting in a steepening of the

electron temperature gradient. The absorbing layer is also visible on the electron density plots in figures 4-5.

Then, the backward front expansion velocity is strongly reduced and an hydrodynamic shock wave separates from the ionization front. This wave originates in the primary breakdown as a result of the expansion of the ionized gas against the background pressure. The shock parameters follow closely the Hugoniot relations, particularly for the wavefront velocity which is within 10 % of the Hugoniot value over the range of processes covered in figures 4-5. A secondary hydrodynamic wave separates from the forward front at a much smaller velocity. These hydrodynamic compression waves have little effect on the electron density in front of the plasma. As they propagate in the cold gas, their energy is not sufficient to heat the gas (and hence the electrons) above the threshold temperature for appreciable ionization.

The shape of the laser absorption front in the two pressure cases considered here may be explained by balancing the ionization and laser absorption times against the volume expansion time. At high neutral densities, ionization is fast and hydrodynamic expansion is slow resulting in a relatively narrow absorption front, as depicted on figure 4b. At lower initial gas densities, the increasing plasma pressure in the laser heated zone rapidly exceeds the pressure of the surrounding regions. As a result, the laser heated region expands mainly in the forward direction where the plasma mass density is lower due to initial breakdown expansion. This is well illustrated on figure 5b where the absorption front motion is seen to be almost stopped in the backward direction while the rear side of absorption front propagates rapidly in the forward direction.

The physical situation described above, where, as the velocity of expansion decreases, the dense hot plasma streams reverse into the rarefaction zone has yet been observed in Q-switched laser-induced breakdown of helium [58, 59]. This process is strongly enhanced in CO₂ laser-induced breakdown where the laser energy contained in the pulse tail is still absorbed at times very far from breakdown.

As the size of the absorption front increases, laser energy is deposited in deeper regions of the breakdown plasma. This plasma is highly ionized with a majority of He⁺ ions. The density of He²⁺ ions depends strongly on the electron temperature behind the front, as given by the collisional-radiative model. Thereby, laser energy is preferably converted in thermal energy rather than ionization energy during this period. A gradual increase of electron temperature results at times later than 200 ns, as depicted on figures 4b and 5b.

Experimentally, this phase of plasma evolution coincides with the onset of the He II line spectrum, particularly the 4 686 Å ($n = 4 \rightarrow n = 3$) transition. Figure 6 shows the measured intensity in

Fig. 6. — Absolute light intensities measured at the centre of the $n = 4$ to $n = 3$ He II (4 686 Å) transition as a function of axial position and time. Laser from left to right focused at $z = 11.66$ mm. 300 arbitrary units correspond to

$$2.5 \times 10^{-4} \text{ erg.s}^{-1} \text{ ster}^{-1} \text{ Hz}^{-1} \text{ cm}^{-2}.$$

$\text{erg.s}^{-1} \text{ ster}^{-1} \text{ Hz}^{-1} \text{ cm}^{-2}$ emitted at 4 686 Å by a CO₂ laser-induced helium plasma at atmospheric pressure as a function of time and position along the laser axis. Details of the experimental equipment and procedure may be found elsewhere [11]. Light intensities were calibrated absolutely against a tungsten ribbon lamp. The laser pulse, which propagates from left to right, is focused around $z = 11.66$ mm. The plasma front expands rapidly toward the lens during the first 300 ns. Then, the expansion velocity decreases strongly, while the absorption front grows larger in the forward direction. The initial peak is mainly continuum radiation rather than He II emission, as shown by companion experiments [49]. At later times, He II emission is dominant when the energy contained in the laser tail contributes mainly to heat the steady plasma front.

Figure 7 shows the computed intensity emitted at 4 686 Å using input parameters close to the experimental conditions. The radiation transport scheme devised for integrated-light scans along the laser axis has been used to treat approximately self absorption effects. Semi quantitative agreement with the experimental expansion is found at times greater than

Fig. 7. — Absolute light intensities computed at the center of the He II (4 686 Å) transition as a function of axial position (focus at 11.6 mm) and time. Laser from left to right.

roughly 200 ns. More important is the very good agreement between the absolute values of the measured and computed light intensities which are within a factor of three over most of the expansion.

4.1.2 Breakdown versus laser energy. — In a second series of runs, the initial gas density was kept constant at $N_0 = 3 \times 10^{19} \text{ cm}^{-3}$ (1.12 atm) and the laser energy contained in the pulse was varied. However, the laser pulse shape used in the computation was kept the same. Again, a laser tail 20 % of the peak for about 1 000 ns was used, to simulate the real laser (the influence of the laser tail on the plasma ionization state and expansion will be treated separately in the next section). Figure 8 shows the electron density and temperature as a function of time and position along the laser axis at a laser energy of 810 mJ, that is three times the laser energy used in the simulations given in figures 4-5.

Fig. 8. — Space- and time-resolved evolution of a) electron density and b) electron temperature at an initial neutral density of $3 \times 10^{19} \text{ cm}^{-3}$ but at a laser energy of 810 mJ. Note the increase in plasma extension and the higher peak electron temperature.

During the first 100 ns, the breakdown starts sharply during the rising edge of the laser pulse. A very fast ionization wave grows toward the lens at a velocity roughly two times larger than the front velocity shown in figure 4. During the laser tail (later than 100 ns), the absorption front development follows qualitatively the conclusions of the previous

section. This is particularly true for the complicated system of heating waves established in the plasma as depicted in the temperature plot of figure 8b. The total expansion of the plasma is more dependent on laser energy than on gas pressure. This is explained qualitatively by the fact that both the initial ionization wave velocity (which scales linearly with laser energy) and the thermally-driven plasma expansion velocity are increasing functions of laser energy. On the contrary, according to the discussion of sect. 4.a, these two phases of plasma development have opposite variations of their characteristic velocities with gas density.

The detailed comparison of the longitudinal plasma expansion as a function of incident laser energy is an important test of our model with experiment, providing the unidimensional description of the plasma is valid. This is done on figure 9 for two laser energies in a 3 : 1 ratio. The axial position of the plasma front (toward the lens) is plotted as a function of time after the beginning of the laser pulse. In theory and experiment, distances along the laser axis, are referred to the focus position. The uncertainty in the experimental time scale origin compared with the origin in the test run is ± 10 ns, the resolution of the experiment. Error bars on the experimental points comprise the spatial resolution (0.1 mm) and shot-to-shot reproducibility of the plasma position at a given time. The agreement between experiment and the code results is quite good at 270 mJ. Experimental errors in laser energy measurements ($\approx 10\%$) and focusing parameters of the lens ($\approx 30\%$) may explain the small discrepancy at 810 mJ.

Fig. 9. — Plasma front position along laser axis as a function of time and laser energy in helium at an initial density of $5 \times 10^{19} \text{ cm}^{-3}$. Continuous line : model results. Points : experimental results. Dashed lines are an aid to the eye.

4.1.3 *Influence of the laser tail on breakdown.* — The preceding simulations were performed with a monotonically decreasing tail in the laser pulse. This is the case for a CO_2 laser operating with a low relative concentration of N_2 in the gas mixture. However, by increasing the N_2 partial pressure, it is possible to create a large bump in the laser tail, and to investigate post-breakdown laser heating effects [49]. However, in a real experiment, it is difficult to isolate the influence of the laser tail on the experimental results. Any change in the gas mixture composition produces a modification of the discharge conditions and hence a change in both the laser peak and tail output. The breakdown simulation gives a data set free of these shortcomings. In addition, the model provides a large number of virtually unmeasurable plasma parameters necessary for characterizing the laser-heated plasma. Results are shown on figure 10. The light emitted on the He I 5 876 Å ($3^3\text{D}-2^3\text{P}$) transition has been plotted as a function of axial position and time together with the laser output pulse shape. Case a) shows plasma evolution without tail and case b) shows the effect of a bump in the tail approximately 40 % the peak amplitude, centred at 300 ns. In the first case, the initial break-

Fig. 10. — Computed light intensities emitted on the He I (5 876 Å) transition as a function of axial position and time. Laser from left to right, focus at 11.66 mm. The temporal shape of the laser pulse is shown in both cases a) without laser tail ; b) with a bump in the tail at 300 ns. Note the prominent forward wave at $t > 400$ ns in case b).

down wave slows down quickly after the end of the pulse. In the second case, the motion of the plasma front is clearly accelerated by laser absorption and more light is emitted due to elevated electron temperatures. A maximum in light emission, correlated to the bump in the tail, occurs around 300 ns. The later decrease in the He I light emission is correlated with the apparition of He II lines. In both cases, the rear side of the plasma front streams reverse in the forward direction.

4.2 RADIAL RESULTS. — The results presented in this section are taken from cylindrical simulation runs. The main objective in simulating light intensities emitted along a chord of the cylindrical plasma was to provide a direct comparison between the simulation model predictions and parameters which can be experimentally measured. We confined ourselves in the modelisation of plasma expansion well behind the ionization front. In this region, the laser has been attenuated exponentially by absorption in the over-dense plasma moving toward the lens.

Initial plasma parameters are determined through an axial simulation at the wanted gas pressure and laser intensity. The radial simulation begins at the time where laser cutoff has been reached in the axial simulation. The initial electron temperature and density profiles are deduced from experiments, wherever possible. They generally consist of a cosine squared or a rectangular peak centred around $r = 0$ with the measured half width at half maximum (HWHM). Electron density and temperature, ion and neutral densities and temperature on axis ($r = 0$) are initialized using the values obtained in the axial run, near the laser focus.

Figure 11 shows the electron density and temperature as a function of time and radial position at an helium pressure of 1.12 atmosphere. On axis, the initial electron density was $1.51 \times 10^{19} \text{ cm}^{-3}$ and electron temperature was 3 eV, as deduced from the axial run (this corresponds to the results on figure 8 at 30 ns). A cosine squared profile (HWHM = 250 μ) was superimposed on an uniform background (see sect. 4.1). By 30 ns, the electron density and temperature profiles are nearly rectangular (Fig. 11a). This, again, is a result of the ionization dynamics incorporated in the model. In the core, plasma parameters are in collisional radiative equilibrium since they are directly issued from the axial run at times where this equilibrium prevails. The outer plasma parameters, which are initialized too far from equilibrium, are brought to equilibrium in a time very much shorter than the characteristic expansion time. Time resolved experiments of the first 50 ns after breakdown in helium at low [10] and high [60] pressures have shown such nearly rectangular profiles.

After 30 ns, the plasma begins to expand in the radial direction. The electron temperature and density are minimum on-axis, with a maximum behind the

Fig. 11. — Computed electron density *a*) and temperature *b*) as a function of time and radial position at $N_0 = 3 \times 10^{19} \text{ cm}^{-3}$. Initial parameters described in the text.

radially expanding ionization wave. An hydrodynamic shock wave, with a compression ratio of 3:1, propagates in the unionized gas. No additional ionization is created by this wave since its associated electron temperature jump is only a fraction of eV. The shock velocity is in good agreement with a similarity solution of the hydrodynamic equations for a cylindrical blast wave [10].

At late times, electron density and temperature profiles are nearly rectangular. Space and time resolved measurements of electron density in a CO₂ laser-produced plasma of helium [49] have shown such uniform profiles. Measured electron densities are in satisfactory agreement with the prediction of the model except for the shape of the outer density gradient which is much smoother. Also, the radial extension of the plasma is well reproduced by the model.

The accuracy of the present simulation model in predicting the actual plasma state may be illustrated at best in the comparison of the computed Stark-broadened line spectrum to experiment. Results obtained in our laboratory are shown on figure 12a at $N_0 = 5 \times 10^{19} \text{ cm}^{-3}$ and at laser energy about 1 J. Measurements were made along a chord on a cylindrical cross-section of the plasma at $r = 0.4 \text{ mm}$

Fig. 12. — Measured *a*) and computed *b*) spectrum around the $3^3\text{D}-2^3\text{P}$ He I (5876 Å) transition at $z = 1.3$ mm and $r \approx 0.5$ mm. Note the good agreement between absolute intensity values : in the experiment 1 unit = $1.2 \times 10^{-4} \text{ erg.s}^{-1}.\text{ster}^{-1}.\text{Hz}^{-1}.\text{cm}^{-2}$.

and $z = 1.3$ mm in the direction of the lens over a 300 Å wavelength interval centred around the $3^3\text{D}-2^3\text{P}$ (5876 Å) He I transition. It takes about 50 ns (after the beginning of the laser pulse) for the plasma front to reach the observation position. The line is strongly broadened during the first 100 ns due to the elevated electron density. Continuum radiation (free-free and free-bound) contributes to the total intensity of the line during the first 200 ns. The time variation of the electron density is clearly displayed by the narrowing of the lines. They are also shifted to the blue, in agreement with Griem's stark-broadening theory [57]. Significant self-absorption at 5876 Å (on the red side of the line) is seen on the experimental profiles. Results of a simulation experiment with the initial conditions given for figure 11 are shown on figure 12*b* at $r = 0.5$ mm. The agreement between the predicted and measured line profiles is excellent both for the time variations and linewidths. Also, the computed light intensity reproduce, within a factor a two, the absolutely

calibrated intensities from experiment. Moreover, self absorption is clearly apparent in the simulation run. More detailed comparison between the computed 2^3P population and experimental data [49] obtained from time-resolved laser-based absorption measurements have shown a satisfactory agreement.

5. Conclusion. — By combining a sophisticated time-dependent calculation of the ionization and radiation dynamics of a dense inhomogeneous plasma with a lagrangian hydrodynamic computer code, we have studied the radiation emission and plasma dynamics of a medium pressure laser-induced spark in helium gas.

The major conclusions suggested from the calculations and their comparison to experiment are two fold. First, the use of approximate ionization models based on the LTE or coronal approximations is clearly inappropriate for the densities and temperatures encountered here, particularly for He II ionization. One consequence of this result is that the LTE model yields incorrect evaluation of the temperature at which a given ionization stage and line radiation peak, invalidating the usefulness and accuracy of spectroscopic techniques employed to diagnose the temperature of plasmas. Second, the common method of breakdown analysis using ionization front velocity measurements is misleading. Laser spark evolution is clearly governed by several, simultaneous processes where the balance between kinetic, internal and radiation energies play a fundamental role.

As radiation emission serves as an important diagnostic tool in the experimental studies of the interaction of pulsed laser light with a gas, the simulation model was primarily designed to calculate the time-resolved spectrum including radiation-trapping effects. Particularly, radiation losses were coupled back to the ionization dynamics by an appropriate radiation transfer scheme. The satisfactory agreement between measured and computed absolute intensities, linewidths, spectra and plasma extension is gratifying, owing to the relative simplicity of the one-dimensional calculations. Although this study specifically concerns CO_2 laser-induced plasmas of helium, the flexibility of the ionization dynamics and laser absorption description allows for its application to multiply-charged ions in heavier gases and to other laser wavelengths.

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