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HEATING BY DISSIPATION OF SOUND WAVES IN THE INTERSTELLAR GAS

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ABSTRACT

The equilibrium resulting in a recombining plasma with arbitrary metallicity Z, heated by a mean radiation field E as well as by sound waves dissipation due to thermal conduction, dynamic and bulk viscosities is analyzed. Generally, the heating by acoustic waves' dissipation induces drastic changes in the range of temperature where the thermochemical equilibrium may exist. An additional equilibrium state appears which is characterized by a lower ionization and higher gas pressure than the equilibrium resulting when the wave dissipation is neglected. The above-mentioned effects are sensible to the values of the gas parameters as well as the wavelength and intensity of the acoustic waves. Implications in the interstellar gas, in particular, in the high velocity clouds are outlined.

Keywords: Recombining astrophysical plasma, interstellar medium (ISM), high velocity clouds (HvCs), metallicity, mean radiation field.

INTRODUCTION

It has been well established (Landau and Lifshitz, 1987; Mihalas *et al*., 1984; Stix, 1992) that when the gas dynamic equations are linearized, assuming small disturbances, a rotational mode and a potential mode result by the way that they are independent of each other. If viscosity is accounted for, it produces damping on the potential mode and if thermal conduction is taken into account an additional damped thermal mode appears (Landau and Lifshitz, 1987).

The above two dissipative effects have been invoked as one of the mechanisms of heat input in different astrophysical plasmas, particularly, in the solar one, and more generally in stellar atmospheric plasmas (Bird, 1964; Narain *et al*., 1990; Stein *et al*., 1972; Stein *et al*., 1974) in the interstellar medium (ISM) (Spitzer, 1978; Spitzer, 1982, Spitzer, 1990) and more recently in the intracluster gas (Fabian *et al*., 2003; Fabian *et al*., 2005; Ruszkowski *et al*., 2004; Ferland *et al*., 2009; Ibañez *et al*., 2005; Fajardo *et al*., 2021).

Most of the works on this subject does not consider bulk viscosity dissipation, none the less the corresponding dissipation can be larger than both the dynamic viscosity and the thermal conduction dissipation when chemical processes are present (including ionization and recombination processes) (Ibañez *et al*., 1993, 2019). In addition to the quantitative changes in the acoustic dissipation, the bulk viscosity also introduces important qualitative changes, in particular, due to the fact that the bulk viscosity coefficient is dispersive (depends on the wave frequency) as it will be shown later on.

The present work is aimed to analyze how much a thermochemical equilibrium is modified if heating by sound wave dissipation is taken into account in an homogeneous recombining plasma able to cool down by the well-known cooling function for plasmas with arbitrary metallicity Z, ionized and heated by a mean radiation field E as well as by dissipation of the acoustic waves. Implications in the structure of the interstellar medium (ISM) and in high velocity clouds (HvCs) will be outlined.

Inhomogeneity effects will be neglected and only three dissipative mechanisms will be considered, i.e. the thermal conduction κ, the dynamic viscosity η, and the bulk viscosity ζ. Molecular formation which appears once the hydrogen molecule H_2 forms (in gas phase or adsorption on cool grains) in particular, the formation of the CO molecule (which is a much more stronger coolant than the corresponding atomic cooling) as well as the strong cooling by solid grains and their important opacity effects will not be taken into account at the present approximation. These additional effects completely change the radiative heating and cooling of the interstellar gas. These complications as well as

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magnetohydrodynamic effects will be carried out elsewhere (Fajardo *et al*., 2021).

BASIC EQUATIONS

For a fluid where a chemical reaction can have the form $\Sigma_i b_i C_i = 0$, the equations of gas dynamics can be written in the following form:

$$\frac{\partial \rho}{\partial t} + \frac{\partial \rho v_i}{\partial x_i} = 0 \tag{1}$$

$$\rho \left(\frac{\partial v_i}{\partial t} + v_j \frac{\partial v_i}{\partial x_j} \right) = - \frac{\partial p}{\partial x_j} + \frac{\partial}{\partial x} \left[\eta \left(\frac{\partial v_i}{\partial x_j} + \frac{\partial v_j}{\partial x_i} - \frac{2}{3} \delta_{ij} \frac{\partial v_k}{\partial x_k} \right) \right] + \frac{\partial}{\partial x} \left(\zeta \frac{\partial v_j}{\partial x_j} \right) \tag{2}$$

$$\frac{d\,T}{d\,t} + X(\rho, T, \xi) = 0 \tag{3}$$

$$A(\xi) \frac{d\,T}{d\,t} - \frac{p}{\rho^2} \frac{d\,\rho}{d\,t} + RTB(\xi, T) \frac{d\,\xi}{d\,t} + \mathcal{L}(\rho, T, \xi) - \frac{1}{\rho} \boldsymbol{\nabla}.(\kappa\, \nabla\, T) = 0 \tag{4}$$

$$p = \frac{R\,\rho\,T}{\mu(\xi)} \tag{5}$$

where d/dt is the convective derivative, $X(\rho, T, \xi)$ is the net rate function, $L(\rho, T, \xi)$ the net rate of cooling per unit mass and time, which are functions of the degree of ionization ξ (parameter determining the advance of the reaction), the density ρ, and the temperature T. Also, κ is the thermal conduction coefficient, η is the dynamic viscosity, and ζ the bulk viscosity (complex parameter in this case). Furthermore, an ideal gas state in equation (5) has been assumed. R is the universal gas constant, and the coefficients A, B, and μ are defined as follows:

$$A(\xi) = \sum_i \frac{b_i\, \xi + x_i^0}{\gamma_i - 1}, \quad B(T) = \sum_i \left(\frac{b_i}{\gamma_i - 1} + \frac{1}{k_B\, T}\, b_i\, \epsilon_i^0 \right), \quad \frac{1}{\mu(\xi)} = \sum_i b_i\, \xi + x_i^0 \tag{6}$$

where x_i^0 is some initial concentration, γ_i and ϵ_i^0 are the specific heat ration and zero-point energy of the *i-esime* gas component, respectively, and k_B is the Boltzmann constant. Additionally, for a plasma with the ionization degree ξ, the thermal conduction becomes

$$\kappa = 2.5 \times 10^3\, (1 - \xi)\, T^{\frac{1}{2}} + 1.84 \times 10^{-5} \frac{\xi\, T^{5/2}}{\ln \Lambda_s(\rho, T, \xi)} \tag{7}$$

where $\ln[\Lambda_s(\rho, T, \xi)]$ is the relation between the Debye screening and the impact particles' parameter (Braginskii, 1965; Spitzer, 1962). The dynamic viscosity following the work by Braginskii (1965) and Spitzer (1962) takes the following form:

$$\eta = 2.21 \times 10^{-15} \frac{T^{5/2}}{\ln \Lambda_s(\rho, T, \xi)} \tag{8}$$

On the other hand, the complex bulk viscosity ζ becomes

$$\zeta = \frac{\rho\,\tau}{1 - i\omega\tau} [c_\infty^2 - c_0^2] \tag{9}$$

where τ is the chemical relaxation time

$$\tau = \left(\frac{\partial X(\rho, T, \xi)}{\partial \xi} \right)^{-1} \tag{10}$$

and the sound velocities c_∞^2 and c_0^2 are defined as follows:

$$c_\infty^2 = \left(\frac{\partial p}{\partial \rho} \right)_\xi \text{ and } c_0^2 = \left(\frac{\partial p}{\partial \rho} \right)_{eq} = \left(\frac{\partial p}{\partial \rho} \right)_\xi + \left(\frac{\partial p}{\partial \xi} \right)_\rho \left(\frac{\partial \xi_0}{\partial \rho} \right) \tag{11}$$

Here ξ_0 is the value of the chemical parameter at chemical equilibrium and the derivative $\partial\xi_0/\partial\rho$ is also at equilibrium (Landau and Lifshitz, 1987). In this work, the c.g.s. unit system is used.

LINEAR WAVES

If the velocity field of sound disturbances with the wavenumber k and the frequency ω is assumed to be $v = v_x = v_1 \cos(kx - \omega t)$, where t is time and $v_y = v_z = 0$ as a first approximation, the time averaged in some volume V of energy dissipation by the bulk viscosity ζ of sound waves becomes

$$\frac{1}{V} \int \overline{\zeta \frac{\partial v}{\partial x}}\, dV = \zeta_R \frac{v_1^2}{2} k^2 \tag{12}$$

And as can be realized easily, the real part ζ_R of the bulk viscosity ζ is

$$\zeta_R = \frac{\rho\,\tau}{1+(\omega\tau)^2}\left[c_\infty^2 - c_0^2\right] \tag{13}$$

On the other hand, we have that

$$c_\infty^2 - c_0^2 = \left(\frac{\partial p}{\partial \xi}\right)_\rho \tag{14}$$

where $(\partial\xi/\partial\rho)_{eq}$ is calculated at equilibrium. This result will allow some calculations of the bulk viscosity coefficients for reacting gases, as it will be shown later on.

Therefore, the heat inputs due to the dynamic viscosity as well as the bulk viscosity (Landau and Lifshitz, 1987) and the thermal conduction becomes

$$\Gamma_\omega(\rho,T,\xi) = \gamma_d\,\frac{v_1^2}{2}\,k^2 \tag{15}$$

where

$$\gamma_d = \left[\left(\frac{4}{3}\eta + \zeta_R\right) + \frac{(\gamma-1)\kappa}{c_P}\right] \tag{16}$$

and c_P is the specific heat at constant pressure.

THERMAL EQUILIBRIUM OF A PHOTOIONIZED GAS

Commonly the heat/loss function for a low-density plasma at equilibrium temperatures (neglecting molecular formation as quoted out in the introduction) in the range 30 [K] < T < 3 × 10^4 [K] can be written (dimension [erg × cm^{-3} × s^{-1}]) as follows:

$$\rho\mathcal{L}(\rho,T,\xi) = \Lambda(\rho,T,\xi) - \Gamma_0(\rho,T,\xi) - \Gamma_\omega(\rho,T,\xi) \tag{17}$$

where $\Lambda(\rho, T, \xi)$ is the cooling rate and $\Gamma_0(\rho, T, \xi)$ is the heating input (different from wave dissipation) per unit volume and time, respectively. So, at thermal equilibrium we have $L(\rho, T, \xi) = 0$, i.e. the cooling rate becomes

$$\Lambda(\rho,T,\xi) = \Gamma_0(\rho,T,\xi) + \gamma_d\,\frac{v_1^2}{2}\,k^2 \tag{18}$$

For an optically thin hydrogen plasma with the ionization rate ς and the metallicity Z heated and ionized by a background radiation field of mean photon energy E, the net rate function $X(\rho, T, \xi)$ and the cooling and heating rates per unit volume and time are respectively given by Corbelli *et al*. (1995) as follows:

$$X(\rho,T,\xi) = N_0\,\rho\,[\xi^2\alpha - (1-\xi)\,\xi\,\gamma_c] - (1-\xi)(1+\phi)\,\varsigma \tag{19}$$

$$\Lambda(\rho,T,\xi) = (N_0\,\rho)^2[(1-\xi)Z\Lambda_{HZ} + \xi Z\Lambda_{eZ} - (1-\xi)\Lambda_{eH} + \xi^2\Lambda_{eH^+}] \tag{20}$$

and

$$\Gamma_0(\rho,T,\xi) = N_0\,\rho\,(1-\xi)\,\varsigma\,[E_h + (1+\phi)\,\chi_h] \tag{21}$$

where ϕ denotes the number of secondary electrons, E_h the heat released per photoionization (Shull *et al*., 1985), Λ_{HZ}, Λ_{eZ}, Λ_{eH}, and Λ_{eH+} respectively are the cooling efficiencies by collisions of neutral hydrogen-ions and metal atoms (Launay *et al*., 1977; Dalgamo *et al*., 1972), electrons-ions and metal atoms (Dalgamo *et al*., 1972), *Lya* emission by neutral hydrogen (Spitzer, 1978), and hydrogen recombination, on the spot approximation (Seaton, 1959).

From the results obtained in the paper by Corbelli *et al*. (1995), it follows that at ionization equilibrium state $X(\rho, T, \xi) = 0$ we have

$$\left(\frac{\partial\xi}{\partial\rho}\right)_{eq} = -\frac{1}{2}\,\frac{(1+\phi)\,\varsigma\,\left[\gamma_c + 2\alpha - \sqrt{F} + {}^{(1+\phi)\,\varsigma}\!\big/_{(N_0\,\rho)}\right]}{(\alpha+\gamma_c)\,N_0\,\rho^2\,\sqrt{F}} \tag{22}$$

where

$$F = \gamma_c^2 + 2(\gamma_c + 2\alpha)\frac{(1+\phi)\,\varsigma}{N_0\,\rho} + \left(\frac{(1+\phi)\,\varsigma}{N_0\,\rho}\right)^2$$

On the other hand, the damping scale length l_d is given by

$$l_d = \frac{2\,\rho\,c_0}{\gamma_d\,k^2} \tag{23}$$

where equation (18) holds as far as the damping scale-length is larger than the sound wave length $\lambda = 2\pi/k$, i.e.

$$\frac{\lambda}{2\,\pi\,l_d} = \frac{\sqrt{\gamma_d\,\phi}}{\sqrt{2}\,\epsilon\,c_0} \tag{24}$$

where $\epsilon = v_1/c_0$ (notice that the parameter ϵ does not have units). So, the dimensions of the region heated by sound waves' dissipation is larger than the sound wavelength and the above-written approximation holds.

APPLICATIONS

Interstellar gas (ISM)

At the present subsection the results obtained previously will be applied to a gas with characteristic values of the parameters representative of the interstellar medium, i.e. $N_0\rho = 1$, $Z = 1$, and $E = 10^2$ [eV], and in the range of temperatures where the hydrogen ionization-recombination takes place.

Figure 1 shows the following parameters in dependence on the temperature T: the ionization ξ (a), the pressure $p/k_B\varsigma$ (b), the heat input Γ_0 (c), and the dissipative coefficients $(\gamma - 1)\kappa/c_P$ (the dash line in Figure 1d), $4\eta/3$ (the point line in Figure 1d), ζ_R and γ_d (the thick line in Figure 1d) without heat input by sound waves' dissipation, ($\epsilon k = 0$ [cm^{-1}]), and for a plasma at equilibrium, i.e. $X(\rho, T, \xi) = 0$ and $L(\rho, T, \xi) = 0$.

For the above-mentioned particular values of density, metallicity, and mean photon energy, the thermal equilibrium may only exist for $T < 2.254 \times 10^4$ [K], and the nature of the equilibrium for this particular case has been analyzed, see some approximation by Ibañez (2009). As it can be seen in Figure 1d, the dynamic viscosity term (point line), and the thermal conduction term (dash line) are very small respect to the bulk viscosity ζ_R (thick line) which determines the total dissipation γ_d that is undistinguished from ζ_R at the scale shown in Figure 1d.

When one considers sound waves' dissipation the above results can drastically change depending on the value of the sound wavelength, more exactly on the value of ϵk. In fact, for the above-written values of parameters of the gas, sound waves with $\epsilon k \lesssim 10^{-18}$ [cm^{-1}] only produce small changes on variables the of interest, in particular at high temperatures $T > 10^3$ [K] as it is apparent in Figure 2 that

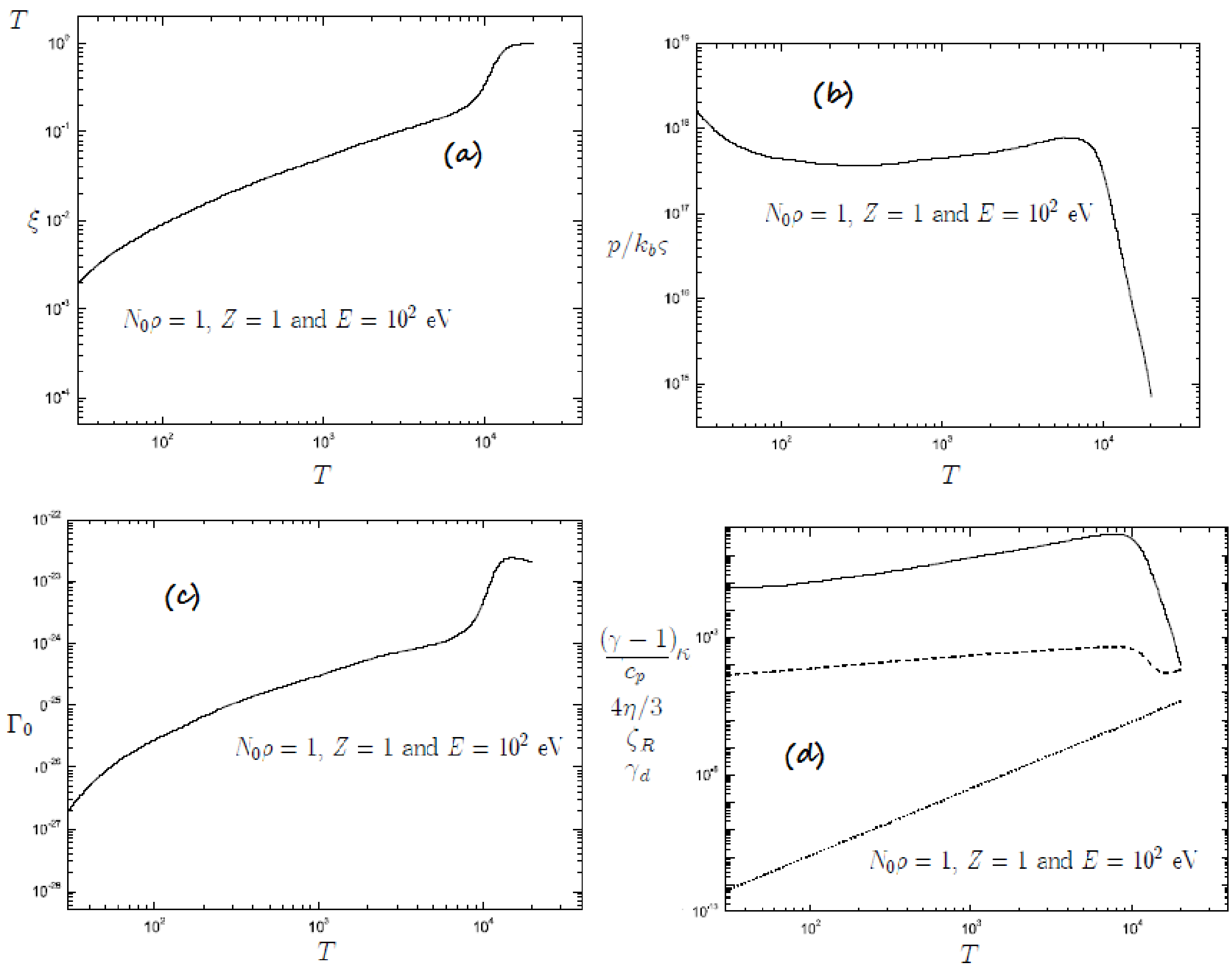

Fig. 1. The ionization ξ (a), the pressure at equilibrium (b), the heating input Γ_0 (c), and four dissipative terms (d) for a gas with $N_0\rho = 1$, $Z = 1$, and $E = 10^2$ [eV] when $\epsilon k = 0$ [cm^{-1}] for a plasma at equilibrium, i.e., $X(\rho, T, \xi) = 0$ and $L(\rho, T, \xi) = 0$. The dissipation γ_d is undistinguished from the values of the real part of the bulk viscosity ζ_R at the scale shown in Figure 1d.

shows the plots as those shown in Figure 1 but for ϵk = 10^{-18} [cm^{-1}]. This is due to the fact that the heat input Γ_0 by dissipation of sounds shown by the dash line in Figure 2c, at these wavelengths, is one order of magnitude smaller than the magnitude of the radiative heating Γ_ω shown by the thick line in Figure 2c.

The results shown in Figure 3 are similar to the ones shown in Figure 2 but for $\epsilon k = 10^{-17}$ [cm^{-1}]. At this shorter value of the acoustic wavelength, qualitative changes appear. In addition to a modified equilibrium ionization given by the upper branch in Figure 3a there is some comparison with the case when wave heat input is neglected. Also, an additional ionization branch appears (lower branch in Fig. 3a) but this ionization corresponds to a higher pressure (upper branch in Figure 3b) which occurs in the following narrow interval of temperatures: 7×10^3 [K] < T < 10^4 [K].

These two values of ionization and pressure produce two values of the heating rates $\Gamma_0(\rho, T, \xi)$ and $\Gamma_\omega(\rho, T, \xi)$ shown in Figure 3c as well as two values in the dissipation coefficients κ, η, and ζ_R shown in Figure 3d in the above-written interval of temperatures. From the physical point of view, the above-obtained results imply that when the heat input by acoustic waves in a recombining gas becomes important, the gas can be at equilibrium in two different phases: one at high ionization and low pressure, and the another one at low ionization and high pressure for the same value of the density.

The effects shown in Figure 3 shift towards higher values of the temperature when the value of ϵk increases as can be seen in Figure 4 that shows the results shown in Figure 3 but for the value of $\epsilon k = 10^{-16}$ [cm^{-1}]. For this particular value of ϵk, the ionization (Fig. 4a), pressure (Fig. 4b), both heating rates (Fig. 4c), as well as the dissipative terms (Fig. 4d) show some bifurcation point at certain temperature $T = 1.14 \times 10^4$ [K].

A low ionization branch and high-pressure branch appear in the range of temperatures 1.14×10^4 [K] < T < $1.268 \times$

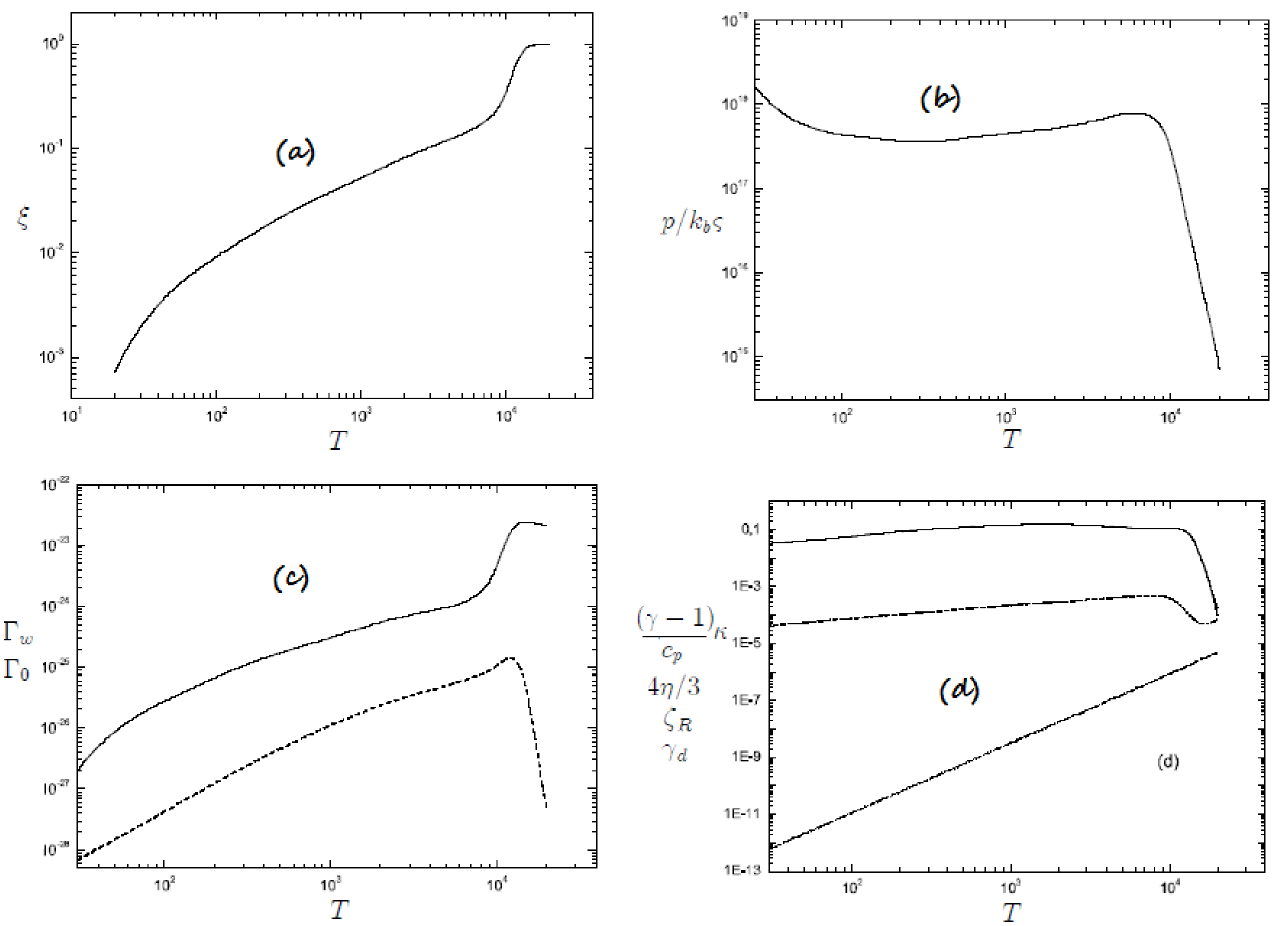

Fig. 2. The ionization ξ (a), the pressure (b), the radiative heating Γ_ω and the heating input Γ_0 (c), and four dissipative terms (d) for a gas with $N_0\rho$ = 1, Z = 1, and E = 10^2 [eV] when ϵk = 10^{-18} [cm^{-1}]. The dissipation γ_d again is undistinguished from the values of the real part of the bulk viscosity ζ_R at the scale shown in Figure 2d.

10^4 [K]. For these equilibrium states the heating $\Gamma_0(\rho, T, \xi) < \Gamma_\omega(\rho, T, \xi)$. In Figures 4a, 4b, and 4c, the radiation heating is higher than the wave heating, contrary to the high ionization (and pressure) branch. For this case, the thermal conduction dissipation terms dominate over the viscosity terms. The low ionization branch corresponds to the total dissipative branch γ_d higher than the corresponding higher ionization shown in Figure 4d.

When the value of ϵk increases the interval of temperature where this equilibrium may exist decreases in such a way that the equilibrium cannot exist for $\epsilon k = 5 \times 10^{-16}$ [cm^{-1}]. Figure 5 shows the plot of the variables under discussion that have been plotted for $\epsilon k = 5 \times 10^{-16}$ [cm^{-1}]. For this particular value, the ionization and pressure become one-valued functions of temperature (as in the case $\epsilon k = 10^{-18}$ [cm^{-1}]) but the gas pressure shows a minimum value at $T = 1.849 \times 10^4$ [K] shown in Figure 5b. At this limiting value the acoustic dissipation $\Gamma_\omega(\rho, T, \xi)$ becomes higher than the radiation heating $\Gamma_0(\rho, T, \xi)$ as it is apparent from the results shown in Figure 5c. Note that in this case, the dissipation is mainly due to the thermal conduction, as for the previous ϵk-value, please, see in Figure 5d.

The values of the damping length scale l_d shown in Figure 6 were obtained for the following four values of the magnitude ϵk [cm^{-1}], i.e., $\epsilon k = 10^{-18}$ [cm^{-1}] (a), $\epsilon k = 10^{-17}$ [cm^{-1}] (b), $\epsilon k = 10^{-16}$ [cm^{-1}] (c), and $\epsilon k = 5 \times 10^{-16}$ [cm^{-1}] (d). The temperature dependences of the damping length scale l_d (equation (24)) are shown in Figures 6a, 6b, 6c, and 6d, respectively.

Figure 7 shows the plots for the corresponding ratios $\lambda/(2\pi l_d)$. For large wavelengths when the sound wave heat input is much smaller than the radiative heating shown in Figure 6a, some regions of the order of parsecs up to kiloparsecs can exist in thermochemical equilibrium in a wide range of temperatures. When the wave number decreases this temperature range decreases (as well as the dimensions of the regions where such an equilibrium may exist) and the two equilibria become apparent.

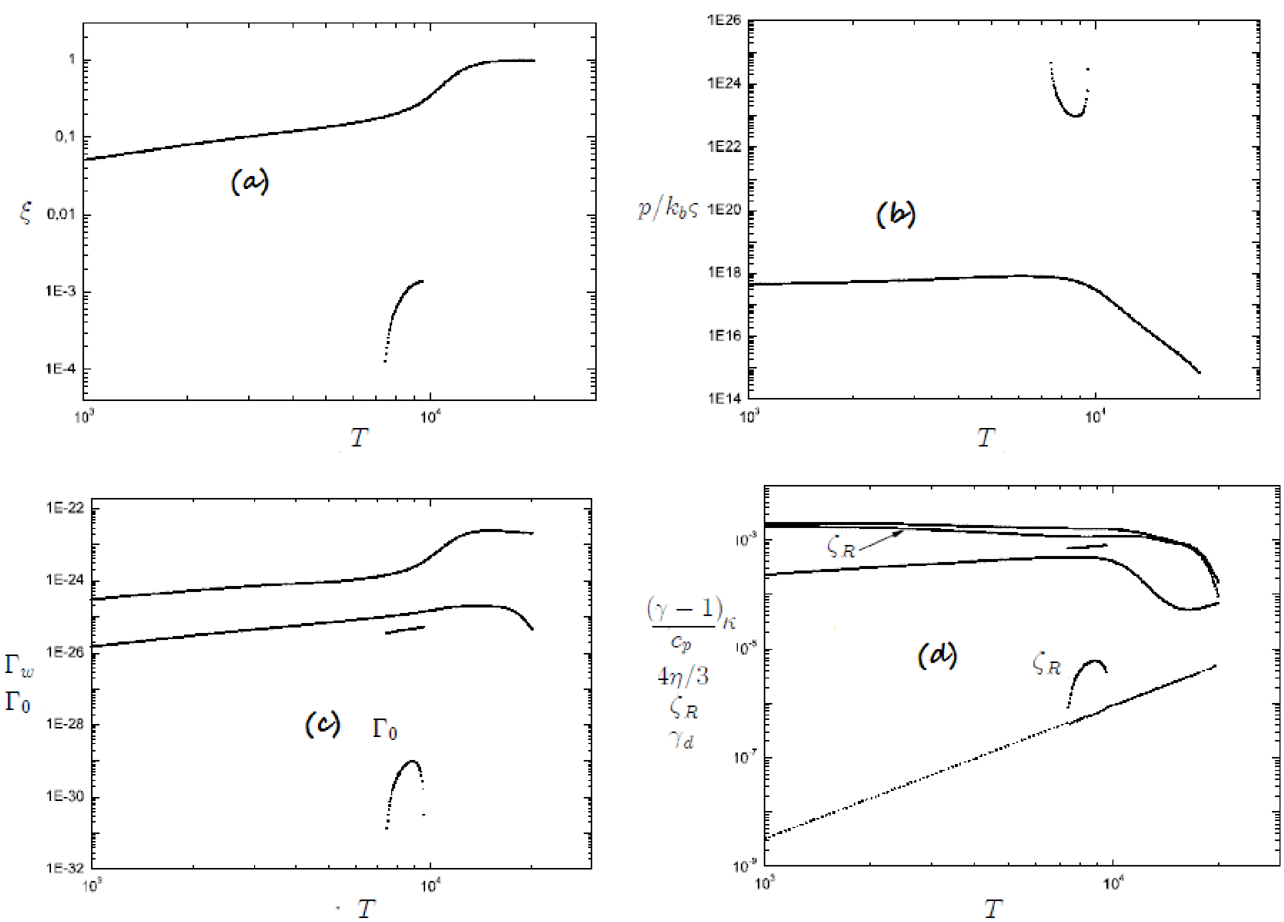

Fig. 3. The ionization ξ (a), the pressure (b), the radiative heating Γ_ω and the heating input Γ_0 (c), the four dissipative terms (d) for a gas when $\epsilon k = 10^{-17}$ [cm^{-1}].

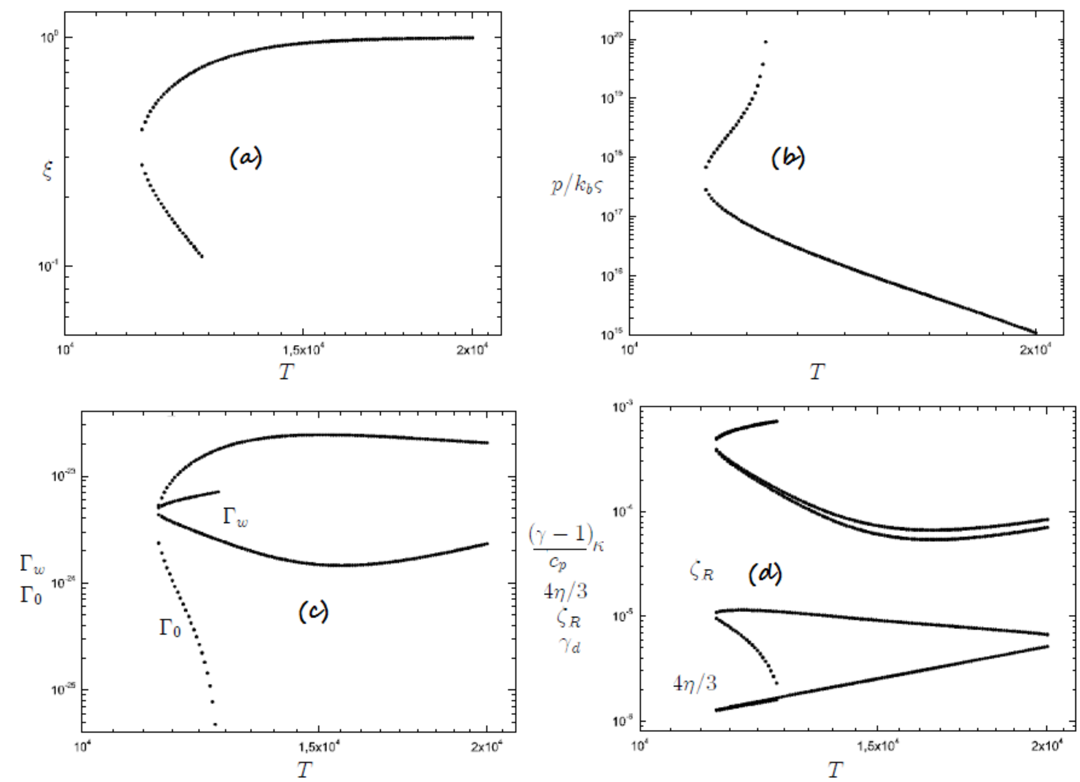

Fig. 4. The ionization ξ (a), the pressure (b), the radiative heating Γ_ω and the heating input Γ_0 (c), the four dissipative terms (d) for a gas when $\epsilon k = 10^{-16}$ [cm^{-1}]. A bifurcation point is seen in all four Figures for $T = 1.14 \times 10^4$ [K].

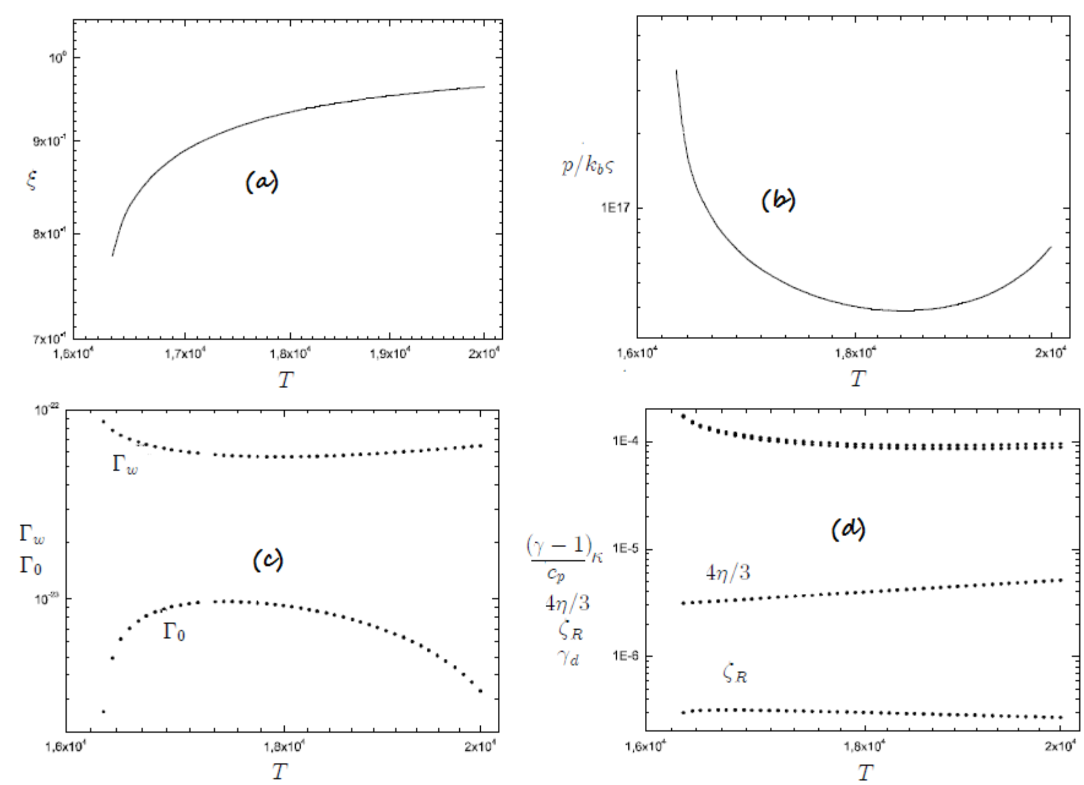

Fig. 5. The ionization ξ (a), the pressure (b), the radiative heating Γ_ω and the heating input Γ_0 (c), the four dissipative terms (d) for a gas when $\epsilon k = 5 \times 10^{-16}$ [cm^{-1}]. In Figure 5d, the higher lines correspond to the thermal conduction.

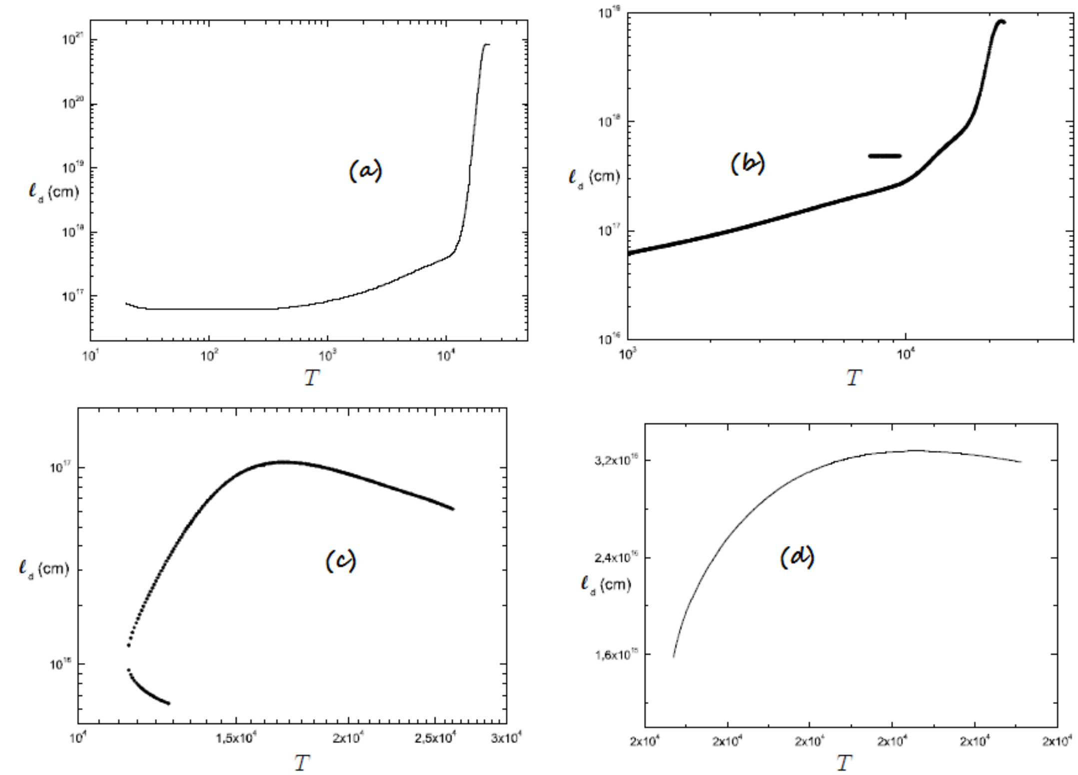

Fig. 6. The damping length scale l_d as the function of the temperature T, for the following four values of the magnitude ϵk [cm^{-1}], i.e., $\epsilon k = 10^{-18}$ (a), $\epsilon k = 10^{-17}$ (b), $\epsilon k = 10^{-16}$ (c), and $\epsilon k = 5 \times 10^{-16}$ (d).

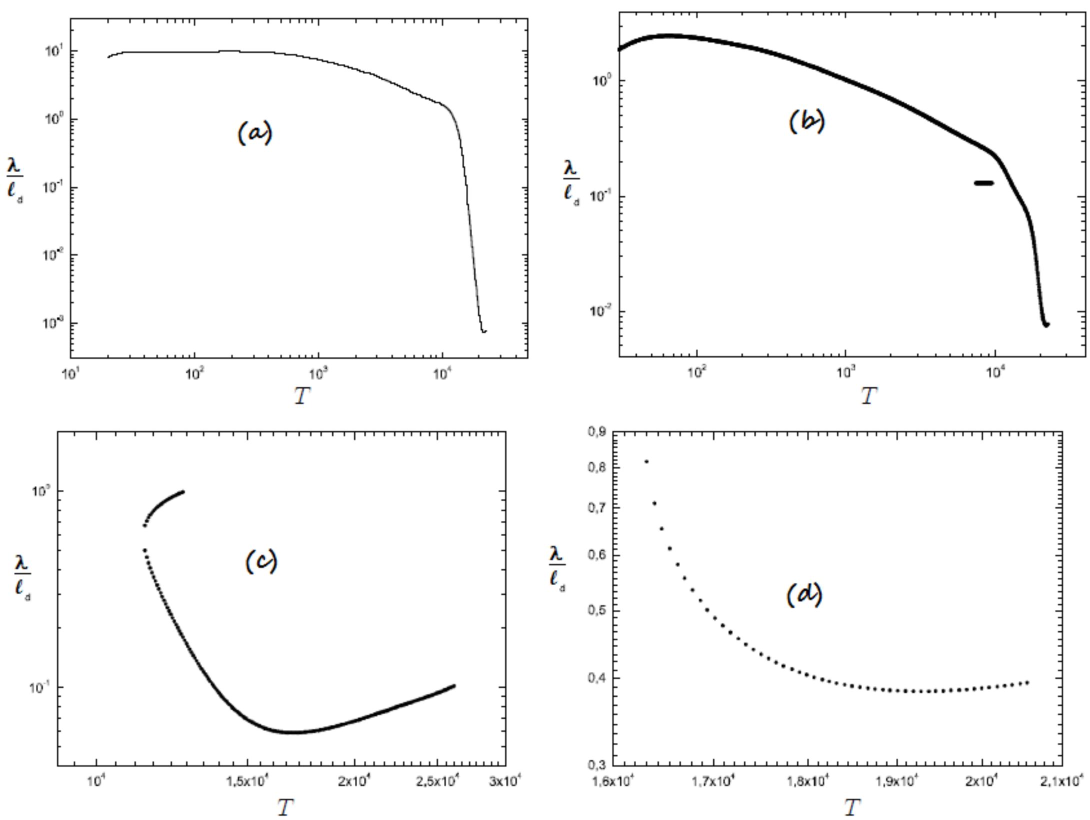

Fig. 7. The ratios λ/l_d for $\epsilon k = 10^{-18}$ (a), $\epsilon k = 10^{-17}$ (b), $\epsilon k = 10^{-16}$ (c), and $\epsilon k = 5 \times 10^{-16}$ (d).

Regions of high pressure are larger than those at lower pressure for $\epsilon k = 10^{-17}$ [cm^{-1}] (Fig. 6b) but the opposite occurs for $\epsilon k = 10^{-16}$ [cm^{-1}] (Fig. 6c). Close to the limiting value where the equilibrium may exist, only small regions of the order of 10^{-3} [pc] can exist in a narrow interval of temperatures. As it is apparent from the results shown in Figure 7, the present approximation holds, $\lambda/l_d < 1$ in the range of temperatures where the recombination-ionization becomes important.

High velocity clouds at high galactic latitude (HvCs)

Observations suggest that for the hydrogen gas at large galactic latitude and in the halos of other galaxies, the metallicity ranges are between $0.03 \lesssim Z \lesssim 0.3$ and the mean photon energy ones are between 0.3 [keV] $\lesssim E \lesssim$ 2 [keV] (Collins *et al*., 2004; Maller *et al*., 2004; Miller *et al*., 2015; Wakker *et al*., 2001; Wakker, 2004; Tripp *et al*., 2003; Olano, 2008; Lockman *et al*., 2008, Ibañez *et al*., 2011) and (Conde and Sandra, 2010. Calentamiento de plasmas por ondas hidrodinamicas. M. Sc. Thesis, University of Los Andes, Center for Fundamental Physics). Thus, equilibrium states may exist for the values of ϵk smaller than the threshold value of $(\epsilon k)_{thr}$ that depends on the exact values of Z and E (Fig. 8).

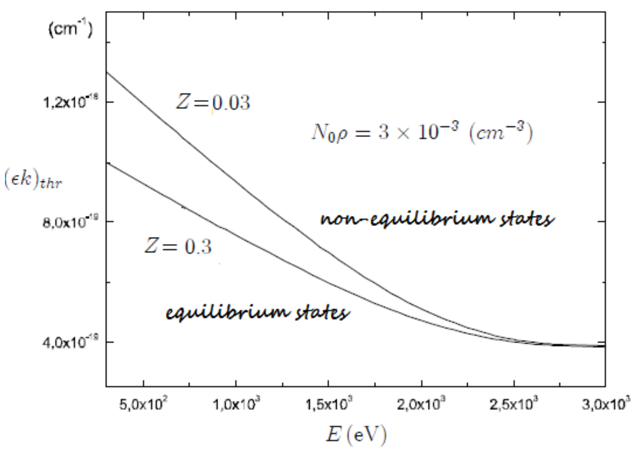

Fig. 8. The threshold value of $(\epsilon k)_{thr}$ as a function of the mean radiation field E for the metallicity values of Z = 0.03, 0.3 in HvCs with $N_0\rho = 3 \times 10^{-3}$ [cm^{-3}].

For $\epsilon k > (\epsilon k)_{thr}$, equilibrium states cannot exit. Therefore, the hydrogen gas into High velocity Clouds (HvCs) can be in thermochemical equilibrium as far as the HvCs have dimensions $l > l_{thr}$, where $l_{thr} = 3.24 \times 10^{-19} \times (\epsilon k)_{thr}$ [pc] is the above-mentioned range of values of metallicity and mean photon energy. Strictly speaking this threshold value has to be taken as a limiting value, for real situations one should expect thermochemical equilibrium for gas clouds with dimensions $l >> l_{thr}$. This conclusion can also be applied to the intracluster clouds. However, this important case will be analyzed elsewhere.

Additionally, depending on the particular values of Z and E, temperature gaps appear where the equilibrium cannot exist. Say for instance, for Z = 0.3 and E = 0.3 [keV] the gap is situated between the following temperatures: 2.405×10^4 [K] < T < 8.149×10^5 [K], and for Z = 0.03 it is between the ones: 2.297×10^4 [K] < T < 2.255 × 105 [K]. For this particular value of Z, only one equilibrium exists just at $T = 2.255 \times 10^5$ [K]. From the above results one may conclude that HvCs at thermal equilibrium may only exist in well-defined ranges of temperature (depending on the E, Z, and $N_0\rho$) and with dimensions $l >> l_{thr}$.

CONCLUSION

Generally, the dissipation of acoustic waves by the thermal conduction, dynamic and bulk viscosities can introduce drastic quantitative and qualitative changes in the equilibrium of the interstellar gas. Depending on the particular values of the metallicity, the mean radiation energy input, and the particle density, we can state that:

- The bulk viscosity can be the most important dissipation mechanism.
- Several equilibrium branches may appear in certain temperature ranges that depend on the particular values of ϵk.
- For large enough values of ϵk, the range of temperature where equilibrium may occurs becomes very narrow.
- For the ϵk-values larger than the threshold value of $(\epsilon k)_{thr}$, the equilibrium does not exist that imposes a limiting value to the scale length of equilibrium structures, in particular, to the dimensions of HvCs.

Strictly speaking, in any plasma, in particular, in a very complex interstellar medium (ISM) a spectral distribution (discrete one as well as continuous one depending on the origin of the waves) for sound waves should be present. However, it is an important issue, namely determining for which particular values of the ϵk of such spectrum the equilibrium ($X(\rho, T, \xi) = 0$ and $L(\rho, T, \xi) = 0$) is modified by the waves' dissipation.

We emphasize that this work is essentially restricted to a linear approximation, in which the wave dissipation is a second order phenomenon (Landau and Lifshitz, 1987). For this, the results depend on the ϵk that allows one to get a "first" information on the dependence on the wave amplitude.

On the other hand, turbulence is by itself a nonlinear phenomenon, and this problem out the scope of the paper. Certainly, turbulence should be present in the different "regions" of the ISM, which is a typical (very complex) plasma where many relaxation time scales go into play. Models beyond the qualitative ones (Kolmogorov-like) (Braun *et al*., 2012) do not exist, and the understanding of turbulent plasmas under realistic physical conditions is an open question.

Authorship contribution statement

P. Ibañez: Conceptualization, methodology, software, investigation, validation, writing (original draft), and supervision.
S. Conde: Methodology, validation, investigation, writing (original draft).
P. Contreras: Investigation, data curation, visualization, writing (review and editing).

Declaration of competing interest

The authors declare that there are no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

REFERENCES

Bird, GA. 1964. A gas-dynamic model of the outer solar atmosphere. The Astrophysical Journal. 139(2):684-689. DOI: http://dx.doi.org/10.1086/147794.

Braginskii, SI. 1965. Transport processes in a plasma. Reviews of Plasma Physics. 1:205-217.

Braun, H. and Schmidt, W. 2012. A semi-analytic model of the turbulent multi-phase interstellar medium. Monthly Notices of the Royal Astronomical Society. 421(3):1838-1860. DOI: https://doi.org/10.1111/j.1365-2966.2011.19889.x.

Corbelli, E. and Ferrara, A. 1995. Instabilities in photoionized interstellar gas. The Astrophysical Journal. 447(2):708-720.

Collins, JA., Shull, JM. and Giroux, ML. 2004. Highly ionized high-velocity clouds toward PKS 2155-304 and Markarian 509. The Astrophysical Journal. 605(1):216. DOI: https://doi.org/10.1086/382269.

Dalgarno, A. and McCray, RA. 1972. Heating and ionization of HI regions. Ann. Rev. Astron. Astroph. 10:375-426.

Fabian, AC., Sanders, JS., Allen, SW., Crawford, CS., Iwasawa, K., Johnstone, RM., Schmidt, RW. and Taylor, GB. 2003. A deep Chandra observation of the Perseus cluster: shocks and ripples. Monthly Notices of the Royal Astronomical Society. 344(3):L43-L47. DOI: https://doi.org/10.1046/j.1365-8711.2003.06902.x.

Fabian, AC., Reynolds, CS., Taylor, GB. and Dunn, RJH. 2005. Monthly Notices of the Royal Astronomical Society. 363(3):891-896. DOI: https://doi.org/10.1111/j.1365-2966.2005.09484.x.

Fajardo M., JJ., Ibañez S., MH. and Contreras E., PL. 2021. An eigenvalue analysis of damping in optical thin plasmas. Rev. Mex. Fis. 67(6):061502 (13 pages). DOI: https://doi.org/10.31349/RevMexFis.67.061502.

Ferland, GJ., Fabian, AC., Hatch, NA., Johnstone, RM., Porter, RL., van Hoof, PAM. and Williams, RJR. 2009. Collisional heating as the origin of filament emission in galaxy clusters. Monthly Notices of the Royal Astronomical Society. 392(4):1475-1502. DOI: https://doi.org/10.1111/j.1365-2966.2008.14153.x.

Ibáñez S., MH. 2009. The stability of optically thin reacting plasmas: Effects of the bulk viscosity. The Astrophysical Journal. 695(1):479-487. DOI: https://doi.org/10.1088/0004-637X/695/1/479.

Ibañez S., MH. and Escalona, OB. 1993. Propagation of hydrodynamic waves in optically thin plasmas. The Astrophysical Journal. 415:335-341. DOI: https://doi.org/10.1086/173167.

Ibañez S., MH., Shchekinov, Y. and Bessega, MC. 2005. On the stability of subsonic thermal fronts. Physics of Plasmas. 12(8):082108. DOI: https://doi.org/10.1063/1.1993287.

Ibañez S., MH. and Conde C., SM. 2011. Heating of a recombining plasma by bulk viscosity dissipation of sound waves. Rev. Mex. Astron. Astrofis. Ser. Conf. 40:188.

Ibañez S., MH. and Contreras E., PL. 2019. Damping of sound waves by bulk viscosity in reacting gases. Can. J. Pure Appl. Sci. 13(3):4897-4906.

Landau, LD. and Lifshitz, EM. 1987. Fluid Mechanics. Pergamon Press, London, UK.

Launay, JM. and Roueff, E. 1977. Fine structure excitation of carbon and oxygen by atomic hydrogen impact. Astronomy and Astrophysics. 56(1-2):289-292.

Lockman, FJ., Benjamin, RA., Heroux, AJ. and Langston, GI. 2008. The Smith cloud: A high-velocity cloud colliding with the Milky Way. The Astrophysical Journal. 679(1):21L. DOI: https://doi.org/10.1086/588838.

Maller, AH. and Bullock, JS. 2004. Multiphase galaxy formation: high-velocity clouds and the missing baryon problem. Monthly Notices of the Royal Astronomical Society. 355:694-712.

Mihalas, D. and Mihalas, BW. 1984. Foundations of Radiation Hydrodynamics. Oxford University Press, Oxford, UK.

Miller, MJ. and Bregman, JN. 2015. Constraining the Milky Way's hot gas halo with 0 vii emission lines. The Astrophysical Journal. 800(1):14. DOI: https://doi.org/10.1088/0004-637X/800/1/14.

Narain, U. and Ulmschneider, P. 1990. Chromospheric and coronal heating mechanisms. Space Science Reviews. 54(3-4):377-445. DOI: http://dx.doi.org/10.1007/BF00177801.

Olano, CA. 2008. Distribution of the high-velocity clouds in the Galactic halo. Astronomy and Astrophysics. 485(2):457-473. DOI: https://doi.org/10.1051/0004-6361:20077556.

Ruszkowski, M., Brüggen M. and Begelman, MC. 2004. Cluster heating by viscous dissipation of sound waves. The Astrophysical Journal. 611(1):158-163. DOI: https://doi.org/10.1086/422158.

Seaton, MJ. 1959. Radiative recombination of hydrogenic ions. Monthly Notices of the Royal Astronomical Society. 119(2):81-89. DOI: https://doi.org/10.1093/mnras/119.2.81.

Stein, RF. and Schwartz, RA. 1972. Waves in the solar atmosphere. II. Large-amplitude acoustic pulse propagation. The Astrophysical Journal. 177:807-828. DOI: http://dx.doi.org/10.1086/151757.

Stein, RF. and Leibacher, J. 1974. Waves in the solar atmosphere. Annual Review of Astronomy and Astrophysics. 12:407-435. DOI: http://dx.doi.org/10.1146/annurev.aa.12.090174.002203.

Shull, JM. and van Steenberg, ME. 1985. X-ray secondary heating and ionization in quasar emission-line clouds. The Astrophysical Journal. 298(1, Part 1):286-274. DOI: http://dx.doi.org/10.1086/163605.

Spitzer, L. 1978. Physical Processes in the Interstellar Medium. Interscience, John Wiley and Sons Inc., New York, USA. Pp.333. DOI: http://dx.doi.org/10.1002/9783527617722.

Spitzer, L. 1982. Acoustic waves in supernova remnants. The Astrophysical Journal. 262(1, Part 1):315-321. DOI: http://dx.doi.org/10.1086/160423.

Spitzer, L. 1990. Theories of the hot interstellar gas. Annual Review of Astronomy and Astrophysics. 28:71-101. DOI: https://doi.org/10.1146/annurev.aa.28.090190.000443.

Spitzer, L. 1962. Physics of Fully Ionized Gases. Interscience, John Wiley and Sons Inc., New York, USA. pp.61.

Stix, TH. 1992. Waves in Plasmas. Springer, Berlin, Germany.

Tripp, TM., Wakker, BP., Jenkins, EB., Bowers, CW., Danks, AC., Green, RF., Heap, SR., Joseph, CL., Kaiser, ME., Linsky, JL. and Woodgate, BE. 2003. Complex C: A low-metallicity, high-velocity cloud plunging into the Milky Way. The Astronomical Journal. 125:3122-3144.

Wakker, BP. 2001. Distances and metallicities of high- and intermediate- velocity clouds. Astrophysical J. Supplement. 136:463-535.

Wakker, BP. 2004. High velocity clouds and the local group. IAU Symposium Series. 217:2.

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