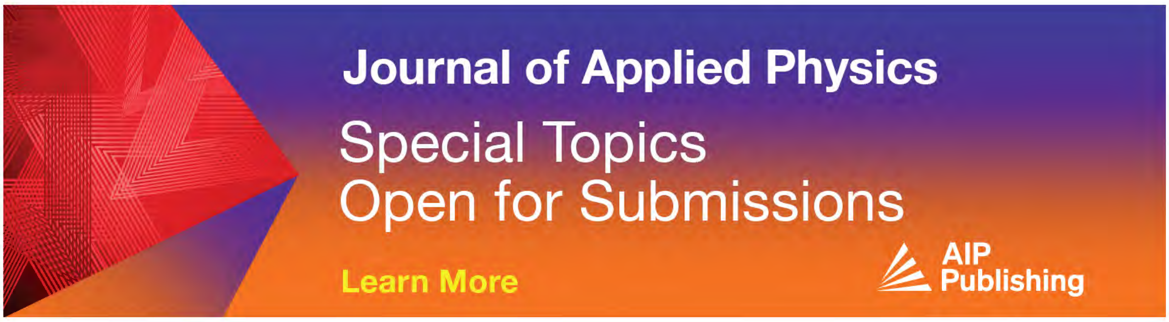
RESEARCH ARTICLE | NOVEMBER 01 1985

A model for time dependence in shock-induced thermal radiation of light

Mark B. Boslough

J. Appl. Phys. 58, 3394–3399 (1985)
https://doi.org/10.1063/1.335756



A model for time dependence in shock-induced thermal radiation of light

Mark B. Boslough
Sandia National Laboratories, Albuquerque, New Mexico 87185

(Received 7 May 1985; accepted for publication 24 July 1985)

High-speed optical pyrometry has seen increasing application in the measurement of shock temperatures in initially transparent solids and liquids; however, the information contained in the time-dependent intensity of the emitted light has frequently been overlooked. A model has been developed for this time dependence in the observed intensity of light emitted from materials undergoing high-pressure shock loading. Most experimental observations of this time dependence can be explained on the basis of geometric effects only, without having to invoke intrinsic time dependences of the source intensity (due to changes in temperature, emissivity or shock-wave structure). By taking advantage of this fact, observed time dependences can be used to determine the absorption coefficient of shocked materials and their effective emissivities, thereby providing more precise temperature measurements. The model is invoked under various limiting conditions to explain time dependences previously observed in NaCl, CaO, Mg_2SiO_4 (forsterite), SiO_2 (quartz), MgO, and $CaAl_2Si_2O_8$ (anorthite) glass. As an example, the linear absorption coefficient at 650 nm of NaCl shocked to 75 GPa is found to be 13 cm^{-1}, close to previously published values based on a similar but less general model.

I. INTRODUCTION

The optical measurement of light radiated from material which has been compressed by strong shock waves has been used increasingly to characterize the high-pressure state. In most cases, the light intensity is used to determine the temperature of the shocked material, a variable which is much more sensitive to energetics and phase transitions than pressure and density, the variables determined by more conventional shock Hugoniot experiments. This type of experiment has been successfully applied to studies in condensed matter physics and chemistry, geophysics, and related areas.

Work in this area was pioneered by Kormer,[1-3] who found that when initially transparent materials are shocked to sufficiently high pressure, they radiate light. By measuring at two different wavelengths, the intensity of light which escapes through the transparent, unshocked layer of a sample (Fig. 1), Kormer and co-workers where able to determine the color temperture of shocked alkali halides. Similar techniques have been used more recently by other shock-wave researchers,[4-10] using optical pyrometers[11,12] with wavelengths in the visible and near infrared, and fitting the data with a graybody distribution function to obtain the temperature and emissivity of materials shocked to pressures in the 100 GPa range. These workers have determined shock temperatures for silicates and oxides of geophysical interest, as well as for water, sodium chloride, and organic liquids. The optical pyrometers used by these investigators are capable of making time-resolved measurements, but in many cases advantage has not been taken of the observed time-dependent intensity. The purpose of the present paper is to model the source of the observed time dependence, to compare the limits of the model to observations, and to illustrate the usefulness of measured time dependences within the scope of the model.

II. MODELS FOR TIME DEPENDENCE

A. General model

Time dependences in the intensity of light radiated by shocked materials can be put into two classifications. "Geometric" time dependences are due to the homogeneous absorption and emission of light, and to interface cooling. "Intrinsic" time dependences are due to nonsteady shock waves, or to time-dependent behavior of some physical parameter associated with the shocked material. The latter type of time dependence has been observed, for example, in fused silica at 73 GPa,[8] and in anorthite glass at 85 GPa[11]; in both cases this occurs near inferred phase transformations of these materials, suggesting a phase instability as the cause. Despite its importance, intrinsic time dependence is beyond the scope of this paper.

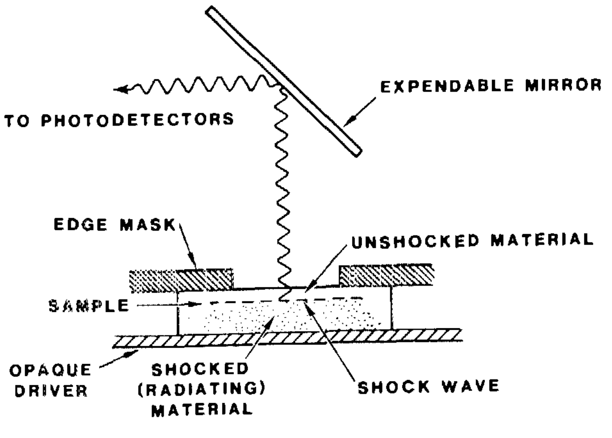


FIG. 1. Configuration of a typical shock-temperature experiment. The projectile is incident from below.

0021-8979/85/213394-06$02.40

Most workers in the field make the assumption that the spectral radiance $I(\lambda, T)$ of light of wavelength λ radiated in a shock-temperature experiment is described by a graybody distribution,

$$I(\lambda, T_s) = \epsilon_s f(\lambda, T_s), \tag{1}$$

where ϵ_s is the emissivity of the shocked layer, T_s is the shock temperature, and

$$f(\lambda, T) = C_1 \lambda^{-5} (e C_2^{/\lambda T} - 1)^{-1} \tag{2}$$

is the Planck function. C_1 and C_2 are constants. The spectral radiance data for most materials shocked above 2500 K fit this function very well, giving good experimental values for ϵ_s and T_s. It is frequently found, however, that when a two-parameter least-squares fit is carried out for ϵ_s and T_s, ϵ_s goes to unphysical values greater than unity. This arises from the fact that errors in the measured spectral radiance weigh more heavily on ϵ_s than on T_s, and these two variables are balanced against each other in the fit. It is therefore important to try to constrain ϵ_s in another way.

In the types of experiments described above, a shock wave is driven into a material with a low initial temperature $T_u \approx 300$ K) and a linear absorption coefficient a_u. The shock carries the material to a highly compressed state, with a much higher temperatrure T_s and a different linear absorption coefficient a_s. The absorbance of the unshocked layer is $A_u = 1 - e^{-a_u d_u}$ and that of the shocked layer is $A_s = 1 - e^{-a_s d_s}$, where d_u and d_s are the thicknesses of the unshocked and shocked layers, respectively. In general the linear absorption coefficients are functions of wavelength, so absorbances will be also. Kirchoff's radiation law[13] states that, at a given wavelength, the emissivity of a layer of nonreflecting material is $\epsilon = A$, where A is the absorbance.

The fraction of light at a given wavelength radiated from the shocked material at temperature T_s in a sample undergoing shock compression at time t will be approximately

$$I_s(t) = f(T_s) A_s (1 - R_s)(1 - A_u)(1 - R_f), \tag{3}$$

where R_s is the reflectivity of the shock front, and R_f is the reflectivity of the sample's free surface. In general these reflectivities are functions of wavelength; however, Eq. (3) is valid for monochromatic light and the λ dependence need not be explicitly referred to. Here the source of the radiation is the function $f(T_s)$, radiated with emissivity $\epsilon = A_s$. It is reduced by reflections at the two surfaces by factors $(1 - R_s)$ and $(1 - R_f)$, and by absorption in the unshocked layer by a factor $(1 - A_u)$. Typical reflectivities of dielectrics subjected to shock temperature experiments are about 0.1. Reflectivities of the shock front can be estimated using the known jump in density and the unknown initial index of refraction with the Gladstone–Dale[14] law. For experiments in the 100 GPa pressure range, this reflectivity is also about 0.1. Multiple internal reflections generate terms on the order of R^2 $(R \approx R_s \approx R_f)$ and are ignored in Eq. (3). The reflectivities can be combined into a wavelength-dependent reflection reduction factor, $r = (1 - R_f)(1 - R_s)$, which is about 0.8. For a sample of thickness d, with shock velocity U_s and particle velocity (velocity of shocked material) u_p, Eq. (3), at a particular wavelength, becomes

$$I_s(t) = r f(T_s)(1 - e^{-a_s(U_s - u_p)t})\, e^{-a_u(d - U_s t)}, \tag{4}$$

where t is time after impact. This equation is valid during shock transit through the sample; at times between 0 and d/U_s.

To generalize this equation further, one must also consider the light radiated from the interface between the driver plate and sample (Fig. 1), or between the flyer and sample in the case of a direct impact. Surface imperfections typically leave a gap between the driver and sample of about 1 μm. Multiple reverberations[15] and shock heating of trapped gas in the interface gap can cause the temperature of the interface to be significantly higher than the shock temperature of either boundary material. The decay of the interface temperature $T_i(t)$ is determined by solving a one-dimensional Fourier diffusion boundary value problem, the solution of which is[16]

$$T_i(t) = T_d + \frac{Q}{2\rho C_v (\pi \kappa t)^{1/2}}. \tag{5}$$

T_d is the shock temperature of the driver material, Q is the unknown quantity of heat deposited at the interface, ρ, C_v, and κ are the density, specific heat, and thermal diffusivity of the shocked driver material, respectively. The shocked dielectric sample is considered to be a thermal insulator. If the interface radiates as a graybody with emissivity ϵ_i, its contribution to the observed spectral radiance is

$$I_i(t) = \epsilon_i r f(T_d + C t^{-1/2}) e^{-a_s(U_s - u_p)t}\, e^{-a_u(d - U_s t)}, \tag{6}$$

where C contains the constants from Eq. (5).

Combining Eqs. (4) and (6) gives the general equation for geometric time dependence of observed spectral radiance at a given wavelength:

$$I(t) = r\big[f(T_s)(1 - e^{-a_s(U_s - u_p)t})\, e^{-a_u(d - U_s t)} + \epsilon_i f(T_d + C t^{-1/2}) e^{-a_s(U_s - u_p)t}\, e^{-a_u(d - U_s t)}\big]. \tag{7}$$

Note the assumption that the spectral radiance from the unshocked layer is negligible because of the low initial temperature. It is instructive to look at Eq. (7) under limiting conditions, and compare the resulting time dependence to experimental data.

B. The "ideal case"

In the "ideal case," there is no interface contribution $[\epsilon_i f(T_d + C t^{-1/2}) \ll f(T_s)]$, the initial state is transparent $(a_u \approx 0)$, and the shocked material is opaque $(a_s \gg 0)$. Equation (7) reduces to

$$I(t) = r f(T_s), \tag{8}$$

which is independent of time. This behavior is illustrated in Fig. 2(b). Note that Eq. (7) is only valid between the time of shock arrival at the interface t_0 and the time of shock wave breakout at the free surface t_1. Prior to t_0, $I(\epsilon) = 0$. After t_1, $I(t)$ depends on the post-shock temperature of the material, a problem which is not addressed in this paper.

C. Time-dependent case 1

For an experiment in which there is no significant interface contribution $[\epsilon_i f(T_d + C t^{-1/2}) \ll f(T_s)]$ and an initial-

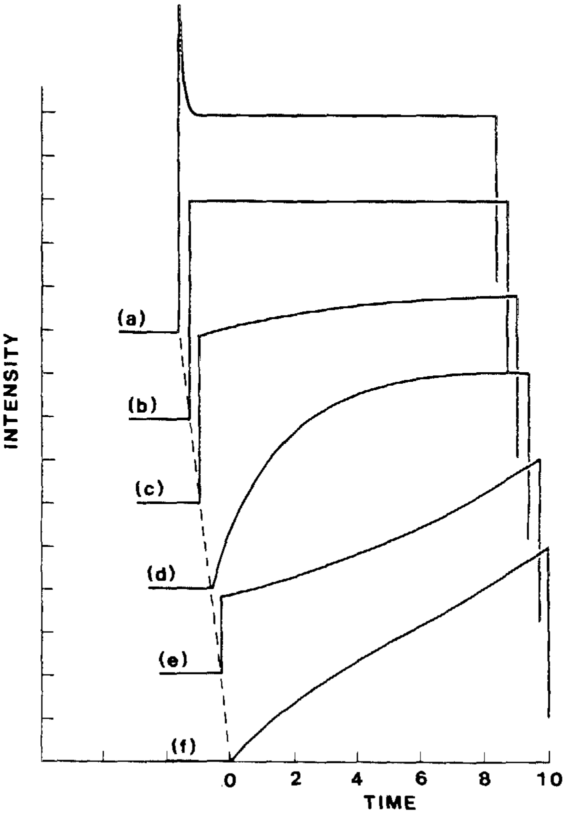


FIG. 2. Time-dependent cases, with $f(T_s) = 10$, $d/U_s = 10$. (a) Case 4 with $\epsilon_i f(T') = 15$, $a_s(U_s - u_p) = 10$; (b) the "ideal" case; (c) case 5, $a_s(U_s - u_p) = 0.2$, $\epsilon_i = 0.8$; (d) case 1, with $a_s(U_s - u_p) = 0.5$; (e) case 2, with $a_u = 0.1$; (f) case 3, with $a_s(U_s - U_p) = 0.5$, $a_u = 0.1$.

ly transparent sample ($a_u \approx 0$), Eq. (7) becomes

$$I(t) = rf(T_s)[1 - e^{-a_s(U_s - u_p)t}]. \tag{9}$$

For certain values of the other parameters the light intensity is seen to exponentially rise to a constant [Fig. 2(d)]. This case was described by Kormer *et al.*[2] and explains the time dependence of their data on NaCl (Fig. 3).

D. Time-dependent case 2

Another possible situation is the case in which the initial absorptivity is nonzero ($a_u \neq 0$) and the shocked material is opaque ($a_s > 0$). Equation (7) becomes

$$I(t) = rf(T_s)e^{-a_u(U_s - u_p)t}, \tag{10}$$

and the time dependence can appear as in Fig. 2(e). This case offers an explanation for the data on CaO,[9] in Fig. 4. The material used in that study had an initially wavelength-dependent, nonzero absorptivity, and the observed time dependence demonstrated a consistent wavelength dependence.

E. Time-dependent case 3

In a combination of cases 1 and 2 there is no interface contribution $[\epsilon_i f(T_d + ct^{-1/2}) \ll f(T_s)]$ but a_s and a_u are both finite and nonzero. Equation (7) becomes

$$I(t) = rf(T_s)(1 - e^{-a_s(U_s - u_p)t})e^{-a_u(d - U_s)t}. \tag{11}$$

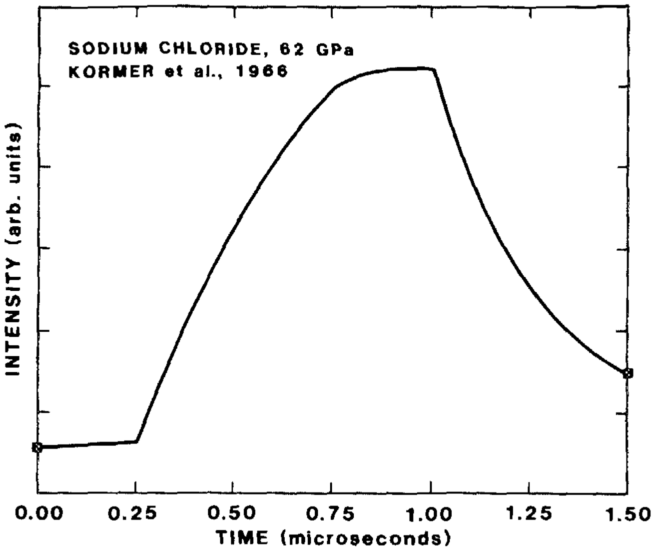


FIG. 3. Time dependence of 478-nm light radiated from NaCl shocked to 62 GPa (Kormer *et al.*, see Ref. 2).

For certain combinations of values for the other parameters, the time dependence can appear to be an almost linear function, as illustrated in Fig. 2(f). However, no data could be found which demonstrated this behavior.

F. Time-dependent case 4

Perhaps the most ubiquitous form of time dependence observed in shock-temperature data is a sharp rise at $t = 0$, followed by an approximate exponential decay to a constant value. This occurs when the interface temperature passes through some finite temperature T' (greater than T_s) at time τ, where τ is the photodetector response time, ($T_d + ct^{-1/2} \approx T'$ for $t \approx \tau$), and when the rate of change of $f(T_d + Ct^{-1/2})$ is much smaller than that of $I_s(t)$ for $t \approx \Delta t$, where Δt is the time for the shock wave to transit one optical

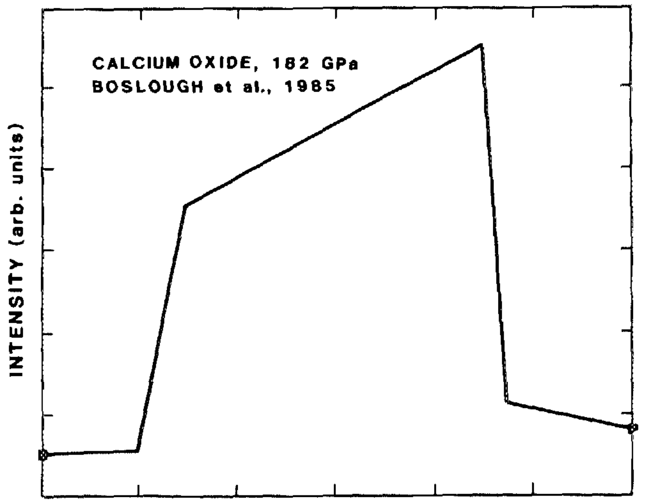


FIG. 4. Time dependence in 508-nm light radiated from CaO shocked to 182 GPa (Boslough *et al.*, see Ref. 9).

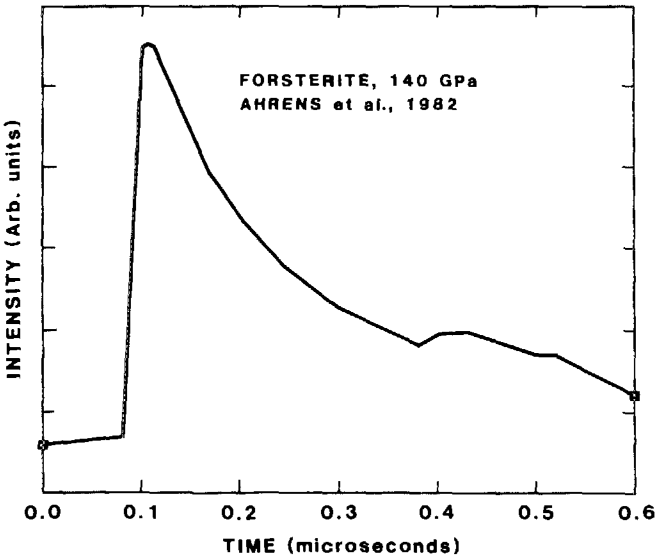


FIG. 6. Time dependence of 790-nm light radiated from forsterite shocked to 140 GPa (Ahrens *et al.*, see Ref. 6).

depth of the shocked material. For an initially transparent sample ($a_u \approx 0$) in the limit as τ goes to zero, Eq. (7) becomes

$$I(t) = r\left[f(T_s)(1 - e^{-a_s(U_s - u_p)t}) + \epsilon_i f(T')e^{-a_s(U_s - u_p)t}\right]. \tag{12}$$

An example of this type of behavior is illustrated in Fig. 2(a), and compared to data for NaCl and forsterite[6] in Figs. 5 and 6, respectively.

G. Time-dependent case 5

Another frequently observed case, which comes closest in appearance to the ideal case, occurs when the interface temperature T_i is approximately equal to the shock temperature ($T_i \approx T_s$) and the interface emissivity is less than unity ($\epsilon_i < 1$). For an initially transparent sample ($a_u \approx 0$), Eq. (7) becomes

$$I(t) = r\left[f(T_s)(1 - e^{-a_u(U_s - u_p)t}) + \epsilon_i f(T_s)e^{-a_s(U_s - u_p)t}\right]. \tag{13}$$

This is illustrated in Fig. 2(c), along with quartz data[7] in Fig. 7 for comparison.

III. APPLICATION OF MODEL

A. Shock and rarefaction in MgO

This model can be applied, in modified form, to the shock luminescence data of Schmitt and Ahrens[17] on MgO. Their wavelength-resolved data were obtained using a 500-channel optical multichannel analyzer (OMA). This instrument does not provide time resolution, so an unfiltered silicon photodiode was also used to record the luminescence history, yielding time-resolved data, but no spectral information. Data from experiment MgO 1, at 60.5 GPa, are plotted in Fig. 8.

It is instructive to make use of an x-t diagram (Fig. 9) to illustrate the configuration of that experiment. The initially transparent MgO specimen was shock loaded with an asymmetric (tungsten on copper) impact. The copper driver was sufficiently thin that a rarefaction had time to reflect off the MgO interface, off the impact interface, and back to the MgO interface, sending a rarefaction wave into the MgO about 280 ns after the initial shock. Because the shock-transit time is about 430 ns, there is 150 ns during which both a shock and a rarefaction exist in the sample.

The behavior to 280 ns can be described by time-dependent case 1 [Eq. (9)]. However, in their wavelength-resolved data, Schmitt and Ahrens found that the spectral dependence is not Planckian (graybody), with a λ dependence given by Eq. (2). Instead the radiation is band type; the light appears to be radiated in discrete wavelength bands. For sufficiently small a_s, Eq. (9) becomes linear

$$I(t) \propto (U_s - u_p)a_s t, \quad 0 < t < 280 \text{ ns}.$$

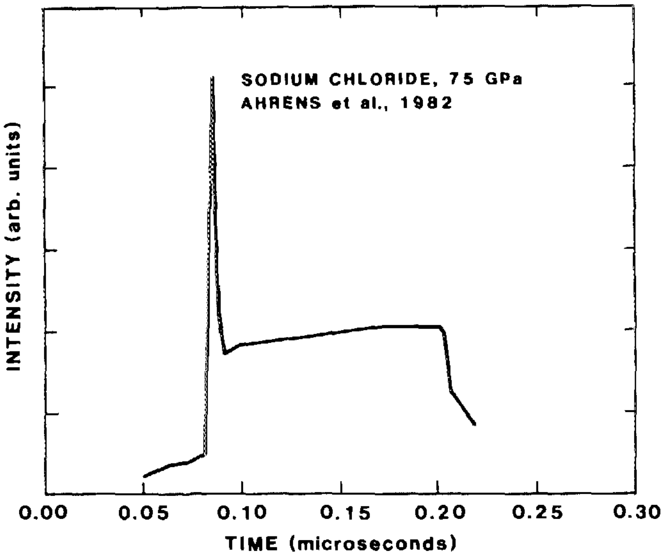


FIG. 5. Time dependence of 650-nm light radiated from NaCl shocked to 75 GPa (Ahrens *et al.*, see Ref. 6).

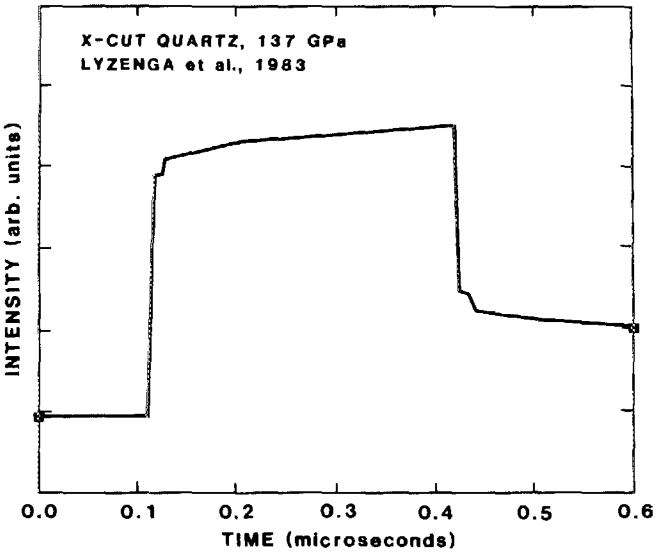


FIG. 7. Time dependence of 790-nm light radiated from quartz shocked to 137 GPa (Lyzenga *et al.*, see Ref. 7).

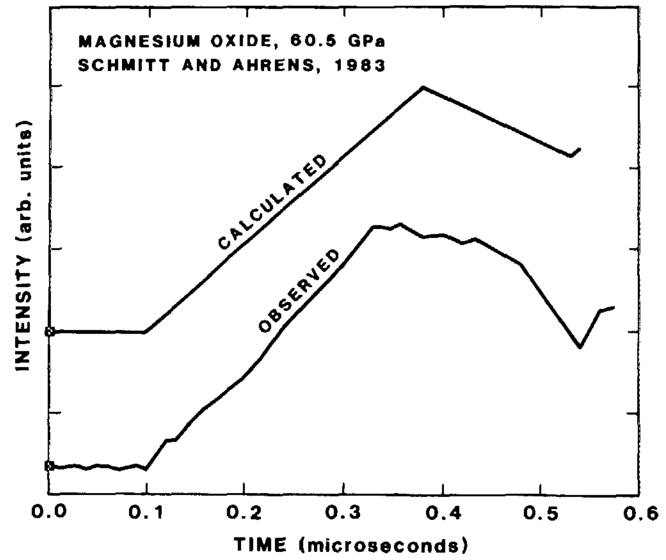


FIG. 8. Time dependence of unfiltered light radiated from MgO shocked to 60.5 GPa (Schmitt and Ahrens, see Ref. 17) compared to that calculated by model presented here.

Assuming that the MgO does not radiate from behind the first rarefaction characteristic, the radiating, shocked layer begins to decrease in thickness after 280 ns and is described by

$$I(t) \propto (U_s - u_r) a_s t, \quad 280 < t < 430 \text{ ns},$$

where u_r is the velocity of the first rarefaction characteristic. The calculated luminescence history is plotted in Fig. 8, showing remarkable agreement with the data for such a simple model.

B. Determination of linear absorption coefficient

In time-dependent cases 1 and 4, the data can be inverted to determine a_s. From Eqs. (9) and (12), one can see that the characteristic decay time is $\tau = 1/a_s(U_s - u_p)$. Knowledge of U_s and u_p, and determination of τ from the time-resolved data is sufficient to obtain a_s as a function of shock pressure. In experiments with both time- and wavelength resolved data, a_s can be determined as a function of wavelength. Making use of Kirchoff's radiation law, one can then obtain the emissivity as a function of wavelength and determine whether the shocked material demonstrates band or Planckian emission. This information is independent of whether the absolute spectral radiances are determined or not.

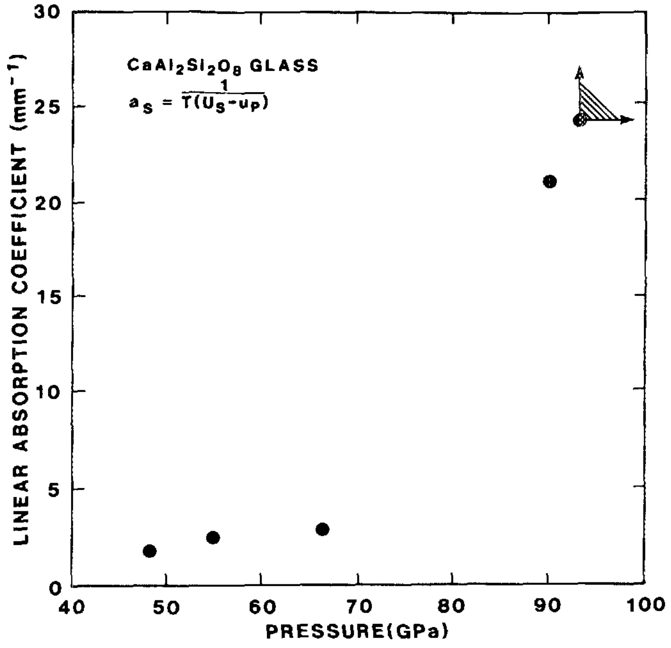


FIG. 10 Linear absorption coefficient as a function of shock pressure determined for anorthite glass by Boslough *et al.* (see Ref. 10).

As an example, the linear absorption coefficient at

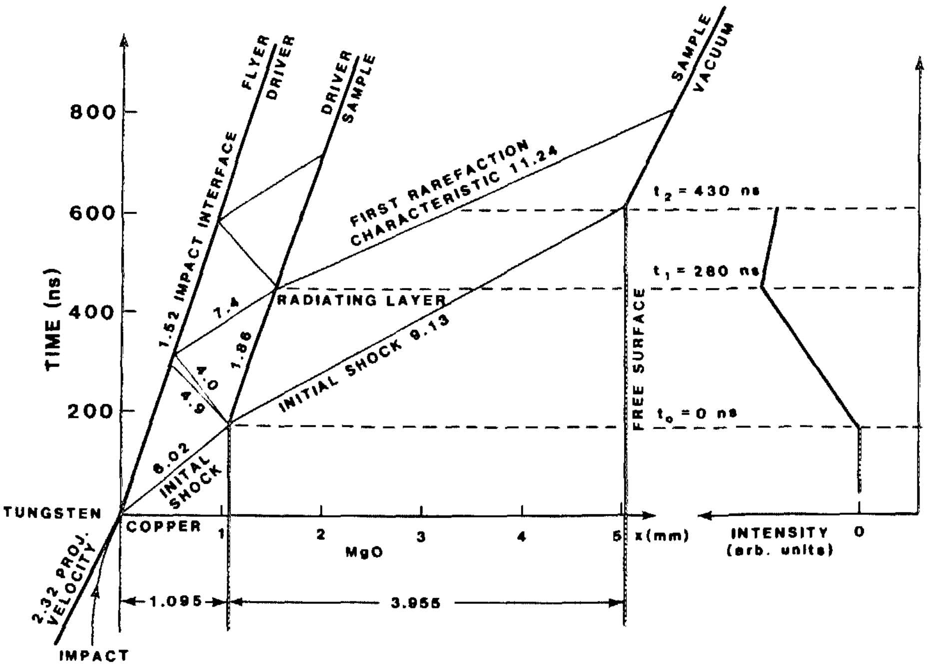


FIG. 9. Distance-time (*x-t*) diagram for MgO experiment of Schmitt and Ahrens. The numbers represent shock and sound velocities in laboratory coordinates in km/s.

$\lambda = 650$ nm can be determined for NaCl shocked to 75 GPa. The $1/e$ decay time can be determined from Fig. 5 to be about 1.5 ns. The shock and particle velocities for NaCl at this pressure are about 9 and 4 mm/μs, respectively.[18] Thus, a_s is about 13 cm^{-1}. Kormer *et al.*[2] determined a_s at two similar pressures; 70.5 and 79.0 GPa, and at two wavelengths: 478 and 625 nm. At both pressures they found a_s ($\lambda = 478$ nm) $= 12\ \text{cm}^{-1}$ and a_s ($\lambda = 625$ nm) $= 10\ \text{cm}^{-1}$, comparable to the value from this study.

Linear absorption were also determined in this way for shocked anorthite glass by Boslough *et al.*[10] for anorthite, a_s is independent of wavelength, which is consistent with the fact that the spectral radiance data could be fit very well with a Planck function. The absorption coefficient is an increasing function of shock pressure (Fig. 10), similar to what was found for NaCl.[2] This is consistent with the fact that the higher defect densities and more excited electronic states that cause larger absorption coefficients generally accompany higher-pressure shock states. Unpublished time-dependent data of Ahrens, Lyzenga, Mitchell, and Nellis indicate that quartz behaves similarly.

IV. SUMMARY

A model has been developed which describes the time dependence of light emission from shocked materials under certain conditions. The model is consistent with a wide variety of shock data, and can be used to determine the wavelength and pressure dependence of linear absorption coefficients of shocked materials. It may thus be possible to use the time dependence of light emission as a tool for measuring band edges and band gaps of shocked materials, or to locate shock pressure at which phase transformations occur, if they are accompanied by a sudden change in absorption coefficient. The model is, however, of limited use in the cases where there is an intrinsic time dependence of light emission, due to changes in temperature or emissivity of the shocked material.

ACKNOWLEDGMENTS

I would like to thank T. J. Ahrens, G. A. Lyzenga, A. C. Mitchell, W. J. Nellis, and D. R. Schmitt for allowing me to use their unpublished data. I appreciate the helpful discussions with R. Grover, N. Holmes, P. A. Taylor, S. C. Schmidt, and P. A. Urtiew, and thank A. C. Mitchell for getting me interested in the problem. This work was performed at Sandia National Laboratories supported by the U.S. Dept. of Energy under contract No. DE-AC04-76DP00789.

[1] S. B. Kormer, M. V. Sinitsyn, G. A. Kirilov, and V. D. Urlin, Sov. Phys. JETP **21**, 689 (1965).

[2] S. B. Kormer, M. V. Sinitsyn, G. A. Kirilov, and L. T. Popova, Sov. Phys. JETP **22**, 97 (1966).

[3] S. B. Kormer, Sov. Phys. Usp. **11**, 229 (1968).

[4] G. A. Lyzenga and T. J. Ahrens, Geophys. Res. Lett. **7**, 141 (1980).

[5] G. A. Lyzenga, T. J. Ahrens, W. J. Nellis, and A. C. Mitchell, J. Chem. Phys. **76**, 6282 (1982).

[6] T. J. Ahrens, G. A. Lyzenga, and A. C. Mitchell, *High-Pressure Research in Geophysics*, edited by S. Akimoto and M. H. Manghnani (Center for Academic Publications, Tokyo, 1982).

[7] G. A. Lyzenga, T. J. Ahrens, and A. C. Mitchell, J. Geophys. Res. **88**, 2431 (1983).

[8] W. J. Nellis, F. H. Ree, R. J. Trainor, A. C. Mitchell, and M. B. Boslough, J. Chem. Phys. **80**, 2789 (1984).

[9] M. B. Boslough, T. J. Ahens, and A. C. Mitchell, J. Geophys. Res. **89**, 7845 (1984).

[10] M. B. Boslough, T. J. Ahrens, and A. C. Mitchell, Geophys. J. R. Astron. Soc. (1985) (in press).

[11] G. A. Lyzenga and T. J. Ahrens, Rev. Sci. Instrum. **50**, 1421 (1979).

[12] M. B. Boslough, Ph.D. Thesis, California Institute of Technology, 1983 (unpublished).

[13] T. R. Harrison, *Radiation Pyrometry and its Underlying Principles of Radiant Heat Transfer* (Wiley, New York, 1960).

[14] J. H. Gladstone and T. P. Dale, Philos. Trans. **153**, 317 (1863).

[15] P. A. Urtiew and R. Grover, J. Appl. Phys. **45**, 140 (1974).

[16] H. S. Carslow and J. C. Jaeger, *Heat Conduction in Solids* (Oxford University, Oxford, 1959).

[17] D. R. Schmitt and T. J. Ahrens, Geophys. Res. Lett. **10**, 1077 (1983).

[18] S. P. Marsh, editor, *LASL Shock Hugoniot Data* (University of California, Berkeley, 1980).