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On the instability parameters of stellar atmospheres

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Abstract. We discuss the significance of the adiabatic exponents Γ_1, Γ_2 and Γ_3 and derive expressions for their calculation in a stellar atmosphere including simultaneous single-ionization of various elements and the presence of an equilibrium radiation field. A discussion is given of the relation of Γ_1 to dynamic instability. It is shown that some parts of some Kurucz models for extreme supergiant atmospheres are dynamically instable as a result of ionization and radiation in the deeper layers.

Key words: gas dynamics – chemical reactions – stars:atmospheres of – stars:supergiant

1. Introduction

In a gas at high temperature and low density, such as atmospheres of supergiants where the density is of the order of 10^{-8} g cm^{-3}, there are very little interactions between particles, and the pressure P, number density n and temperature T are related by the equation of state of a perfect gas. We consider a mixture of various elements, each of which obeys this equation $P_i = n_i kT$, but in which n_i may change as a result of ionization. The equation of state can be written,

$$P_{\text{gas}} = (1 + \bar{x})\rho NkT\,. \tag{1}$$

We consider an enclosure of (specific) volume V, where $V = \frac{1}{\rho}$ and ρ is the mass-density, at temperature T, containing N atoms. In Eq. (1) $\bar{x} = \sum_i \nu_i x_i$ denotes the mean degree of ionization, x_i the degree of ionization of element i, and ν_i the element abundance, where $\sum_i \nu_i = 1$.

The mixture is assumed to be always in thermal and chemical equilibrium. This requires the assumption that the sum of neutrals and ions of each element remains constant (mass conservation), besides the constancy of the element abundances.

We also include the effects of black-body radiation, and neglect the effects of any interactions between gas and radiation. We assume also a static medium, which means : no time-dependence or hydrodynamic motions. The radiation pressure is given by

$$P_{\text{rad}} = \frac{\text{a}}{3}T^4\,, \tag{2}$$

where a is the radiation constant.

Free electrons from elements with low ionization potentials will determine the ionization of hydrogen through the electron density in the Saha equation. Therefore, certain species of atoms can be very significant in stellar atmospheres even though their abundance is very low as compared to that of hydrogen. In particular, variations in metallic abundance from one star to another may have major effects on the physical state of the material, especially its ionization equilibrium (Bowers & Deeming 1984). For this reason, we take the elements heavier than helium also into account in the equation of state. On the other hand, we will restrict our calculations only to single ionization.

2. Adiabatic exponents

Generalized adiabatic exponents Γ_1, Γ_2 and Γ_3 are useful for the study of adiabatic processes in an ionizing gas. They are related to the behaviour of thermodynamic systems under infinitesimal adiabatic changes. These changes are isentropic (entropy S is constant) and hence reversible. In adiabatic sound waves, the temperature of the elementary volumes (typical one wavelength) changes accordingly with the variation in pressure as a result of the change in the density of matter. Another important feature of adiabatic processes is that all forces are conservative (absence of dissipative effects such as viscosity).

In sufficiently simple physical systems (i.e. non-ionizing perfect gases) , the ratio of specific heats equals Γ_1 (Cox & Giuli 1968), where

$$\Gamma_1 \equiv \left(\frac{\partial \ln P}{\partial \ln \rho}\right)_S = \frac{C_p}{C_v}\chi_\rho\,, \tag{3}$$

and $\chi_\rho = 1$ (density exponent, which is related to the coefficient of adiabatic compressibility of the gas). It is of importance to notice that Eq. (3) results from a general relation in thermodynamics that relates thermodynamic properties (ratio of specific heats) with mechanical properties (ratio of isothermal and adiabatic compressibility) of the system, and which is given by

$$\frac{C_p}{C_v} = \frac{\kappa_T}{\kappa_S}\,, \tag{4}$$

where

$$\kappa_T \equiv \left(\frac{\partial \ln \rho}{\partial P}\right)_T\,,\ \kappa_S \equiv \left(\frac{\partial \ln \rho}{\partial P}\right)_S\,, \tag{5}$$

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are the coefficient of isothermal and adiabatic compressibility, respectively (Mihalas & Mihalas 1984).

The gammas are important for the study of stellar interiors. Γ_1 is important in connection with dynamical instability of stars. The second adiabatic exponent is related to the differentials of state variables by the definition

$$\frac{\Gamma_2}{\Gamma_2 - 1} \equiv \left(\frac{\partial \ln P}{\partial \ln T}\right)_S = \frac{C_p}{C_p - C_v}\chi_T\,, \tag{6}$$

where χ_T is the temperature exponent. Γ_2 is important in connection with convective instability in stars. The third adiabatic exponent is defined by

$$\Gamma_3 - 1 \equiv \left(\frac{\partial \ln T}{\partial \ln \rho}\right)_S = \frac{C_p - C_v}{C_v}\frac{\chi_\rho}{\chi_T}\,, \tag{7}$$

and is important in connection with pulsational instability of stars.

In the case of a simple perfect gas with constant mean molecular weight μ, all gammas are the same and equal to the ratio of specific heats ($\frac{5}{3}$ for a monoatomic gas) because $\chi_\rho = \chi_T = 1$.

When black-body radiation is also included $\Gamma_1 > \Gamma_3 > \Gamma_2$, and all of them vary between $\frac{5}{3}$ and $\frac{4}{3}$. For monoatomic gas, $\Gamma_{1,2,3} = \frac{5}{3}$, while for pure radiation $\Gamma_{1,2,3} = \frac{4}{3}$. Without ionization, the density exponent $\chi_\rho = \beta$ (Chandrasekhar 1939) and the temperature exponent $\chi_T = 4 - 3\beta$, where

$$\beta \equiv \frac{P_{\rm gas}}{P_{\rm gas} + P_{\rm rad}}\,. \tag{8}$$

Now, considering a gaseous mixture with simultaneous ionization of the various elements and equilibrium radiation, Mihalas 1965 derived expressions for the specific heats, which take into account these various abundances. In order to calculate the gammas (and also the local speed of sound), we derived expressions for the density and temperature exponent, for this system;

$$\chi_\rho \equiv \left(\frac{\partial \ln P}{\partial \ln \rho}\right)_T = \frac{\beta\left(\bar{x}^2 + \bar{x} + \sum_i \nu_i x_i (1 - x_i)\right)}{(1+\bar{x})\left(\bar{x} + \sum_i \nu_i x_i (1 - x_i)\right)}\,, \tag{9}$$

and

$$\chi_T \equiv \left(\frac{\partial \ln P}{\partial \ln T}\right)_\rho = (4 - 3\beta) + \frac{\beta\bar{x}\sum_i \nu_i x_i (1 - x_i)(\frac{3}{2} + \frac{\chi_i}{kT})}{(1+\bar{x})(\bar{x} + \sum_i \nu_i x_i (1 - x_i))}\,, \tag{10}$$

where χ_i denotes the first ionization potential of each element. These expressions reduce to those of Cox & Giuli 1968 for a gas consisting of one ionizing element with radiation. It is also possible to derive more general expressions for C_p, C_v, χ_ρ and χ_T, including different stages of ionization for a mixture of various elements (for which it can be shown that they reduce to Mihalas' expressions for C_p and C_v and the χ_ρ and χ_T here presented).

Figure 1 shows that Γ_1 and $\Gamma_3 - 1$ reach a local minimum, nearly at the temperatures corresponding to 50 per cent ionization of H and He. The first local minimum is the result of the H ionization and the second of the He ionization. The curves presented here are calculated using solar abundances for 16 elements H, He, C, N, O, Ne, Na, Mg, Al, Si, S, Ar, K, Ca, Cr and Fe. The abundances were taken from Allen (1973). Partition functions at various temperatures were derived according to methods developed by Claas (1951) and Baschek, Holweger and Traving (1966). The mean molecular weight of unionized cosmic material (μ_0=1.26) follows from the abundances.

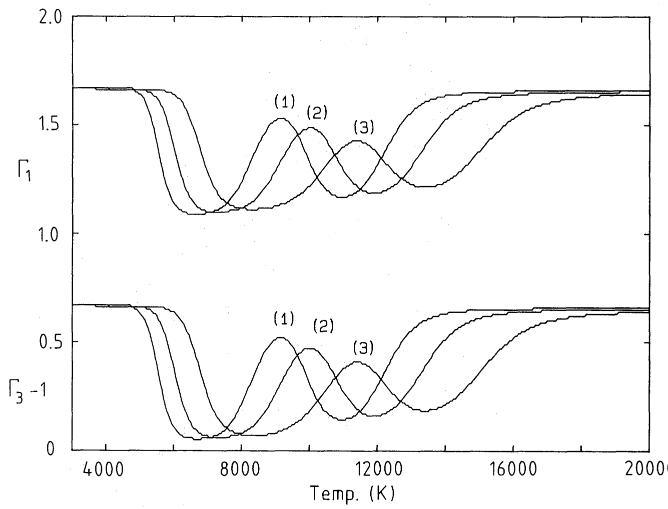

Fig. 1. Generalized adiabatic exponents Γ_1 and $\Gamma_3 - 1$ with ionization only, for three different electron pressures: (1) $P_{\rm e} = 0.1$, (2) $P_{\rm e} = 1.0$, (3) $P_{\rm e} = 20.0$

In an ionization region the mean molecular weight $\mu = \frac{\mu_0}{1+\bar{x}}$ is approximately halved because the released ionization-electrons double the total amount of free particles, assuming single-ionization. In this region, most of the work of compression is transformed into ionization energy, rather than into kinetic energy of thermal motion, so that the temperature will increase less upon compression than it would in a neutral or fully ionized zone.

From the principle of energy equipartition for perfect gases without radiation, it is shown that

$$\frac{C_p}{C_v} = 1 + \frac{r}{f}\,, \tag{11}$$

where $r = 2$ for a non-relativistic gas, and $r = 1$ for a fully relativistic gas. For three-dimensional non-relativistic gases, with the degree of freedom $f = 3$, it follows that $\Gamma_1 \simeq \frac{5}{3}$, and for a fully relativistic gas $\Gamma_1 \simeq \frac{4}{3}$ (outside the ionization regions $\chi_\rho \simeq 1$ in Fig. 3). In the hydrogen ionization zone of Fig. 1, Γ_1 drops to very close near unity which corresponds to a large value of f. When $f \to \infty$, physically, this would correspond to the situation for which *all* energy added to the system goes into 'internal' (i.e. ionization, rotation, vibration, etc.) forms of energy and *none* into kinetic energy.

In absence of radiation it follows from Eq. (8) that β=1, and for this situation of simulateous single-ionization of various elements, we find that $0.89 \le \chi_\rho \le 1$ (Fig. 3). In this case, the ratio of specific heats stays always enough above 1 to give for the product with χ_ρ a value for Γ_1 above unity.

3. Behaviour of Γ_1 when including equilibrium radiation

The sudden decrease of Γ_1 in an ionization region is caused by the low pressure gradient as compared to the density gradient there. This results from energy going into ionization, while the kinetic energy, and hence the gaspressure, increases much less than when compared to the situation outside this region. Upon further compression, without ionization, the pressure gradient

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will raise again, because all added energy returns back into the kinetic energy of ions *and* released electrons. Hence, the adiabatic compressibility κ_S will decrease again.

However, when including equilibrium radiation, things change. Radiation is more compressible than pure gas with the same pressure and temperature. This is because a change of the specific volume upon compression will cause a much smaller temperature-increase for radiation than it would cause for a gas under the same conditions. Since the radiation pressure only depends on the temperature in Eq. (2), the change in this pressure will be accordingly smaller.

Suppose now that the radiation pressure becomes comparable to the gas pressure in an ionization region. The temperature changes little upon compression until the ionization is over, but at that moment the radiation is preventing the total pressure to increase again while the compression goes on. So, in Eq. (3) the density gradient is much larger as compared to the gradient in the total pressure, than it would be if radiation were absent. Consequently, Γ_1 can drop to a lower value than in the absence of radiation. (Also notice in Fig. 3 that $\chi_\rho \to 0$ when the radiation pressure becomes more important than the gas pressure, or $\beta \to 0$). This effect can be seen in Fig. 2 where Γ_1 *can even drop to below 1 in certain ranges of electron pressure and temperature.* We stress that this is only a result of the combination of radiation *and* ionization, in a certain range of P_e and T -values. The effect appears for $0.1 \frac{\mathrm{dyn}}{\mathrm{cm}^2} \leq P_e \leq 20.0 \frac{\mathrm{dyn}}{\mathrm{cm}^2}$ and $6100 \mathrm{~K} \leq T \leq 9000 \mathrm{~K}$ in the hydrogen ionization region for solar abundance. We find a minimum value of 0.84 for Γ_1.

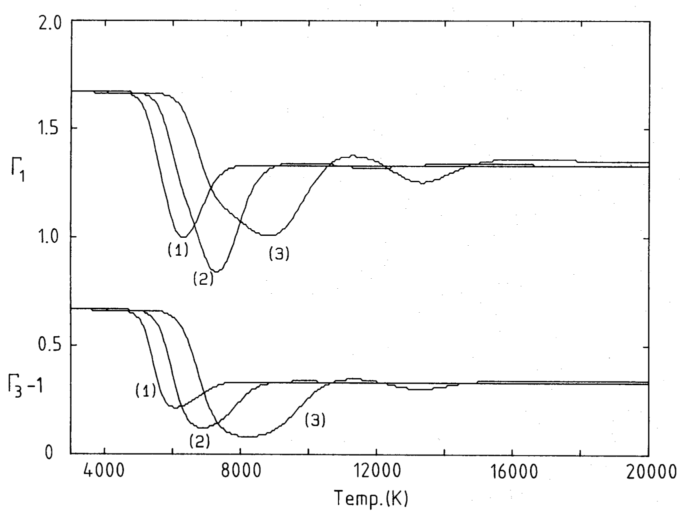

Fig. 2. Generalized adiabatic exponents Γ_1 and $\Gamma_3 - 1$ with ionization and equilibrium radiation, for three different electron pressures: (1) $P_e = 0.1$, (2) $P_e = 1.0$, (3) $P_e = 20.0$

For a perfect gas, this would correspond to the somewhat surprising situation of a negative polytrope, because then $n < 0$ in

$$1 + \frac{1}{n} = \left(\frac{\partial \ln P}{\partial \ln \rho}\right)_S \equiv \Gamma_1 \tag{12}$$

However, when partial ionization or radiation is included, it is shown (Cox & Giuli 1968) that expression (12) has to be replaced by

$$\Gamma_1 = \chi_\rho \left(\frac{1 + n_e}{1 + n_e - \chi_T}\right), \tag{13}$$

where n_e is now called the 'effective polytropic index'. In Fig. 3 we calculate n_e from Eq. (13) as a function of the temperature with an electron pressure of $1.0 \frac{\mathrm{dyn}}{\mathrm{cm}^2}$. Here, n_e always remains positive, even in the region where Γ_1 has values below unity (Fig. 2).

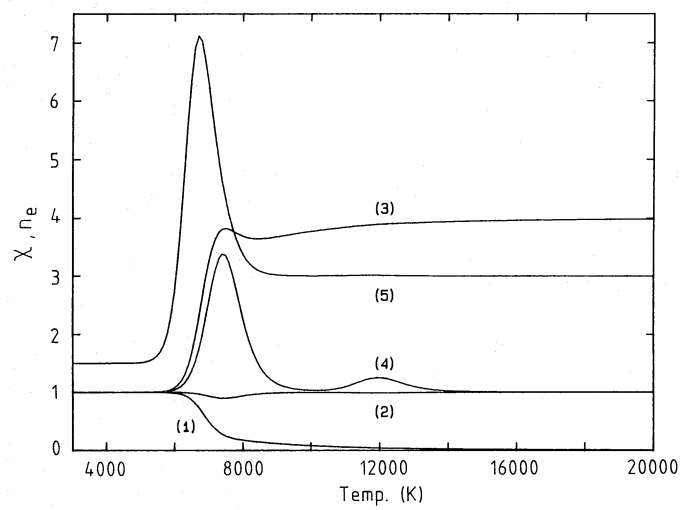

Fig. 3. Density exponent χ_ρ (with equilibrium radiation (1) and without (2)) and temperature exponent χ_T (with radiation (3) and without (4)) as a function of the absolute temperature, in the case of an ionizing gas where $P_e = 1.0$. (5) shows the effective polytropic index n_e when including radiation

4. Γ_1 in relation to atmospheric dynamical instability

Radial adiabatic oscillations about a mean state of hydrostatic equilibrium with conservation of mass can be described by a linear second order differential equation : the Linear Adiabatic Wave Equation (LAWE). It describes a differential equation for a radial displacement $\frac{\delta r}{r}$ caused at a point r by perturbations under certain boundary conditions (δr vanishes at the centre of the star and δP vanishes at the surface). It is assumed that its standing wave solutions are of the form $\zeta(r,t) = \frac{\delta r}{r} = \xi(r)e^{i\sigma t}$, where $\xi(r)$ is a function of r only and σ is a constant.

The simplest solution for this equation is given when assuming adiabatic motion of the fundamental mode of radial oscillation. When $\sigma \neq 0$, the perturbation has an oscillatory part for either sign of σ. This mode represents the case where perturbations vanish at the centre and reach maximum values at the surface. It has the smallest eigenvalue σ_0^2 of an infinite discrete set σ_i^2 corresponding to the eigensolutions ξ_i.

Assuming a homologous motion with $\xi(r)$ constant, it can be shown (Ledoux 1965) that

$$\sigma_0^2 = \frac{\int_0^R (3\Gamma_1 - 4) 3P dV}{\int_0^M r^2 dm} = <\Gamma_1 - \frac{4}{3}> \frac{3\mid\Omega\mid}{I}, \tag{14}$$

where I is the generalized moment of inertia with regard to the centre of the star,$\mid\Omega\mid$ the gravitational potential energy in its equilibrium state, P the unperturbed pressure and Γ_1 is a function of the position in the star.

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When $\sigma_0^2 > 0$ the system is stable because $\frac{\delta r}{r}$ does not grow with time. Stability will depend on the volume-averaged value of Γ_1, and vanishes when this average $\bar{\Gamma}_1 \leq \frac{4}{3}$ or $\sigma_0^2 \leq 0$.

In extensive ionization zones of stellar evelopes, Γ_1 falls below $\frac{4}{3}$ through a significant part of the atmosphere (see next paragraph). Hence, local dynamical instabilities in atmospheric layers may result from H and He ionizations. Ledoux 1965 found that regions where $\Gamma_1 < \frac{4}{3}$ and decreases with increasing radius, have a destabilizing influence. Regions where $\Gamma_1 < \frac{4}{3}$ and increases with increasing radius, have a stabilizing influence. If Γ_1 is higher than $\frac{4}{3}$, the effects of its variations are reversed.

In parts of the atmosphere where $\Gamma_1 > \frac{4}{3}$ and the layers are contracting under a disturbance, its potential energy becomes less negative and is mainly transformed into kinetic energy of the particles (and partly radiated because of the increasing temperature). This leads to a counterworking pressure (low compressibility) that tries to restore the equilibrium of the layer.

Conversely, when $\Gamma_1 < \frac{4}{3}$ in a layer, the potential energy is mainly transformed into 'internal forms of energy' (i.e. ionization energy, dissociation energy of molecules), without increasing the pressure to stop the contraction. This means that for an arbitrary displacement of atmospheric layers from equilibrium, there are forces which tend to push them still further from equilibrium. These forces, that act on a mass shell, when it is displaced from equilibruim, are determined not only locally, but also by conditions in the entire star. It therefore follows, for example, that in some region where Γ_1 is less than $\frac{4}{3}$, this mass shell will still be dynamically *stable if* the whole star is stable or $\sigma_0^2 > 0$ (Cox 1980). This means that it will experience positive restoring forces if displaced from equilibruim. In that case, pulsational driving can occur from the γ-mechanism because $\Gamma_3 - 1$ has very low values in these regions (Fig. 1). This driving is produced by recombination on decompression that releases the ionization energy (Cox & Cahn 1988).

The destabilizing mechanism in the H- and He-ionization zones is considered here because Dziembowski 1977 concluded in his theoretical studies that non-radial oscillations could be excited in the envelopes of several kinds of stars, by the same envelope ionization mechanisms that are responsible for the pulsations of the Cepheids and RR Lyrae variables. From an investigation of a supergiant model for α Cygni (log $T_{\rm eff}$=3.96, logg=1.1) he found that this model is overstable to high-l non-radial p-mode oscillations. Lower-l modes were excited in a cooler supergiant, lying close to the Chepeid instability strip. Ando 1976 performed numerical calculations for non-radial p-modes ($l \geq 10$) in the envelope of four supergiant models ($T_{\rm eff} < 6500$). He found that high-order p-modes ($k \geq 5$) cannot be trapped because the acoustic cut-off frequency is low, resulting from the large atmospheric scale height relative to the stellar radius. The stability of these eigenmodes strongly depends on the position of the driving ionzation region below the outermost node of the p-modes. De Jager, de Koter, Carpay & Nieuwenhuijzen 1990 concluded that gravity waves in supergiant atmospheres are fairly high-mode waves. Also from observations of early-type supergiants there appears to be an indication that these stars are non-radial pulsators (for references see de Jager 1990).

5. Application to the Kurucz models

On the basis of the considerations in Sects. 2 and 4, we want to determine whether parts of the atmosphere of the Kurucz models 1979 for extreme supergiants are dynamically stable agianst compression or rarefaction.

We calculated with Eq. (3) and Eq. (9) the variation of the stability parameter Γ_1 throughout the atmospheres of various Kurucz models. Taking only gas pressure into consideration and increasing the temperature from 5500 K to 8000 K, while choosing for the log g -value the lowest possible present in the tables, Fig. 4 shows that Γ_1 is gradually decreasing in the outer parts of the atmosphere. This is caused by the hydrogen ionization region which is moving outwards at higher temperatures. The first ionization region of He is also visible in the deeper layers, as is shown by the sudden decrease of Γ_1. This minimum is less deep because of the lower abundance of He.

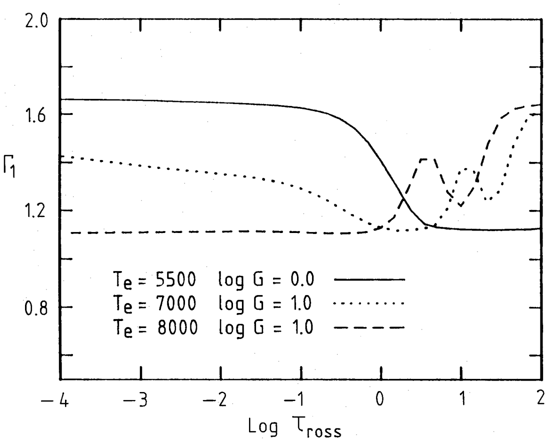

Fig. 4. Γ_1 throughout the atmospheres of extreme supergiants for three different effective temperatures, using Kurucz models with solar abundance. Only a single-ionizing gas is considered

When further increasing the effective temperature to 13000 K, as is shown in Fig. 5, Γ_1 increases again in the outer parts of the atmosphere, which indicates that all hydrogen is already ionized at the surface. From this graph it can be seen that the He ionization region is moving outwards with increasing effective temperature.

In Fig. 6 the effective temperature is increased further to 20000 K. At the temperatures of 14000 K and 15000 K we find that the region of single ionization of He is extending very rapidly towards the surface. He ionization appears to drop Γ_1 to below $\frac{4}{3}$ in a large part of the atmosphere, comparable to the effect of the ionization of H at lower effective temperatures. The effective temperature has to be increased up to 20000 K to find Γ_1 -values above $\frac{4}{3}$ over the whole atmosphere. In these atmospheres H and He are completely singly ionized.

From this we can conclude that parts of the Kurucz models with lowest gravitational acceleration show an enhanced sensitivity to atmospheric dynamical instabilities. This is caused by H ionization for models between 7000 K and 12000 K, and by single He ionization for model atmospheres between 15000 K and 20000 K.

The question if these models are unstable can only be answered by evaluation of Eq. (14). At large optical depths (log $\tau_{\rm Ross} > 0$) $\Gamma_1 > \frac{4}{3}$, and since the pressure is high there, this can give a large contribution to the integral of Eq. (14). Consequently, it will reduce the instability of the atmosphere. However, when radiation pressure is included in the calculations, it appears from Fig. 7 that in the deeper layers Γ_1 decreases to a value close to $\frac{4}{3}$,

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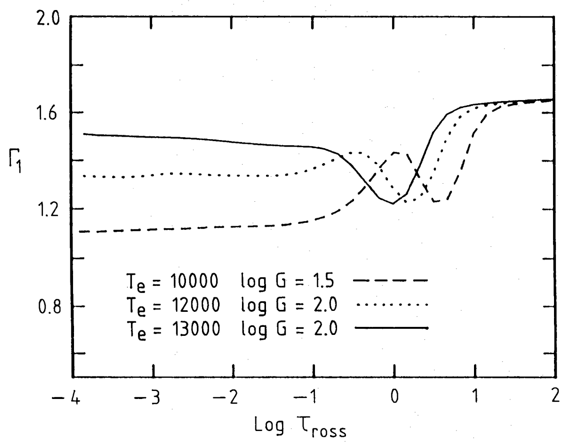

Fig. 5. Γ_1 throughout the atmosphere of extreme supergiants, using Kurucz models with solar abundance. Only a single-ionizing gas is considered and effective temperatures between 10000 K and 13000 K

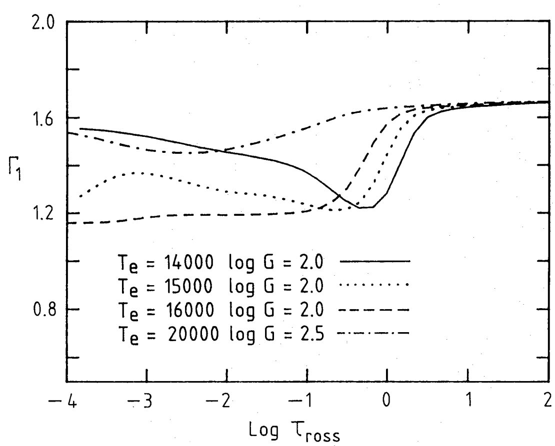

Fig. 6. Γ_1 throughout the atmosphere of extreme supergiants, using Kurucz models with solar abundance. Only a single-ionizing gas is considered and effective temperatures between 14000 K and 20000 K

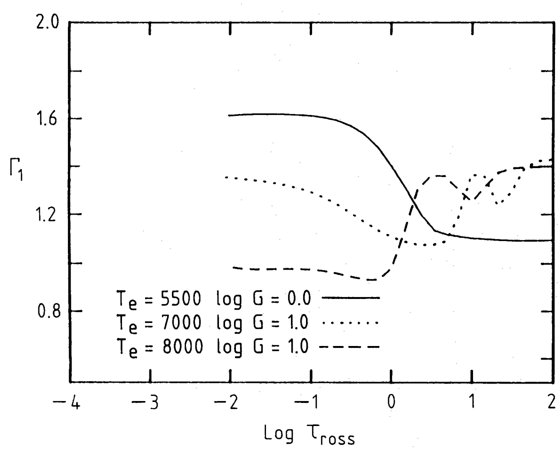

Fig. 7. Γ_1 in the deeper layers of atmospheres for extreme supergiants, using Kurucz models with solar abundance. The total effect of single-ionizing gas and black-body radiation is considered, but notice that the inclusion of the radiationterm is not correct for $\tau_{\mathrm{Ross}} \leq 1$

which will increase the instability. This graph is not valid at low optical depths because there the radiation cannot be treated as black-body radiation. An equilibrium between gas and radiation is then not possible.

We remark that the presence of a non-equilibrium radiation field at very low densities makes the material non-adiabatic (NLTE conditions) and thermodynamic quantities become meaningless (Mihalas & Mihalas 1984). However, Ledoux & Walraven 1958, gave an extensive description of the stability theory applied to atmospheric layers of supergiant variables (like η Aquila with logg=1.0 and T_{eff}=6000 K) where thermodynamic quantities as Γ_1 are considered at very low densities.

In a series of three articles by Hummer & Mihalas 1988; Mihalas, et al. 1988; Dăppen et al. 1988, a much more sophisticated treatment of the thermodynamic properties of a partially ionized multicomponent gas is developed for stellar envelopes. In the third paper the thermodynamic quantities are calculated at these low densities, and here also values for Γ_1 are found to fall below $\frac{4}{3}$ in the thermal (single) ionization regions over the absolute temperature regions of Fig. 4, 5 and 6. Here, it is indirectly shown that the location of the ionization zones is pushed deeper down in the atmosphere because their formalism explicitly deals with excited states of H and He (Dăppen 1988). This results in an enhanced local instability of the deeper layers.

Notice in Fig. 7 that Γ_1 can drop below 1, even at fairly high optical depths ($\tau_{\mathrm{Ross}} \simeq 1$). When compared with Fig. 4, Γ_1 decreases from 1.1 to 0.9 as a result of the inclusion of black-body radiation. This corresponds to a change $\triangle\Gamma_1 \simeq 0.1$ (18 %) which directly affects the freqencies of simple adiabatic radial oscillations and the sound velocity.

Christensen-Dalsgaard & Thompson 1991 investigated the effect of increasing the solar envelope abundance on Γ_1 in the He^{+}-ionization region. Here, changes of $\triangle\Gamma_1 \simeq 0.004$ are found to be significant for the helioseismic determination of the helium abundance. This implies that a detailed knowledge of the contribution of equilibrium radiation, at the basis of stellar atmospheres, is indispensible since $\triangle\Gamma_1$=0.01 when it is only a hundreth of the gas pressure (β=0.99).

Notice also the steep gradient in Γ_1 which results from this effect. It corresponds to a sudden increase in the compressibility of the medium for sound waves moving outwards from a region where hydrogen is fully ionized to a region where this ionization takes place. On this edge the local velocity of sound suddenly decreases since

$$S_{\mathrm{ad}}^2 = \frac{P}{\rho}\Gamma_1\,, \tag{15}$$

(Landau & Lifshitz 1959) which enhances the origin of shock waves.

6. Conclusions

There is a clear indication that the Kurucz atmospheric models for certain extreme supergiants are dynamically unstable in certain ranges of optical depths. This is mainly caused by the ionization of the abundant elements, combined with the effects of equilibrium radiation in the deeper parts of these atmospheres. We show that the first ionization of He reduces also the local stability over a large part of the atmosphere, similar to the effect of H ionization for models with lower effective temperatures. Since the contribution to the adiabatic exponents by simultaneous ionization of various elements and equilibrium radiation

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can be calculated directly, this allows for a better understanding of stability problems in general, but in particular also on dependence of the metallic abundance.

It is shown that the first adiabatic exponent Γ_1 can have values below unity, while this is expected not to be the case in many theoretical deductions for shock waves. This effect results from the combination of ionization and equilibrium radiation at high optical depths. It has to be related to strong dynamical or atmospherical instabilities under certain atmospherical conditions. From the Kurucz models of extreme supergiants, these conditions appear to be realistic.

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