POSSIBLE THERMAL HISTORIES OF INTERGALACTIC GAS

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ABSTRACT

By using simple expressions to represent the time dependence of the heating of intergalactic gas, the thermal history of an initially cold gas is followed and the resulting radiation field is computed. If one requires that the gas must have been strongly heated and ionized at a redshift of about 2 to avoid strong Ly-α absorption in the spectra of QSO (here assumed to be cosmologically distant) then, at a final gas density of the order of $10^{-5}/\text{cm}^3$, inadmissibly large amounts of X-ray emission in the 5–10-Å region can be avoided only for a very narrow range of heating parameters. Measures at 25–30 Å with comparable sensitivity would be decisive. Models with a high degree of clumpiness in the gas can already be excluded.

Koehler's reported value of an intergalactic hydrogen density of $n_{\text{H}} \sim 10^{-7}$ cannot be reconciled with any of the models investigated, and it appears that one either has to drop the assumption of cosmological distances for the QSO or else consider a model in which the ionization is caused by a burst of star formation in the early stages of galaxy formation. However, a specific model of this type, assuming stars of $T_{\text{eff}} = 21000^\circ$ K encounters the difficulty that the total light from the galaxies at this epoch should be far more intense than the observations allow.

I. INTRODUCTION

In this paper we wish to examine the possible thermal histories of a smooth component of intergalactic gas of present density $n \sim 2 \times 10^{-5}$ particles/cm^3 and to compute the resulting radiation fields.

A recent paper by Ginzburg and Ozernoy (1966) considered the run of temperature versus density for intergalactic gas under the influence of various possible heating mechanisms together with the effects of cooling by radiation and the general cosmic expansion. It was suggested previously (Weymann 1966; referred to hereinafter as "Paper I") that calculations similar to Ginzburg and Ozernoy's ought to be repeated to include the effects of line radiation, since this significantly alters the radiative properties of hydrogen. In addition, Gould and Ramsay (1966) have recently called attention to the significant modification of the thermal properties of the gas that result from the presence of helium and have considered the resulting thermal and ionization *equilibrium*.

It now appears (Howell and Shakeshaft 1966; Thaddeus and Clauser 1966) that over the range of about 3 mm to 21 cm the departures from a Planck function of the "cosmic background radiation" are very slight. Accordingly, from Paper I we would conclude that no significant heating of the gas took place prior to the decoupling between the matter and the radiation which occurs when the gas becomes neutral at about 3000° K. Thus any subsequent heating acted upon gas which was originally *cold* and *neutral*.[1] Reference should also be made to the similar work of Rees and Sciama (1966), who considered the temperature and degree of ionization likely in a *steady-state* cosmology (though we feel such a cosmology to be excluded by the 3° K phenomenon).

There are at least three possible points of contact between observations and the models considered here: (1) Isotropic X-ray emission in the 25–1-Å region (Field and Henry 1964); (2) 21-cm absorption measures of distant radio sources (Koehler 1965); (3) Ly-α absorption (or rather the lack thereof) in the spectra of QSO (Gunn and Peterson 1965). At the present time the question of whether these objects are actually at cosmological distances is open; we shall assume for the rest of this paper that they are in fact cosmologically distant.

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[1] It would seem also quite likely therefore that some of the more interesting phases of condensation and fragmentation of galaxies took place under these conditions.

II. ASSUMPTIONS AND COMPUTATIONAL DETAILS

a) Basic Equations and Heat-Input Assumptions

The basic equation governing the radiation field is given in Paper I (eq. [3]) by

$$\frac{\partial U_\nu(\nu_0,t)}{\partial t} = -3\,\frac{U_\nu}{R}\frac{dR}{dt} + \epsilon_\nu(\nu_0,t)\,, \tag{1}$$

where $U_\nu(\nu_0,t)$ is the energy density of photons whose frequency is ν_0 at the present epoch, ϵ_ν is the *net* radiative emission coefficient (in ergs/cm^3/sec/Hertz), and $R(t)$ is the usual cosmological expansion parameter in the standard form of the line element. For the equation governing the temperature of the gas we have (cf. eq. [4] in Ginzburg and Ozernoy 1966)

$$\frac{dU}{dt} = \frac{P}{\rho^2}\frac{d\rho}{dt} + (E_H - E_R)\,. \tag{2}$$

Here E_H represents the heat input per *gram*, E_R is the net rate of radiative cooling per gram, and U is the internal energy per gram, including energy of ionization, of the matter.

The equations governing the state of ionization are

$$\frac{dx}{-d\ln n} = \left|\frac{n}{dn/dt}\right| \{n_e[(1-x)R_1 - xC_1] - x\Gamma_1\}\,, \tag{3a}$$

$$\frac{dy}{-d\ln n} = \left|\frac{n}{dn/dt}\right| [n_e(zR_2 - yC_2) - y\Gamma_2]\,, \tag{3b}$$

$$\frac{dz}{-d\ln n} = \left|\frac{n}{dn/dt}\right| \{n_e[(1-y-z)R_3 - zR_2 + yC_2 - zC_3] + y\Gamma_2 - z\Gamma_3\}\,, \tag{3c}$$

where x is the fraction of hydrogen nuclei that is neutral, y the fraction of helium nuclei that is neutral, and z the fraction that is singly ionized; R, C, and Γ are the recombination, collisional-ionization, and photo-ionization rate coefficients, the latter being given by

$$\Gamma_i = \int_{(\nu_0)_i}^{\infty} \frac{4\pi J_\nu a_\nu d\nu}{h\nu}\,.$$

These equations are derived by simply applying the continuity equation of hydrodynamics to the ion species in question as well as to the total nucleon number.

Finally, the rate of expansion is taken from Paper I, equation (19) (with β, the ratio of the energy density of radiation to the total rest energy of the matter, set equal to zero—a valid approximation for the stages in the expansion of interest). The parameter K appearing in this equation can be expressed in terms of the familiar deceleration parameter q_0 by the relation $K = 1 - \frac{1}{2}q_0$, and throughout most of this work we have considered the case $q_0 = +1$.

The possible sources of heat input and their variation with time have been discussed by Ginzburg and Ozernoy as well as by Gould and Ramsay. Among the mechanisms considered by these authors are the "ionization losses" by relativistic and subrelativistic particles as they pass through the gas, dissipation of plasma jets and shock waves generated by explosions within galaxies, and dissipation of turbulent motions in the intergalactic medium.

To this might be added a fourth[2]—the heating of the gas by the *supersonic* passages

[2] I am grateful to Dr. E. A. Spiegel for a discussion of this mechanism.

of galaxies or protogalaxies through the gas. Obviously we are very ignorant of any of the details—or even the fact of the existence—of any of these mechanisms. Therefore the point of view taken here is simply to describe the heating mechanism by either one of two very simple formulae.

Case a.—For densities larger than n_0 the heating rate *per gram* was taken to be zero. Between n_0 and n_1 the heating rate was taken as a constant and for densities less than n_1 was taken to be zero again. Such a model would correspond roughly to the heating produced by outbursts from active galaxies, whose main activity is confined to a distinct epoch in the evolving universe, as well as to the Spiegel model. However, in this latter case there is a temperature dependence upon the heating, since if the gas were to be heated to the point where the peculiar motion of the galaxies became subsonic the heating would markedly drop. If this mechanism were sufficiently powerful, we should then expect the gas to be kept at a temperature such that $V_{\text{sound}} \sim V_{\text{peculiar}}$. Since the peculiar velocities of non-relativistic particles drop as $1/R$, this would imply a $T \sim n^{2/3}$ dependence. Taking the present epoch peculiar velocities to be of the order of 100 km/sec, we would thus expect present epoch temperatures of the order of $T \sim 10^6$.

Case b.—For densities less than 2 n_0, the heating is zero, commences at a value of $\frac{1}{2}E_H$, and increases inversely as the density to a value of E_H at a density n_0 and thereafter drops directly as the density. This latter phase corresponds to a situation in which high-energy particles heat the gas, with the depletion of these particles being negligible. (In the case of the subrelativistic particles, however, the heating rate will actually drop less steeply than in this model because, as the energy of the individual particles drops during the expansion, the particles heat more effectively). The initial rise in the heating corresponds to the spread in the heating time discussed below.

Most of the results described below are based upon case *a*, but we have also experimented with case *b*.

Under the assumption of constant heat input per gram the relative importance of heating compared to radiative cooling *increases* as the density drops, since the radiative cooling per gram is proportional to the density. The temperature dependence of the radiative cooling function, E_R (see the discussion below and Fig. 1) is characterized by a rather sharp maximum at moderately low temperature, arising mostly from collisional excitation of Ly-α. The gas cannot reach high temperatures until the heating rate exceeds this maximum radiation rate. Starting from a cold gas at density n_0, the heating will quickly raise the gas to a temperature of around 10000° K where it slowly rises until it jumps this "thermal barrier" at around 20000° K (for pure hydrogen) and quickly runs to higher temperature, giving tracks as in Figure 2 below (§ III). In this simple model then, for given E_H, the value of the turnon density, n_0, has virtually no effect on any observable properties provided only that this turnon density is earlier than the density at which the thermal barrier is surmounted for a given E_H.

For case *b*, however, the situation is different in that radiative cooling and heating keep pace with each other beyond n_0. Thus, when E_H is less than a certain critical amount for given n_0, the gas will never surmount the thermal barrier. While the choice of the heating law could be quite general, it is felt that it must conform to the following restriction: It cannot be sharply peaked in time but must be spread over a time corresponding to the interval in which the mean density drops by a factor of the order of 2–4. This restriction is based upon the following points: (1) Whatever the heating mechanisms are (certainly including all the ones mentioned above except possibly the third) they very likely involve the development of galaxies and protogalaxies. (2) Following in a rather general way the ideas developed by Peebles (1965) it is supposed that density fluctuations develop in the course of the expansion with a characteristic spread in amplitude ranging from ϵ to $\epsilon/2$. However, these are arrested from further growth by some process until a critical epoch n_0 when some of the matter (under the present assumption only a very small fraction in fact) further condenses to galaxies or clusters of galaxies.

The spread in condensation times (and accordingly the minimum duration of significant heating) is then estimated on the basis of an extremely simple model as follows:

We ignore all forces except gravity entirely; consider a cosmological model with K in equation (19) of Paper I equal to zero. Consider at some instant t_0 a spherical perturbation with density excess $\epsilon = (\delta\rho/\rho)_0$ but with no perturbation in the velocity field. Let $\xi(t)$ be the separation of two test particles in the perturbed flow, and $\eta(t)$ be the corresponding quantity in the unperturbed flow. Letting time be measured in units of $(\frac{8}{3}\pi G\rho_0)^{-1/2}$, we can write, on the basis of simple Newtonian mechanics for radially moving particles,

$$\left(\frac{d\eta}{dt}\right)^2 = 1/\eta\,, \tag{4}$$

$$\left(\frac{d\xi}{dt}\right)^2 = \frac{1+\epsilon}{\xi} - \epsilon\,, \tag{5}$$

or

$$\eta^{1/2}\, d\eta = \pm[\xi/(1+\epsilon-\epsilon\xi)]^{1/2}\, d\xi\,. \tag{6}$$

For positive ϵ, ξ will expand to a maximum of $\xi_{\max} = (1+\epsilon)/\epsilon$ and then collapse to a singularity. Integrating equation (6) over these two phases, one finds that the value of η, η_s, at the singular point is given by

$$\eta_s = \left[\!\!\left[1 + \tfrac{3}{2}\epsilon^{-1/2} \left\{ \frac{\pi}{2} + \epsilon^{-1/2} + \left(\frac{1+\epsilon}{\epsilon}\right)^{1/2} \sin^{-1}[\,(1+\epsilon)^{-1/2}\,] \right\} \right]\!\!\right]^{2/3}, \tag{7}$$

so that there has been a drop in the mean density by a factor, ρ_s, of

$$\rho_s = \eta_s^{-3}. \tag{8}$$

We identify the value of $\rho_s(\epsilon)/\rho_s(\epsilon/2)$ as the spread in the mean density during which the bulk of the condensation took place—this involves the assumption that the dispersion in amplitude of the density fluctuations is comparable with their mean amplitude. For $\epsilon = 1$, this density ratio is 2.8 and increases smoothly to 4.0 for small amplitudes. We thus adopt as a conservative estimate, that the final condensation of the bulk of the matter—and hence the minimum duration of significant heating—was spread over a time during which the mean density dropped by *at least* a factor of 2.0.[3]

For the constant heating rate (case *a*) we thus take the total duration of the heating period to correspond to a density drop of a factor of at least 2: $n_0/n_1 \gtrsim 2$, while in case *b* the rise to maximum occupies a similar time—the time during which the supply of cosmic-ray particles is built up.

Finally, it should be explicitly stated that it is assumed that all ionizations are due either to (thermal) collisional ionizations or to photo-ionizations from photons generated by the hot gas itself. We are thus ignoring the quite plausible possibility that a high degree of ionization *without corresponding heating of the gas to high temperature* can occur by means of a large burst of star formation involving only moderately hot stars which marked the first epoch of galaxy condensation. We return to this possibility in § III.

[3] Regarding the epoch at which condensation actually occurred, it has been argued by many investigators that the largest stellar orbits in our Galaxy suggest a "separation" epoch of $n \sim 10^{-3}$. This corresponds to a redshift under our assumptions of about 2.5. At the present time no QSO are known with redshifts of this amount or larger which is also suggestive. If we took $n = 10^4$ as the density at which the condensation could begin to freely proceed due to the disappearance of electrons, we would then infer very small fluctuations at this epoch, $\epsilon \sim 10^{-3}$.

b) Adopted Cross-Sections and Computational Methods

We consider the contribution by hydrogen to the term E_R to be given by

$$n_e n_p B_1(T) + n_e n_H S_1(T) - n_H \int_{\nu_0}^{\infty} 4\pi J_\nu a_\nu d\nu , \qquad (9)$$

where n_p is the proton density, $B_1(T)$ is the total rate of radiation due to free-free and free-bound radiation (including the resulting line emission), and $S_1(T)$ is the energy loss rate due to collisional excitation of Ly-α. Expressions of the same form hold for neutral and singly ionized helium.

All collisional excitation and radiative excitation processes are supposed to take place from the ground state, and the excited levels have been represented by single levels of energy 10.2, 20.55, and 40.8 eV for H, He I, and He II, respectively.

Because the expressions used for the cross-sections are taken from such a variety of sources, it seems best simply to present the resulting values of the cooling curves and the corresponding ionization equilibria in graphical form. This is done in Figure 1, which

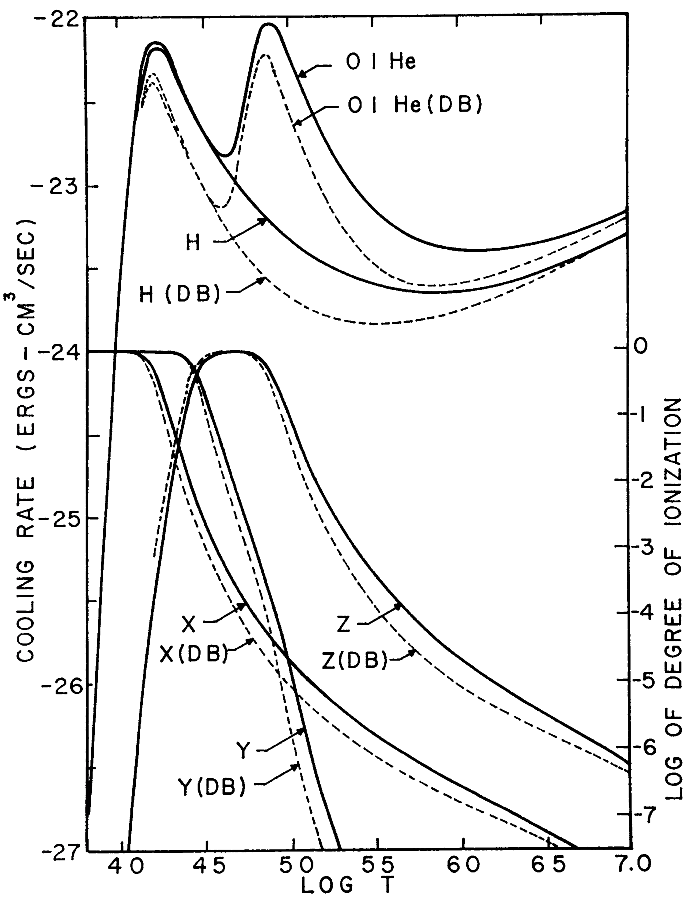

FIG. 1.—Cooling rates and ionization equilibria resulting from the adopted cross-sections. The top four curves (*left scale*) give the cooling rate as a function of temperature for pure hydrogen and a mixture of 10 per cent helium and 90 per cent hydrogen. The lower six curves (*right scale*) give the ionization equilibria set up by collisions and recombinations. X is the fraction of hydrogen nuclei neutral, Y the fraction of helium nuclei neutral, and Z the fraction singly ionized. In all cases the dashed curves marked "D.B." assume detailed balance in the ground state continua.

shows the cooling curves for pure hydrogen and also for 10 per cent helium by number for ionization levels corresponding to the ionization equilibria determined solely by collisional ionization and recombination. Dielectronic recombination for $He^+ \rightarrow He^o$ has *not* been included.

A large grid of frequencies spaced by an amount $\Delta \log \nu = 0.05$ and running from 1 Å to 10 μ (measured at the present epoch) was set up, and the change in intensities of each band was followed as it became redshifted. At these low densities the redistribution in frequency due to Compton scattering can be ignored. The resonance-line radiation arising from collisional excitation and recombination was assumed to escape to the red wing of the line. The line emission was assigned to the point of the grid immediately to the red of the line, with an intensity per unit frequency equal to the total intensity of the line divided by the frequency interval to the next grid point. The integrals appearing in equations (3) and (9) were numerically evaluated using the trapezoidal rule. Equation (1) can be written in the form[4] (eq. [4], Paper I)

$$\frac{n}{N}\frac{\partial N}{\partial n}(\nu_0,n) = -\frac{(j_\nu - \kappa_\nu' J_\nu)4\pi R}{3U_\nu(dR/dt)}, \tag{10}$$

where $N(\nu_0,n)$ is the occupation number, and n is the mean density. When we define $(\Delta\tau)_{\nu_0}$ by

$$(\Delta\tau)_{\nu_0} = -\Delta(\ln n)\,\kappa' c/(3R^{-1}dR/dt)$$

and S by

$$S = j\Big/\left(\frac{2h\nu^3\kappa'}{c^2}\right),$$

equation (10) assumes the usual form of the transfer equation. If $(\Delta\tau)_{\nu_0}$ in the ground-state continuum is very large compared to unity over a single integration step, then we assume detailed balance (or in the language of gaseous nebulae we use the on-the-spot approximation) and put $N_\nu = S_\nu$ for those frequencies, so that in such situations only the recombinations to the excited state have any net effect, as noted by Rees and Sciama.

III. RESULTS AND DISCUSSION

In order to test the compatibility of the existence of a gas of the density assumed above with the models of the heating assumed above, we give in Figures 2, 3, and 4 a series of graphs corresponding to the constant heat input per gram model. The turnon time in this case is, as mentioned above, not relevant—the turnoff time has been varied, while the intensity of the heating was as low as possible, consistent with the attainment of a high degree of ionization of the hydrogen at the epoch demanded by the Ly-α observations. Figure 2 gives the resulting run of temperature-density paths. The paths after heating turnoff are indistinguishable from adiabats under these conditions. Figure 3 shows the resulting radiation field at the present epoch. The large peak is due to Ly-α smeared by the redshift.

The relevant X-ray observations have been summarized by Field and Henry (1964). The positive result reported in the 2–8-Å region is indicated in Figure 3 but is regarded as only an upper limit in view of the fact that it is quite likely to be composed of unresolved discrete sources (Byram, Chubb, and Friedman 1966).

Figure 4 represents a kind of fictitious Ly-α absorption spectrum. It gives simply the absorption by intergalactic hydrogen at different epochs corresponding to various redshifts. It can be thought of as the resulting spectrum of a distant QSO (with a normalized continuum), provided attention is confined to rest wavelengths between 1216 and 912 Å. The effects of Lyman continuum absorption are not included in Figure 4. The Ly-α

[4] The $1/N$ on the left-hand side of eq. (4) of Paper I was inadvertently omitted.

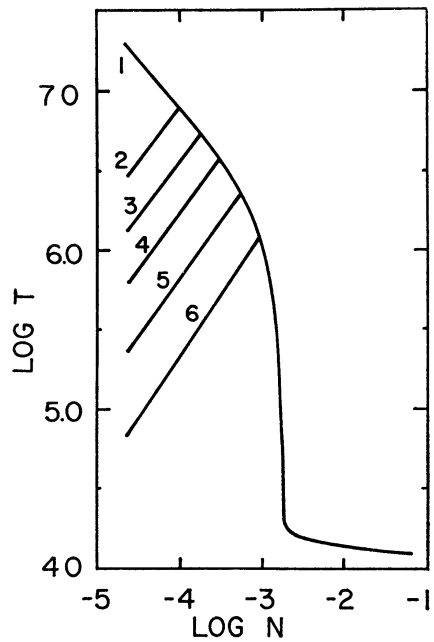

FIG. 2.—Thermal history of intergalactic gas having a final density of about 2×10^{-5} particles/cm³, computed with the case *a* heating model. The level of heating is the minimum required to assure negligible Ly-α absorption in the QSO spectra. The six cases refer to different epochs at which the heating is switched off.

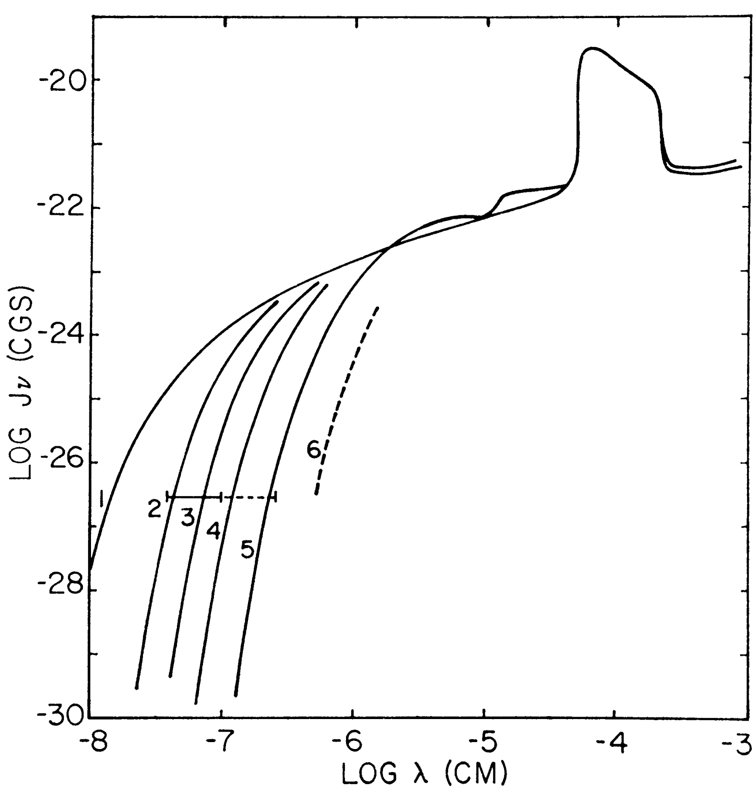

FIG. 3.—The resulting radiation field for the six cases of Fig. 2. The large hump at around 1 μ is redshifted Ly-α resulting from recombination mostly before the jump to very high temperatures. The short solid horizontal line is the approximate limit on X-rays, excluding cases 1, 2, and probably 3 Cases 4 and 5 are permitted but could be excluded if the observations could be extended to 25–30 Å.

absorption discussed by Gunn and Peterson (1965) in 3C 9 is marginal at best, and no convincing evidence for a drop in the continuum seems to have been found in the spectra of other objects with comparable redshifts. For the purposes of this paper we will assume that the optical depth in Ly-α at a redshift of 2.0 is less than 0.5.

To see the effects of helium we show case 3 for pure hydrogen along with a nearly identical set of heating parameters for a 10 per cent addition of helium in Figures 5, 6, and 7.

The second pause in the helium case (Fig. 5) results from the presence of two separate local maxima in the cooling curve for the helium and hydrogen mixture (see Fig. 1). The most conspicuous additional feature in Figure 6 arising from the helium is the He^+ Ly-α line. Lying between about 1500 and 3000 Å, this feature is so bright *that it could very*

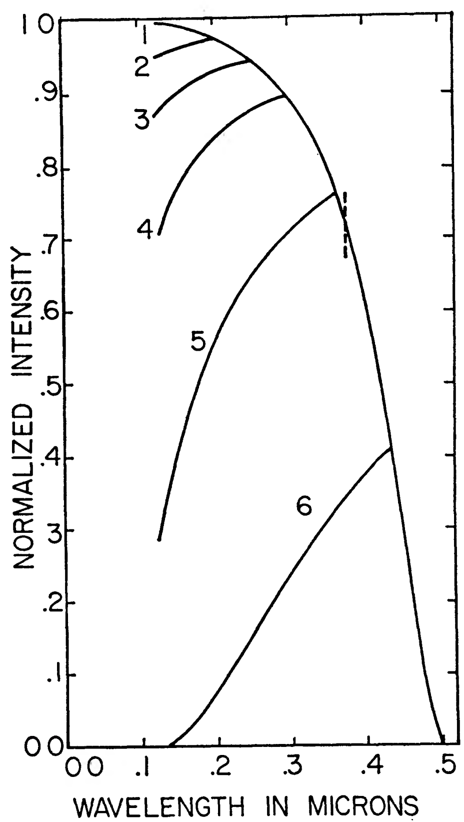

FIG. 4.—Absorption due to intergalactic Ly-α absorption as a function of wavelength. For a distant object, the wavelength interval between the object's rest wavelengths of 912 Å and 1216 Å is given by the curves. The short vertical dashed line gives the position of Ly-α in the most distant QSO so far observed. On this basis, cases 1–5 would be permitted but case 6 would be excluded.

possibly be detected by future OAO experiments provided one could look between individual early-type stars at the galactic poles. The X-ray emission is not very different in the two cases but the He^+ Ly-α line causes the hydrogen ionization to stay very high until this feature has been redshifted past 912 Å.

We may now summarize the answer to the following question: assuming the validity of the constant rate heating model and the existence of a smooth component of intergalactic hydrogen with present density $2.24 \times 10^{-5}/cm^3$, can we find a set of heating parameters which give thermal histories which are not in conflict with both the X-ray data and the lack of strong Ly-α absorption? (And, as always, assuming the QSO to be cosmological.)

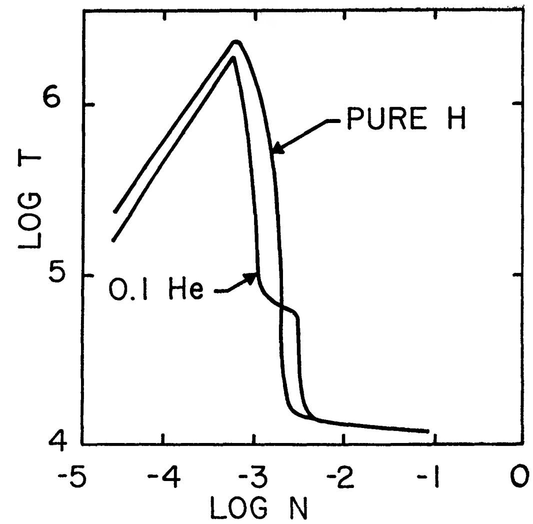

FIG 5.—As in Fig. 2, where case 3 of Fig 2 (pure hydrogen) is compared with similar heating applied to a mixture with 10 per cent helium.

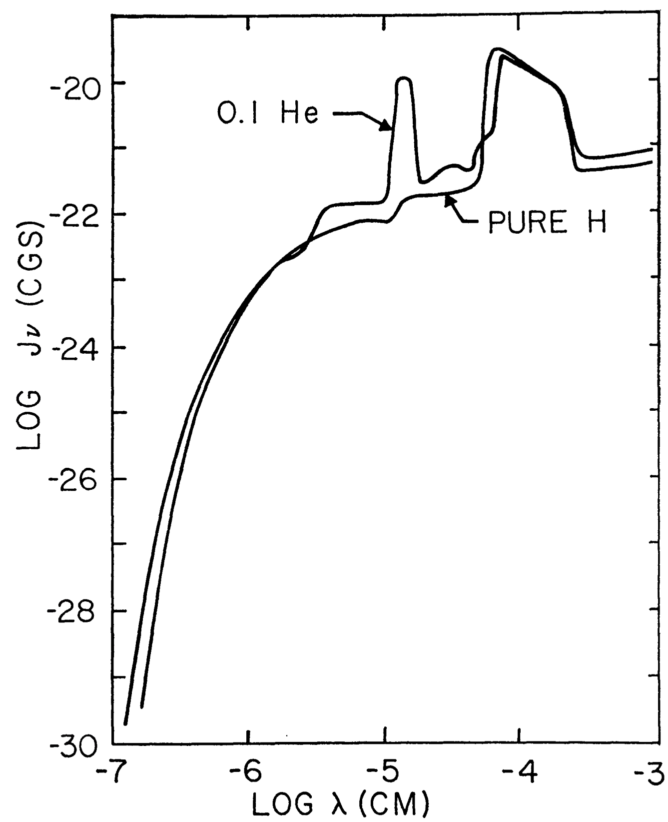

FIG. 6.—The resulting radiation fields for Fig. 5. The sharp spike longward of 1000 Å is due to He^+ Ly-α.

At the present time, cases 1, 2, and 3 (Fig. 2) are excluded by the X-ray data. Case 6 can be excluded by the Ly-α absorption. In this case, the heating is turned off so quickly after the thermal barrier is surmounted that the gas never attains a high enough temperature to become thoroughly ionized, and the recombination accompanying the subsequent drop in temperature leads to unacceptably large amounts of absorption (see Fig. 4). (In addition, shorter turnoff times would be excluded by the demand that the spread in heat input time not be too small.)

Cases 4 and 5 are still permitted. As Field and Henry have also emphasized, because of the exponential factor in the free-free emission, it would be highly desirable if the background limit could be pushed to longer wavelengths (at the same sensitivity as the

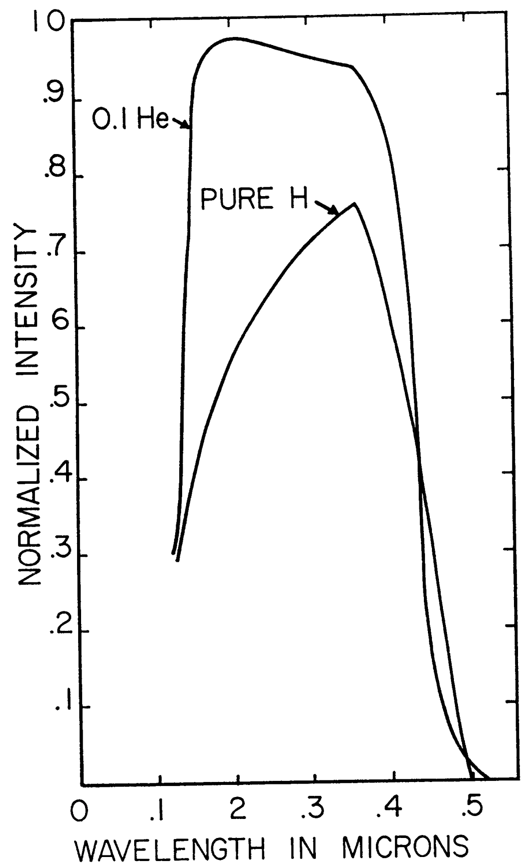

FIG. 7.—As in Fig. 4, showing the effects of helium on the degree of ionization of hydrogen

2–8-Å measures)—up to perhaps 25 Å (where intragalactic absorption is perhaps not prohibitively strong). If this were done with no positive results, *then no permissible case could be found.*

The addition of helium raises the hydrogen ionization to an extent that gives slightly more leeway in the choice of heating parameters, but the intense He^+ Ly-α line probably could be detected in OAO night-sky experiments. [Furthermore (Sargent and Searle 1966), evidence is now accumulating that the initial helium abundance was very low.]

We now wish to discuss in a semi-quantitative way the following additional points: (*a*) How critical is the nature of the heat-input model to the above results? (*b*) How are these conclusions altered if the gas is clumpy, not smooth? (*c*) If we significantly lower the postulated density by a factor of, say, 10, how much more relaxed do the restrictions on the thermal histories become? (*d*) Recently (Koehler 1965) it has been reported that 21-cm absorption in Fornax A indicates the existence of an intergalactic component of neutral hydrogen, with $n_H \sim 10^{-7}$. Are any of the above models compatible with this?

A grid of integrations using the case *b* heating assumption was performed covering a number of different values of the peak heating rate as well as the epoch at which this peak occurs. As mentioned above, unless the thermal barrier is surmounted during the rising part of the heat input expression, it will never be surmounted. If the barrier is surmounted, however, the resulting rapid increase in temperature (and corresponding drop in the emission rate) causes the radiation loss to be unimportant compared to the heat input, even during the declining part of the heat input expression. Consequently the rest of the path in the log n–log T plane is governed by the competition between the declining heating and the cooling due to expansion and the paths thus consist of a very rapid initial rise followed by a slow rise to a maximum temperature with a gentle decrease thereafter. On the basis of this grid of integrations, an estimate of the minimum radiation in the X-ray region for cases satisfying the Ly-α requirement was made, and the results are not very different from the case *a* heating model: J_ν at 10 Å must be at least about 5×10^{-28} cgs units, while the value at 25 Å is about 4×10^{-25}. Once again therefore, present X-ray results cannot exclude the existence of a heating path leading to a hot gas with $n_e \sim 2 \times 10^{-5}/\text{cm}^3$, though extension of equally sensitive results to the 25 Å region could. On the other hand, just as in the case *a* models, the range of parameters which do not give inadmissibly large amounts of 5–10-Å radiation is so limited as to be somewhat implausible.

Thus, since the two heating cases mimic qualitatively the range of heating mechanisms likely to occur, we conclude that a smooth, general intergalactic gas of density $n \sim 2 \times 10^{-5}/\text{cm}^3$—whether hot or cold—is rather unlikely.

There is at the moment no sound, purely theoretical, reason for supposing that any uncondensed gas would be smoothly distributed. On the contrary, the well-known discrepancy between the masses of galaxies in clusters determined by means of the virial theorem and those determined by other means has led to widespread speculation that considerable uncondensed gas exists in clusters of galaxies. This possibility, together with the very low or non-existent amounts of neutral hydrogen present led Woolf (1967) to see what observational limits could be placed on a hot, ionized constituent. He concluded that sufficient ionized gas to stabilize the clusters cannot yet be excluded by the observations (but see Koehler 1965). Accordingly, investigations were made to see what modification in the temperature-density path and in the X-ray emission would result if intergalactic material existed in clouds. Two cases were considered. In one case, the ratio of cloud density to mean density was considered fixed, the cloud thus expanding with the universe. In the second case, the cloud was supposed to be gravitationally stabilized following its formation. Since the general conclusions reached in the first case do not differ markedly from those in the second case, and since the second case seems somewhat more probable, only it will be discussed.

The existence of the clumps has two distinct effects on the resulting radiation field. First, if we denote by α the ratio of gas density to mean density at any instant, the emission coefficient *within the cloud* is increased by a factor of α^2, but because the clouds occupy only $1/\alpha$ of the volume, the effective emission coefficient is increased by α. Second, there is no longer any cooling due to expansion. Thus, after the rapid rise to high temperature, the only cooling is due to radiation, which even at $n = 10^{-3}$ particles/cm^3 is only moderate during the rest of the expansion at temperatures of a few million degrees. The resulting slower temperature drop compared to the adiabatic cooling critically affects the resulting emission, especially at 5–10 Å. The result of these two effects based upon the case *a* heating model is that stable clouds with a density of 10^{-3} particles/cm^3 and with a mean density of 2×10^{-5} particles which might have been formed during the heating epoch at a mean density of $n \sim 10^{-3}/\text{cm}^3$ can be excluded.

The results obtained above have referred entirely to a density of about 2×10^{-5} particles/cm^3, which is predicted for cosmological models (with zero cosmical constant) having $H_0 = 100$ km/sec/mpc and $q_0 = +1$. At the present time, a q_0 of, say, $+0.1$ implying a mean density of around 10^{-6} particles/cm^3 is by no means observationally excluded, but in such a case one would have to give up an oscillating universe, which has a certain appeal. Essentially the same two reasons that make a clumpy distribution of material easier to detect than a smooth distribution make a final density of, say, 10^{-6} particles/cm^3 very much more difficult to exclude than the $n = 10^{-5}$ case. In addition to the decreased emission coefficient due to the lower density permitted at high ionization, the maximum temperatures reached can be substantially lower in the low density case, resulting in very much lower emission shortward of 25 Å. This comes about because in the low-density case the heating necessary to surmount the thermal barrier at a Z of around 2 can be very much less (by roughly a factor of 10) with the result that the cooling effects of the expansion are much more important. Values of J_ν of around 3×10^{-29} can be achieved even at 25 Å with heating parameters which satisfy the Lyman-α requirement. Accordingly, the arguments of this paper coupled with observations obtainable in the reasonably near future probably cannot be extended to exclude mean densities very much lower than those required to give an oscillating universe.

After most of this work was completed, the author's attention was called to the work of Koehler (1965). Koehler reports evidence for a small amount of neutral hydrogen in the Virgo Cluster, as well as a weak general absorption feature in Fornax, interpreted as intergalactic neutral hydrogen, although the weakness of the feature is such that perhaps some caution should be used in its interpretation. A discussion of the spin temperature leads Koehler to an estimate of $n_{\rm H} \sim 10^{-7}$, with a limit of $n_e \sim 10^{-5}$, and, in fact, on the basis of the Virgo Cluster measures, he goes on to argue against n_e any higher than $\sim 10^{-6}$. (Some of the quantitative arguments concerning the spin temperature should be revised to take account of the influence of the 3° K radiation field at 21 cm.) However, the main point to be discussed is whether early heating could have ionized the gas to the extent that it would have been undetectable in the QSO, with subsequent recombination to the extent demanded by Koehler's observations. Such a possibility is in fact discussed very briefly by Gould and Ramsay (1966), quoting a suggestion by G. Field.

Granting the validity of the above heating models, *such a possibility is completely excluded by the calculations described above.* The final density of neutral hydrogen for the models discussed here is typically about 10^{-11}/cm^3. Very low final *total* gas densities make the situation even more difficult. Aside from adopting the local hypothesis for the QSO (in which case no heating of the gas is required at all) there seems at first sight to be at least one possible way of resolving this discrepancy. As discussed below, there is ample time for the gas to recombine provided it is not kept ionized by some mechanism. However, in case 5 of Figure 2 it is seen that the final temperature is still about 3×10^5, and a glance at Figure 1 shows that this is still so hot that collisions alone keep the gas highly ionized. The referee has suggested that contamination of the intergalactic material with a small amount of metals from, say, a density of 10^{-4} to the present epoch might sufficiently enhance the cooling to such an extent that this difficulty could be avoided. However, it is estimated that one would need to enhance the cooling rate by a factor of about 20, and this would demand a degree of contamination such that the metal/hydrogen ratio in intergalactic space be comparable to what it is in our Galaxy. Moreover, at the present epoch the radiation generated by the hot gas for case 5 accounts for about as much ionization as do the collisions, so that the injection of the impurities would have to seriously modify the thermal histories very soon after the thermal barrier was surmounted. Thus, a second possibility would be to appeal to an entirely different mechanism to ionize the gas, which does not at the same time heat the gas to very high temperatures. An intense radiation field with only a moderately high color temperature arising perhaps from high B-star formation rates in the very early stages of galaxy formation would not be unexpected. These first generation stars could completely ionize both the inter- and intragalactic hydrogen. Fairly shortly following their deaths, any residual Lyman continuum radiation would be redshifted beyond 912 Å and the gas would recombine. A computation of the recombination of initially highly ionized hydrogen was made under the following assumptions: (*a*) the initial temperature of the fully ionized gas was 10000° K and dropped adiabatically thereafter; (*b*) the initial gas density corresponded to that expected at the epoch with a Z of 2 (and one can therefore specify instead the final gas density as a parameter); (*c*) only recombinations to excited states were considered. The result is that a neutral hydrogen density of $n_{\mathrm{H}} > 1 \times 10^{-7}$ hydrogen atom/cm^3 is possible for a final *total* gas density of greater than 5×10^{-7} particle/cm^3. In fact for a final density of 10^{-5}, the recombination is nearly complete. This in itself raises a problem, since in this case it is difficult to see how to prevent *too much* neutral hydrogen from occurring, although perhaps extremely hot stars would have so much radiation in the extreme ultraviolet that large redshifts and hence long times would be required before the gas could recombine—long after the deaths of the ionizing stars.

An additional difficulty arises which is illustrated by the following specific numerical example. Assume that all the ultraviolet radiation is supplied by B0 stars of mass $M/M_{\odot} = 17$, $R/R_{\odot} = 7.6$, radiating like black bodies at $T_{\mathrm{eff}} = 21000°$ K. In order to produce sufficient ionization at a redshift of 2.5 to make the Ly-α absorption negligible (i.e., $n_{\mathrm{H}}/n_e \sim 10^{-7}$ at a density of 10^{-3}), it is necessary to have a dilution factor of $\sim 8 \times 10^{-15}$. This in turn requires that about 10 per cent of all *galactic* matter be in the form of such stars. [In making this calculation, only galaxies within a sphere of radius R were considered, with R being given by the condition that the rest wavelength λ_0 at which $J(\lambda_0) \approx \frac{2}{3}J(912\ \text{Å})$ just be redshifted to 912 Å. This value of R is only slightly smaller than the value of R corresponding to a light-travel time equal to the stellar lifetimes.] If one considers galaxies of $M = 10^{11}\ M_{\odot}$, 10 per cent of which consists of such B0 stars, and whose dimensions are of the order of 10 kpc, then one verifies that the remaining galactic material is so highly ionized that no significant absorption occurs within the Galaxy.

Such supergalaxies would of course be extremely luminous—more luminous even than the QSO themselves. Using the expressions given by Sandage (1961), one finds that for $q_0 = +1$ the (metric) angular size is about 6″, so that they should be nearly stellar in appearance with visual magnitudes of about +17.4. There would be very large numbers of them—about 2.1×10^3/square degree—which is about 2 orders of magnitude larger than the number of such objects observed at the galactic poles. Therefore this model is impossible. Calculations have not yet been made to see at what final gas density (if any) this difficulty disappears, or to what extent it can be avoided by going to extremely hot stars or supermassive stars, but these would seem worth carrying out.

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