Mon. Not. R. Astron. Soc. **312**, 747–752 (2000)

Is the radiation temperature–redshift relation of the standard cosmology in accordance with the data?

J. A. S. Lima,[1★] A. I. Silva[2★] and S. M. Viegas[2★]

[1]*Universidade Federal do Rio Grande do Norte, Departamento de Física, CP 1641, 59072-970 Natal, RN, Brazil*

[2]*Instituto Astronômico e Geofísico, USP, Av. Miguel Stefano 4200, 04301-904 São Paulo, SP, Brazil*

Accepted 1999 October 11. Received 1999 August 4

ABSTRACT

The radiation temperature–redshift relation for Friedmann–Robertson–Walker geometries is rediscussed in connection with recent observational data based on the fine-structure splitting of atomic and singly ionized carbon lines in quasar absorption-line systems. Indirect measurement of $T(z)$ is one of the most powerful cosmological tests available because it may exclude even the presence of a cosmological constant. Unlike recent claims, we argue that the temperature at high z may be smaller than the standard prediction, thereby opening a window to alternative (big bang) models. By including new ingredients like a phenomenological decaying vacuum energy density and gravitational 'adiabatic' photon creation as well as late inflationary models driven by a scalar field, a new temperature law is deduced and its predictions are compared with the standard result.

Key words: galaxies: distances and redshifts – cosmic microwave background – cosmology: theory.

1 INTRODUCTION

The existence of the cosmic microwave background radiation (CMBR) is regarded as the best evidence for a primeval expanding state of the Universe. The present-day temperature of the CMBR is accurately measured by the *COBE* FIRAS experiment to be $T_{COBE} = 2.728 \pm 0.002\,\mathrm{K}$ (1σ confidence level; Fixsen et al. 1996). The *COBE* experiment also provides data with remarkable sensitivity on the isotropy (Smoot et al. 1992) and blackbody nature (Mather et al. 1994) of the spectrum. However, this result by itself does not constitute a proof of the cosmological nature of the CMBR. If the universe really started from a homogeneous and isotropic collapsed state, the simplest prediction of such a model is that the temperature of this blackbody radiation has always been spatially uniform and higher in the past. Naturally, this cannot be inferred from *COBE* experiments, and requires measurements of the CMBR temperature at moderate and high redshifts.

How does the CMBR temperature increase when we look back in time? The answer to this question is immediate only for Friedmann–Robertson–Walker (FRW) models where the radiation fluid expands adiabatically, especially after decoupling. Actually, if the average number of photons is conserved, the scaling law of the entropy, $S \sim T^3R^3$, implies that the temperature grows proportionally with the redshift. However, in a more general framework, the function $T(z)$ depends on the ingredients present in the FRW geometry or, equivalently, on the possible mechanisms acting upon the radiative component. If photon creation takes place, for instance, because of a continuous decaying vacuum energy density or some alternative mechanism of quantum gravitational origin, the standard linear law is necessarily modified. In this case, measurements of the temperature at high redshifts constrain the free parameters of these models. As we shall see, the $T(z)$ relation is a potential candidate for a new source of difficulties for the standard model. If confirmed by the observations, such difficulties cannot be resolved either by the presence of a pure cosmological constant or any non-luminous (separately conserved)'quintessence'.

2 TEMPERATURE–REDSHIFT RELATION

In order to emphasize the main assumptions involved in this problem, we first discuss how the temperature law $T(z)$ may be derived in a framework more general than the familiar photon conserved picture.

We shall restrict our analysis to cosmologies described by the FRW line element ($c = 1$):

$$\mathrm{d}s^2 = \mathrm{d}t^2 - R^2(t)\left(\frac{\mathrm{d}r^2}{1-kr^2} + r^2\,\mathrm{d}\theta^2 + r^2\,sin^2\theta\,\mathrm{d}\phi^2\right), \quad (1)$$

where R is the scalefactor and $k = 0, \pm 1$ is the curvature parameter.

For the sake of generality, we assume that the material medium of the universe is a mixture of three different components: a γ fluid, which includes radiation and matter as particular cases, a

★ E-mail: jas@dfte.ufrn.br(JASL); alexsilv@iagusp.usp.br (AIS); viegas@agn1.iagusp.usp.br (SMV)

'quintessence' x component and a 'bare' cosmological constant. In this case, the Einstein field equations are

$$8\pi G(\rho + \rho_x) + \Lambda_0 = 3\frac{\dot{R}^2}{R^2} + 3\frac{k}{R^2}, \tag{2}$$

$$8\pi G(p + p_x) - \Lambda_0 = -2\frac{\ddot{R}}{R} - \frac{\dot{R}^2}{R^2} - \frac{k}{R^2}, \tag{3}$$

where a dot means time derivative, Λ_0 is the cosmological constant, and ρ, ρ_x, p, p_x are the densities and pressures of the γ fluid and x component, respectively. The energy conservation law ($u_\alpha T^{\alpha\beta}_{;\beta} = 0$) and the balance equation for the particle number density can be written as

$$\dot{\rho} + 3(\rho + p)H = C_x, \tag{4}$$

$$\dot{n} + 3nH = \psi, \tag{5}$$

where $H = \dot{R}/R$ is the Hubble parameter, n is the particle number density and ψ is the particle source term, which is directly related to the presence of the x component. For convenience we have also introduced the notation

$$C_x = -\dot{\rho}_x - 3(\rho_x + p_x)H. \tag{6}$$

This quantity depends only on the x component, and works essentially like a source term for the γ-fluid energy. It is worth notice that the C_x term may describe several situations of current cosmological interest. The most important cases are the following.

(i) A time-dependent cosmological constant $\Lambda(t)$ (Ozer & Taha 1986, 1987; Freese et al. 1987; Chen & Wu 1990; Carvalho, Lima & Waga 1992; Lima & Maia 1994; Lima 1996; Lima & Trodden 1996; Silveira & Waga 1997; Overduin & Cooperstock 1998). In this case, $\rho_x = -p_x = \Lambda(t)/(8\pi G)$ and using (6) we find $C_x = -\dot{\Lambda}/(8\pi G)$.

(ii) Thermogravitational quantum creation (Prigogine et al. 1989; Calvão, Lima & Waga 1992; Lima & Germano 1992; Zimdhal & Pavón 1994; Lima, Germano & Abramo 1996; Lima & Alcaniz 1999). Now $\rho_x = 0$, p_x is a creation pressure, and from (6) we have $C_x = -3p_xH$. Although quite different on thermodynamical grounds, this approach resembles the old bulk viscosity description for matter creation. For this kind of models there is no second fluid component, because only an additional negative pressure takes place in the medium.

(iii) Scalar field cosmologies. The x component for this type of model is defined in terms of a Langrangian density, $L = \partial^\mu\phi\partial_\mu\phi - V(\phi)$, for which $\rho_x = \dot{\phi}^2/2 + V(\phi)$ and $p_x = \dot{\phi}^2/2 - V(\phi)$. As one may check, $C_x = -\dot{\phi}\ddot{\phi} - 3H\dot{\phi}^2 - V'(\phi)\dot{\phi}$, where the dash denotes a derivative with respect to the field ϕ. Usually, the field is assumed to be very weakly coupled to produce ordinary matter and radiation (Ratra & Peebles 1988; Frieman et al. 1995; Caldwell, Dave & Steinhardt 1997; Frieman & Waga 1998). In principle, we consider here a more general framework. Even at late stages, when the field rolls down the potential, it is continuous and slowly transfers energy and particles to the γ-fluid component.

Naturally, if $C_x = 0$, both the x fluid and the normal component are separately conserved. In this case, we must have $\psi = 0$ and the standard photon conserved picture is maintained.

The macroscopic quantities of the γ fluid are related by the Gibbs law:

$$nT\,\mathrm{d}\sigma = \mathrm{d}\rho - \frac{\rho + p}{n}\mathrm{d}n, \tag{7}$$

where σ is the specific entropy and T is the temperature. Now, because $\mathrm{d}\sigma$ is an exact differential, the above expression yields the well-known thermodynamic identity

$$T\left(\frac{\partial p}{\partial T}\right)_n = \rho + p - n\left(\frac{\partial \rho}{\partial n}\right)_T. \tag{8}$$

The thermodynamic behaviour of the γ component is completely defined by equations (4), (5), (7) and (8). By adopting the pair (T, n) as thermodynamic independent variables, the variation rates of T and σ are given by (Lima 1996)

$$\frac{\dot{T}}{T} = \left(\frac{\partial p}{\partial \rho}\right)_n \frac{\dot{n}}{n} - \frac{\psi}{nT\left(\frac{\partial \rho}{\partial T}\right)_n}\left[\rho + p - \frac{nC_x}{\psi}\right], \tag{9}$$

$$\dot{\sigma} = \frac{\psi}{nT\left(\frac{\partial \rho}{\partial T}\right)_n}\left[\rho + p - \frac{nC_x}{\psi}\right]. \tag{10}$$

Although a general analysis might be carried out, in what follows we assume that the γ component is a radiation fluid ($p = \rho/3$). We first remark that the equilibrium blackbody relation, $n \sim T^3$, is possible only if the second term on the right-hand side of (9) vanishes identically, that is

$$C_x = \frac{4\rho}{3n}\psi. \tag{11}$$

On the other hand, from (10) we see that the specific entropy is constant. For the cases (i) and (ii) such a condition has been termed 'adiabatic' photon creation, because the radiation entropy and the total number of photons increases but the ratio $\sigma = S/N$ remains constant (Calvão et al. 1992; Lima 1996; Lima et al. 1996). As a result, the radiation entropy increases proportionally to the comoving number of photons N. Naturally, even for adiabatic creation, there are out-of-equilibrium contributions encoded in the temperature law. Actually, inserting (5) and (11) into (9) one finds the generalized relation

$$\frac{\dot{T}}{T} = -\frac{\dot{R}}{R} + \frac{\psi}{3n}. \tag{12}$$

As expected, for $\psi = 0$ the standard FRW law is recovered:

$$TR = T_0R_0 = \text{constant}, \tag{13}$$

or, equivalently,

$$T = T_0(1 + z). \tag{14}$$

In this form, the above temperature–redshift law holds for the FRW class (with and with no cosmological constant), as well as for any cosmology with an exotic decoupled x component. Particular examples are the scalar field cosmologies proposed by Ratra & Peebles (1988). Therefore, a new thermodynamic cosmological scenario comes out only if ψ (or equivalently C_x) is different from zero. In this case, modelling the quantity ψ, C_x is automatically obtained from (11), and a new thermodynamic scenario may be established by rewriting (5) as

$$\frac{\dot{n}}{3nH} + 1 = \beta, \tag{15}$$

where the convenient parameter $\beta(t) = \psi/(3nH)$ has been introduced. This dimensionless parameter is the ratio between two relevant rates, and is constrained by $0 \leqslant \beta(t) \leqslant 1$. The lower limit yields the photon-conserved picture (adiabatic expansion) and the upper limit also arises in a natural way. It defines an extreme theoretical situation where the matter creation phenomenon is so huge that it compensates for the dilution resulting from expansion. This parameter is a function of the cosmological time because the 'decaying rate' $\psi/3n$, in principle, is not proportional to the Hubble parameter. However, in order to present a class of models slightly more general than the standard photon-conserving scenario, it will be assumed that such a ratio is constant. This hypothesis seems to be physically reasonable, at least to the late stages of the evolution, when the radiation fluid has decoupled. A more general discussion of this subject will be presented soon (Lima, in preparation). For constant β, a straightforward integration of (12) yields

$$T = T_0(1+z)^{1-\beta}. \tag{16}$$

This result means that for a given redshift z, the temperature of the expanding universe is slightly lower than in the standard photon-conserved scenario. Now, motivated by this theoretical possibility, let us discuss whether the current observations give some support to this temperature law.

3 PRESENT OBSERVATIONAL STATUS

As is widely known, the best experimental method available to measure $T(z)$ is provided by absorption lines from molecules, atoms and ions. The lines from cyanogen (CN) are of special interest because of their sensitivity, as are those of carbon (C^0 and C^+), which can be observed locally and also in the spectra of distant quasi-stellar objects (QSOs). We recall that observation of CN absorption lines from diffuse interstellar clouds in the line of sight of bright nearby stars yielded a CMBR temperature of $T_0 = 2.729^{+0.023}_{-0.031}$ K (Roth, Meyer & Hawkings 1993; Roth & Meyer 1995), in perfect agreement with the *COBE* FIRAS result. However, for other values of z, the temperature of the CMBR has been obtained from observations of QSO absorbers, as proposed by Bachall & Wolfe (1968).

The general method is to use the column density ratio of the absorption line originating from the fundamental and first excited states of C^0 and C^+, thus obtaining an excitation temperature defined as

$$T_{\rm exc} = \frac{\Delta E}{k \ln\left(\dfrac{g^*/g}{N^*/N}\right)}, \tag{17}$$

where ΔE is the energy difference between the fine-structure levels, k is Boltzmann's constant, g are the statistical weights and N the column densities (the star denotes the excited level). Naturally, because different mechanisms other than the CMBR contribute to populating the fine-structure excited level, this $T_{\rm exc}$ is just an upper limit to the actual $T(z)$.

The first observations leading to $T(z)$ were very uncertain and just set an upper bound of 45 K at $z = 2.5$ by using C II lines (Bahcall, Joss & Lynds 1973). Later Meyer et al. (1986) used C I lines to place a limit of 16 K in an absorption system at $z = 1.776$, which was re-observed by Songaila et al. (1994). With a total exposure time of 13 h at the Keck I telescope, their spectra allowed one not only to place a more accurate upper limit in the CMBR temperature but also to separate the system into two different components.

During the last few years new results have become available in the literature, although most of them are still upper limits. These results are plotted in Fig. 1, where the references are also listed. Except for two absorption systems ($z = 2.652$ and 2.9034) that are Lyman-limit systems (LLS), all the others are damped Lyα systems (DLS). For three systems ($z = 1.721$,1.77638 and 1.77654) the results come from C I lines, while for the others, all at $z > 2.5$, $T(z)$ comes from C II lines. Notice that for the system at $z = 1.9731$ both C I and C II lines are observed (Ge, Bechtold & Black 1997). Except for the system at $z = 2.652$, the values for the high-z systems are upper limits, because either the absorption line arising from the C^+ excited state is too weak to be

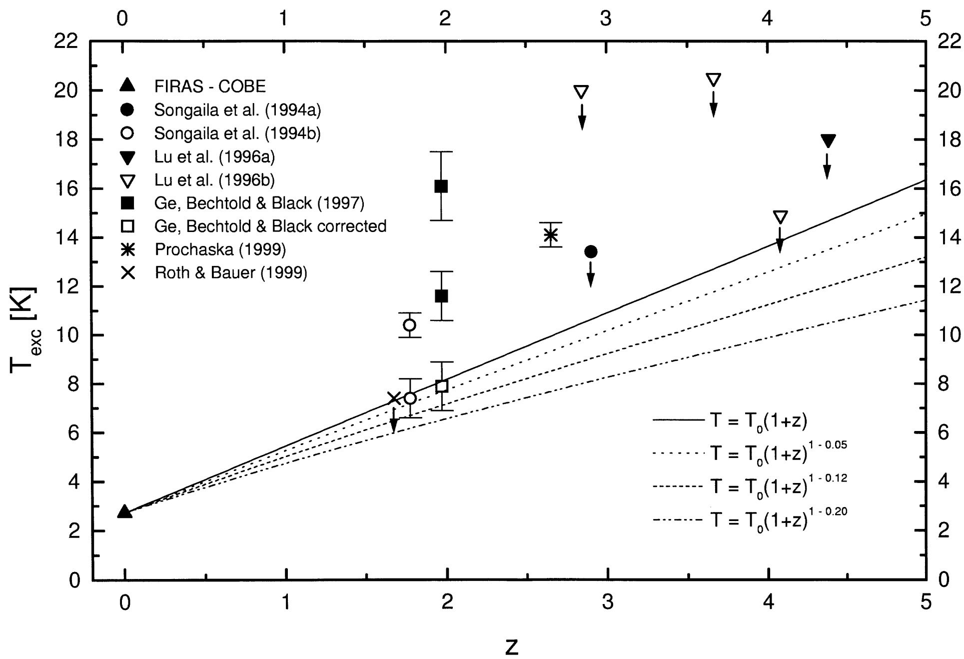

Figure 1. The temperature–redshift relation. The lines correspond to the CMBR temperature given by the standard model (solid line) and models including photon creation characterized by the parameter β (see text) equal to 0.05 (dotted line), 0.12 (dashed line) and 0.20 (dot–dashed line). The points correspond to the excitation temperatures obtained from the column density ratios of absorption lines from the fundamental and first excited levels of C I ($z < 2$) and C II ($z > 2$); error bars are 1σ while upper limits are 2σ.

measured, given the spectral signal-to-noise ratio available, or the ground-state line is too saturated to yield accurate column densities.

In Fig. 1 we show the theoretical prediction of the CMBR temperature, $T(z)$, for all photon-conserving hot big bang scenarios, as well as in the more general framework given by equation (16). The error bars and data points correspond to excitation temperatures that may be converted to the CMBR temperature, after correction for other mechanisms populating the levels (collisions and fluorescence).

4 PHYSICAL MECHANISMS AND TEMPERATURE CORRECTIONS

In the following we discuss the possible corrections to the excitation temperatures and their effect on $T(z)$. Details of the calculations and the atomic data will be presented soon (Silva & Viegas, in preparation).

Neutral carbon is present mainly in neutral low-temperature gas. This, in addition to the CMBR, collisional excitation by H atoms, and fluorescence through higher levels caused by galactic radiation, may populate the C^0 excited state. The collisional effect is illustrated in Fig. 2, showing the C^0 population ratio as a function of the excitation temperature, for different values of the neutral hydrogen density, assuming a gas temperature of 100 K. Two points are plotted for each C I system considering to the same population observed ratio but two different CMBR temperatures. One (on the right) neglects the effect of collisions, i.e. gives the blackbody temperature necessary to explain the population ratio. The other (on the left) corresponds to the standard CMBR temperature at the redshift of the system, which is too low to explain the observed ratio; thus, collisional excitation is necessary and for each system the value of the H density is different, reaching about 45 cm^{-3} for the system at $z = 1.9731$. For this system, Ge et al. (1997) have already proposed a temperature correction (see Fig. 1) assuming H collisions, with $n_{\mathrm{H\,I}} = 21\ cm^{-3}$ and fluorescence resulting from a strong galactic field with an intensity a factor of 17 higher than the Milky Way radiation field (Black 1987). If fluorescence is taken into account in the population calculation, the higher the field intensity the lower the $n_{\mathrm{H\,I}}$ value necessary to reproduce the C^0 population rate. As these C I systems are damped Lyman systems, usually associated with the disc of galaxies, the natural expectation is a higher volumetric density. Hence, the $n_{\mathrm{H\,I}}$ could be as high as indicated in Fig. 2. In this case, if the effect of the galactic radiation field is not negligible, the CMBR temperature would be lower than the value predicted by the standard cosmology at the redshift of the system.

Concerning the C II systems, a similar analysis can be performed. The population ratio versus the temperature is shown in Figs 3 and 4, corresponding to the physical conditions in DLS and LLS, respectively. The effect of fluorescence is negligible for the densities involved. For the Lyman-limit systems, a gas temperature of 10^4 K is assumed and the main excitation mechanism is electronic collisions. As in Fig. 2, each system is represented by two points. Most of the DLS data are upper limits so that the result shown in Fig. 3 is somewhat indicative, and points only to an upper bound for $n_{\mathrm{H\,I}}$. The only 'observed' ratio corresponds to the LLS at $z = 2.652$, for which the electronic density indicated in Fig. 4 is about 0.07 cm^{-3}. This value is higher than the one assumed in models where the LLS is photoionized by the QSO background radiation. However, such models lead to LLS sizes that are too high (Gruenwald & Viegas 1993). An alternative model accounting for a hot halo radiation field

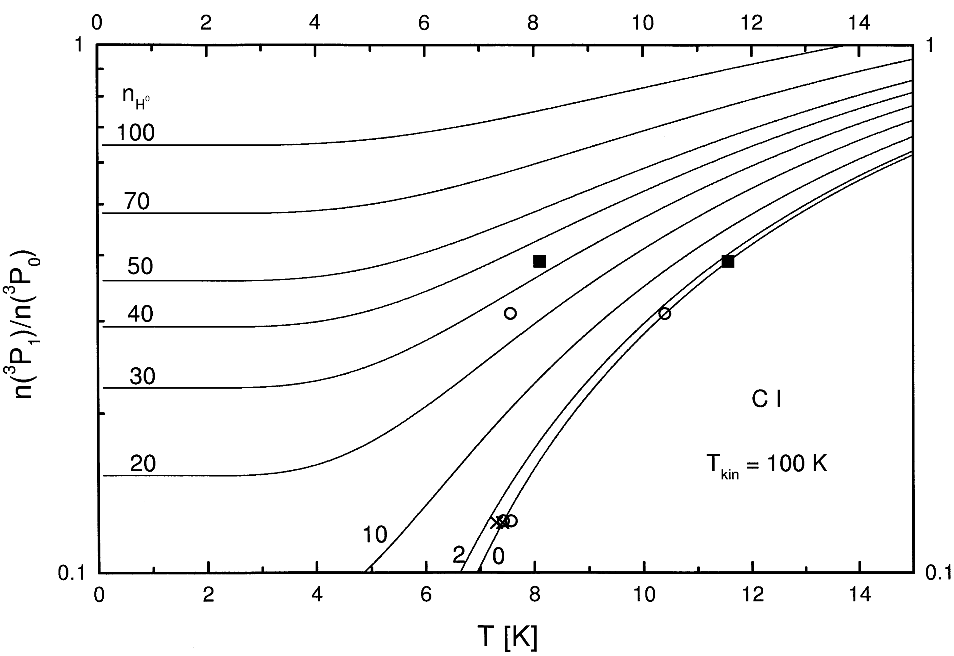

Figure 2. The observed ratio between the population of the first excited and fundamental levels of C^0 as a function of the excitation temperature. The lines are labelled by the value of $n_{\mathrm{H\,I}}$. The points correspond to the observational data for the C I lines, with the same notation as used in Fig. 1. For each absorption system two points are plotted: on the right the result corresponding to $n_{\mathrm{H\,I}} = 0$ (equivalent to the excitation temperature), and on the left the result corresponding to the coupled effect of the CMBR (with the temperature fixed at the value given by the standard model at the redshift of the system) and collisional excitation by H I.

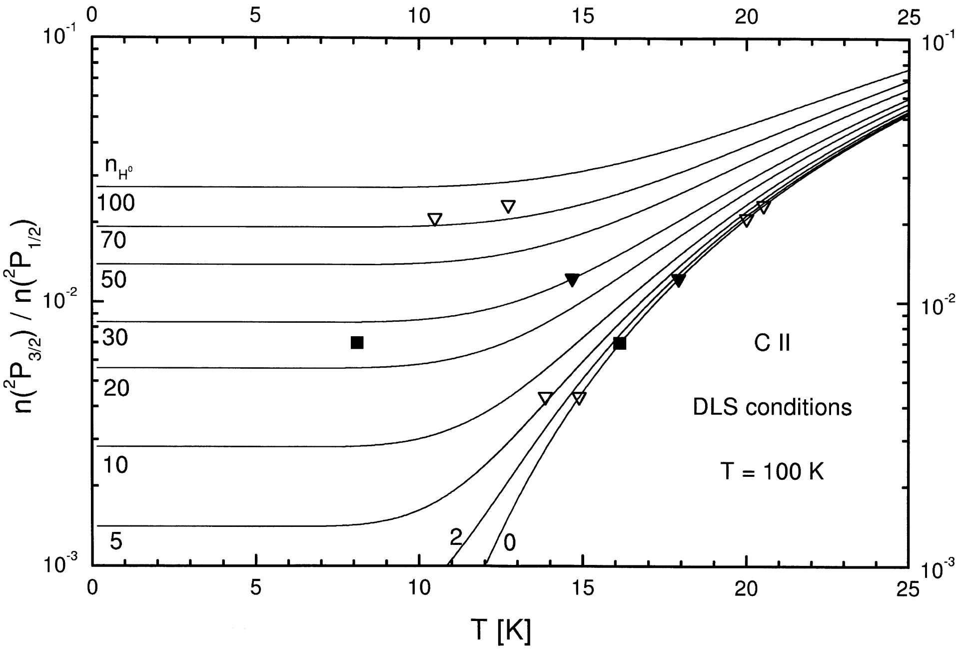

Figure 3. The observed ratio between the population of the first excited and fundamental levels of C^+ as a function of the excitation temperature. The observational data come from damped Lyman systems, assuming a gas temperature equal to 100 K. The curves are labelled by $n_{\rm H\,I}$. As in Fig. 2, for each absorption system two points are plotted.

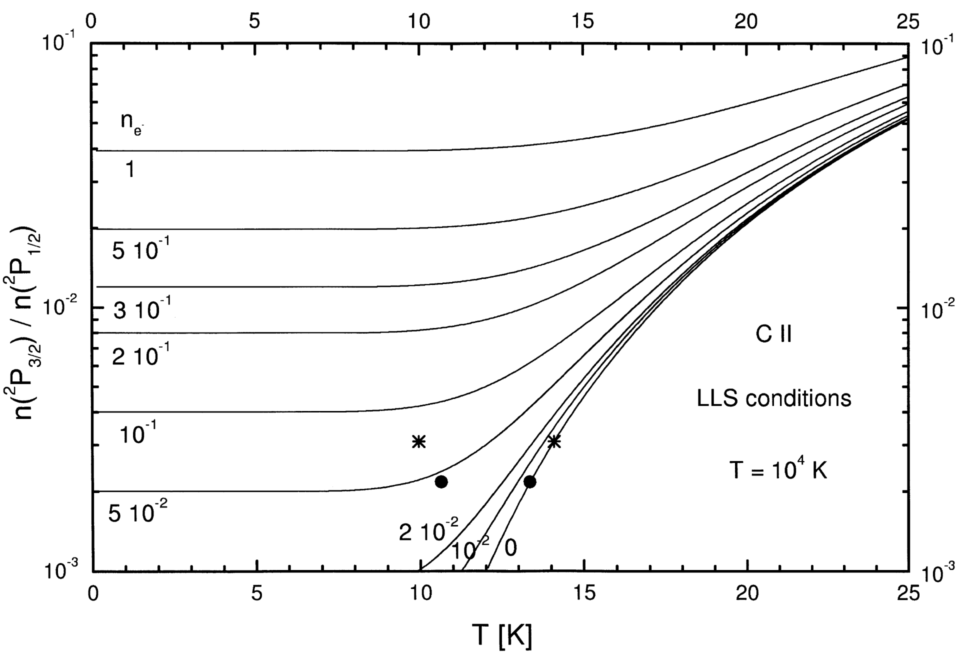

Figure 4. Same as Fig. 3 for Lyman-limit systems, assuming a gas temperature equal to 10^4 K. The curves are labelled by n_e. As in Fig. 2, for each absorption system two points are plotted.

indicates that the system densities may be higher (Viegas & Friaça 1995; Viegas, Friaça & Gruenwald 1999), usually in the range 0.1 to 10 $\rm cm^{-3}$. In this case, the CMBR temperature would also be lower than the standard cosmology value at $z = 2.652$.

Molecular rotational transitions could also provide a $T(z)$ estimate. However only four QSO absorption systems were found showing molecular transitions (Wiklind & Combes 1994, 1995, 1996a,b). $T(z) = 4 \pm 2$ K (the error is 3σ) could be obtained for one of these systems at $z = 0.89$ (Wiklind & Combes 1996b). This value is in agreement with the standard cosmology prediction $[T(0.89) = 5.16\,\rm K]$ only at 2σ confidence level, and may also indicate the need for an alternative model.

As $T(z)$ increases with z, more reliable conclusions would be possible if more precise data at high z became available. This implies better determination of the column densities of the C II absorption lines and an effort to detect the C I absorption lines in damped Lyman systems at $z > 3$, which is presently a challenge to the observers. This brief analysis of the data indicates that the CMBR temperature may be lower than the value predicted by the standard cosmology, consequently opening the possibility of alternative (big bang) models of the Universe. We stress that if the analysis presented here is confirmed by more accurate observations, the standard model, as well as other models in vogue like conserved 'quintessence' or ΛCDM models, will be ruled out by this kind of test.

ACKNOWLEDGMENTS

This work was partially supported by the project PRONEX/FINEP (No. 41.96.0908.00), Conselho Nacional de Desenvolvimento Científico e Tecnológico – CNPq (Brazilian Research Agency) and FAPESP (No. 97/00719-0). We are grateful to Dr Katherine C. Roth for helpful information.

REFERENCES

Bahcall J. N., Wolfe R. A., 1968, ApJ, 152, 701
Bahcall J. N., Joss P. C., Lynds R., 1973, ApJ, 182, L95
Black J. H., 1987, in Hollenbach D. J., Thronson H. A., Jr, eds, Interstellar Processes. Reidel, Dordrecht
Caldwell R. R., Dave R., Steinhardt P. J., 1997, Phys. Rev. Lett., 80, 1582, 1998
Calvão M. O., Lima J. A. S., Waga I., 1992, Phys. Lett. A, 162, 233
Carvalho J. C., Lima J. A. S., Waga I., 1992, Phys. Rev. D, 46, 2404
Chen W., Wu Y. S., 1990, Phys. Rev. D, 41, 695
Fixsen D. J. et al., 1996, ApJ, 473, 576
Freese K., Adams F. C., Frieman J. A., Mottola E., 1987, Nucl. Phys. B, 287, 797
Friemann J., Hill C. T., Stebbins A., Waga I., 1995, Phys. Rev. Lett., 75, 2077
Ge J., Bechtold J., Black J. H., 1997, ApJ, 474, 67
Gruenwald R., Viegas S. M., 1993, ApJ, 415, 532
Lima J. A. S., 1996, Phys. Rev. D, 54, 2571
Lima J. A. S., Alcaniz J. S., 1999, A&A, 348, 1
Lima J. A. S., Germano A. S. M., 1992, Phys. Lett. A, 162, 233
Lima J. A. S., Maia J. M. F., 1994, Phys. Rev. D, 49, 5597
Lima J. A. S., Trodden M., 1996, Phys. Rev. D, 53, 4280
Lima J. A. S., Germano A. S. M., Abramo L. R. W., 1996, Phys. Rev. D, 53, 4287
Mather J. C. et al., 1994, ApJ, 420, 439
Meyer D. M., Black J. H., Chaffee F. H., Foltz C., York D. G., 1986, ApJ, 308, L37
Overduin J. M., Cooperstock F. I., 1998, Phys. Rev. D, 58, 043506
Ozer M., Taha M. O., 1986, Phys. Lett. B, 171, 363
Ozer M., Taha M. O., 1987, Nucl. Phys. B, 287, 776
Prigogine I., Geheniau J., Gunzig E., Nardone P., 1989, Gen. Rel. Grav., 21, 767
Ratra B., Peebles P. J. E., 1988, Phys. Rev. D, 37, 3407
Roth C. K., Meyer D. M., 1995, ApJ, 441, 129
Roth C. K., Meyer D. M., Hawkings I., 1993, ApJ, 413, L67
Roth C. K., Bauer J. M., 1999, ApJL, submitted
Silveira V., Waga I., 1997, Phys. Rev. D, 56, 4625
Smoot G. F. et al., 1992, ApJ, 396, L1
Songaila A. et. al., 1994, Nat, 371, 43
Viegas S. M., Friaça A. C. S., 1995, MNRAS, 272, L35
Viegas S. M., Friaça A. C. S., Gruenwald R., 1999, MNRAS, submitted
Wiklind T., Combes F., 1994, A&A, 286, L9
Wiklind T., Combes F., 1995, A&A, 299, 382
Wiklind T., Combes F., 1996a, A&A, 315, 86
Wiklind T., Combes F., 1996b, Nat, 379, 139
Zimdhal W., Pavón D., 1994, MNRAS, 266, 872

This paper has been typeset from a TEX/LATEX file prepared by the author.