Pregalactic evolution in cosmologies with cold dark matter

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Summary. We discuss the possibility that the first objects forming in a universe dominated by cold dark matter might influence subsequent structure. We argue that the first objects would be stars forming in clusters of mass $\sim 10^5 M_\odot$ which have condensed from the high-$\sigma$ overdensity peaks in the initial fluctuation spectrum. Two simple mechanisms – photoionization and supernovae – may cause feedback effects which can significantly perturb the subsequent evolution from a simple hierarchical clustering progression. We discuss the possibility that effects arising from these initial small-scale clusters might influence regions on galactic scales and suggest tentative ideas which may give clues to the biasing of galaxy formation necessary in an $\Omega=1$ universe.

1 Introduction

For a wide class of cosmogonic schemes, the linear fluctuations that survive until the recombination era have an amplitude that diminishes monotonically towards larger scales. In such schemes, the first bound systems to form could have had masses of order the post-recombination Jeans mass ($\approx 10^5 M_\odot$); larger scales building up later. However, as soon as the first systems have collapsed (maybe into massive stars) there is the probability of an energy input into the medium, which may influence subsequent stages in the cosmogonic process. Indeed, significant generation of heat or bulk motions – by, for example, photoionization or explosions – may occur before more than $\sim 10^{-5}$ of the baryons have condensed into stars: in other words even the rare $3-4\sigma$ peaks may exert important feedback effects before the more typical overdense regions have had time to collapse.

Although the qualitative aspects of our discussion are relevant to any ‘bottom up’ scheme for galaxy and cluster formation, we focus on one particular model: a cosmology where cold non-baryonic matter is dynamically dominant, and where the initial fluctuations are adiabatic with a constant curvature (Harrison–Zeldovich) spectrum. The attractions of this particular model are reviewed in detail by Blumenthal et al. (1984). It has the specific virtue, for our present purposes, of specifying the shape of the fluctuation spectrum: on mass-scales $\approx 10^3 M_\odot$ dark
matter is clustered in such a way that the rms amplitude of $\delta M/M$ decreases only logarithmically with $M$; on larger scales the spectrum ‘rolls over’ towards an $M^{-2/3}$ law. The amplitude is not yet predicted theoretically, and is usually fitted by matching to the observed clustering. When this is done, we find that $\sigma$ peaks in the dark matter distribution on mass-scales $\sim 10^5 M_\odot$ would have collapsed and virialized at redshifts $(1+z)\sim 8\nu\eta$. The factor $\eta (\leq 1)$ allows for the possibility of ‘biasing’ of galaxy formation which leads to enhanced clustering (see Section 4).

The Jeans mass, $M_J$, after recombination falls well below $\sim 10^5 M_\odot$. The first bound systems with $M > M_J$ capture baryons when their expansion turns around. During the subsequent collapse and virialization of the dark matter the baryons are shock heated to the virial temperature. For the baryonic gas to continue to collapse (a prerequisite for fragmentation), the cooling time-scale must be shorter than the dynamical time-scale (otherwise the gas is unable to fragment before being incorporated via a merger into a large scale of the hierarchy). The dominant cooling mechanism, due to molecular hydrogen, is ineffective until the temperature rises above a few hundred degrees. The requirement for continuing collapse – somewhat more stringent than merely requiring $M > M_J$ – is fulfilled for dark matter masses satisfying

$$M > 9.4 \times 10^5/[\Omega_\gamma^{25/24}h_5^{9/5}(1+z)^{11/4}x^{5/8}] M_\odot,$$

(1)

where $h_5$ is the Hubble constant in units of $50\,\text{km}\,\text{s}^{-1}\,\text{Mpc}^{-1}$ and $x$ is the fractional abundance of electrons after recombination; $x = 1.6 \times 10^{-5}/\Omega_\gamma h_5$. This criterion is the same as one given by Bond & Szalay (1983) corrected by Blumenthal (1983, private communication), and discussed in Blumenthal et al. (1984). Thus for $\Omega_\gamma = 0.05$ and $h_5 = 1$ the first mass able to cool at redshifts of about 15–30 will be $\sim 10^5 M_\odot$; (see Fig. 1). Throughout the rest of this paper, we take $H_0 = 50\,\text{km}\,\text{s}^{-1}\,\text{Mpc}^{-1}$ and assume that the total density parameter is $\Omega = 1$.

Once the dark matter has virialized the baryons in systems satisfying (1) collapse dissipatively, forming a self-gravitating cloud in the centre of a ‘minihalo’ of dark matter, and then probably fragment into stars. These stars might have several effects on their environment, depending essentially on their mass. Three general possibilities may be distinguished.

(i) **Low mass stars** – these do not produce ionizing radiation efficiently and cannot affect significant amounts of the surrounding gas.

(ii) **Massive stars** – these have a high luminosity, predominantly in the UV, and thus may ionize large volumes. Over its main-sequence lifetime, such a star produces $\sim 7\,\text{MeV}$ of UV for each of its constituent baryons. If recombinations were negligible it could consequently ionize $\sim 2.5 \times 10^5$ times its own mass in the form of diffuse gas (Couchman 1985).

(iii) Massive stars below a critical mass $\sim 200 M_\odot$ may explode, ejecting heavy elements. A cluster of such stars would generate strong shock waves in the surrounding gas. These shock waves may ‘sweep up’ large amounts of material into thin shells.

We cannot confidently predict the mass function of the first stars; it has often been argued that those forming from the primeval gas would be massive, or even very massive objects (VMOs) of masses $10^3$–$10^5 M_\odot$ (Carr, Bond & Arnett 1984). Conversely, Kashlinsky & Rees (1983) suggest that the first stars could form from gas that had collapsed to densities $\approx 10^{10}/\text{cm}^3$ yielding a low Jeans mass even in the absence of heavy metal coolants. Nor do we know the likely ‘efficiency’ – that is to say, the fraction of the baryons in a system of mass given by (1) which will turn into stars rather than being expelled from the rather shallow potential wells of the first bound systems. This factor, which we call $\psi$, would be a function of the IMF and the depth of the potential well.

Dekel & Silk (1986) suggest that only a small fraction of the gas trapped in a potential well will be able to form into stars; the remainder being blown out by the supernovae explosions of the first stars. This would lead to the parameter $\psi$ having a very low value, if the IMF resembles that in present-day galaxies. An important factor, however, which may influence $\psi$ is the time that

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elapses between the first stars forming and the time at which a sufficient number of explosions occur to remove the remaining gas from the potential well. This may allow significantly more gas to condense into stars, thus increasing \( \psi \).

We consider separately the two main feedback effects: photoionization, which could arise from any population of objects more massive than \(-10 \, M_\odot\) (Section 2); and blast waves, which could come from multiple supernovae (the endpoint of massive stars but not of VMOs) or from strong stellar winds (Section 3). Some cosmogenic consequences of these effects are considered in Section 4.

2 Ionizing radiation

Massive stars \((M \approx 100 \, M_\odot)\) radiate at close to the Eddington limit, and have a main-sequence lifetime of \(\sim 10^{14} \, \text{s}\). This is very much shorter than the Hubble time for \(1 + z < 100\) and less than the recombination time in the gas for \(1 + z < 100 (T/10^4)/(\Omega_\gamma/0.05)^{1/2} h_0^2\). Thus at redshifts of \(\leq 30-40\) it is sufficient to consider that ionized regions are formed instantaneously around each star at some redshift, subsequently cooling and recombining. The dominant cooling process at redshifts greater than about 5 is Compton cooling; \(t_{\text{cc}} = 7.1 \times 10^{19}/(1 + z)^4 \, \text{s}\). This cooling time is longer than the lifetime of massive stars for \(1 + z \leq 30\).

If the spectrum of a star of \(\gtrsim 100 \, M_\odot\) can be approximated by a \(10^5 \, \text{K}\) blackbody the temperature attained in the ionized region, allowing for the presence of He, is roughly \(6 \times 10^4 \, \text{K}\). The zone lasts for a time of order the Compton cooling time, since \(t_{\text{rec}} \approx t_{\text{cc}}\) at these initial temperatures. The gas recombines after cooling (via the Compton effect) to \(\approx 10^3 \, \text{K}\), but will still be on a higher adiabat than before the reheating.

The ionization front, which never comes into equilibrium since \(t_{\text{rec}} \sim 10^{14} \, \text{s}\), advances very much faster than the sound speed and thus no substantial movement of the gas takes place. (Normally some concentration of matter occurs near the ionization front in a steady Strömgren sphere.)

Fig. 1 shows how a rise in the Jeans mass could lead to a pause in the cosmogenic process until larger scale perturbations collapsed (see caption for further details). Smaller stars with masses anywhere in the range \(10-100 \, M_\odot\) produce qualitatively similar effects. A star with mass \(M = M_2 \approx 100 \, M_\odot\) in this range has a luminosity \(L = L_{\text{ED}} M_2^2\), a lifetime \(t_{\text{MS}} = 10^{14}/M_2^{1/2} \, \text{s}\) and an effective blackbody temperature \(T = 10^5 M_2^{0.3} \, \text{K}\) (Ezer & Cameron 1971). A star of mass \(M_2\) produces \(\sim M_2^2\) as many ionizing photons over its lifetime as a \(100 \, M_\odot\) star. For \(t_{\text{MS}} < t_{\text{rec}}\) stars of mass \(M_2\) in a cluster of fixed total mass would thus ionize a mass of gas \(M_2^{1/2}\) times that ionized if the stars had been of mass \(\approx 100 \, M_\odot\).

It is worth noting that dust, if present in sufficient quantities, could thermalize much of the UV radiation from massive stars, thereby preventing the photoionization of the intergalactic gas. At the epoch when the first stars form, however, the relative spatial influence of supernovae explosions to that of photoionizing radiation is very small; thus, at these redshifts, the environment around a radiating massive star will not have been contaminated by the dust from an explosion (see Section 4).

3 Explosions

The other important mechanism by which massive stars can affect their environment is by supernova explosions. Shocks from these explosions can sweep up large amounts of material into shells which may subsequently fragment. The explosions can also enrich the shocked gas with metals thereby affecting its physical properties.

Three distinct phases in the expansion of a supernova shell may be identified: initially the ejecta expand freely until a mass of the order of the ejecta mass has been shocked. Subsequently the
shock expands according to the usual energy-conserving ‘Sedov’ solution until cooling becomes important for the swept up gas. The shock is then no longer driven by the heated gas and becomes momentum conserving, satisfying the familiar ‘Oort snowplough’ solution. The shock radius is then given by

$$r = r_c \begin{cases} \left( \frac{\Delta}{\Delta_c} \right)^{2/5} & \Delta \leq \Delta_c \\
\left( \frac{8}{5} \frac{\Delta}{\Delta_c} - \frac{3}{5} \right)^{1/4} & \Delta \geq \Delta_c \end{cases}$$

where $r_c = 1.15 (E_{t_0}/\Omega_0) \Delta_c^{2/5} \Delta_c^{2/3}$ is the shock radius after a cooling time, $\Delta = (t - t_0)/t_0$, $\Delta_c = t_{cc}/t_0$, $t_0 = 4 \times 10^{17}/(1 + z)^{3/2} h_{50}$ is the Hubble age when the explosion occurs and $\rho_c = (6 \pi G \rho)^{-1}$ is the critical density appropriate for an $\Omega = 1$ universe. $E$ is the energy released in the explosion; it has been estimated to be roughly $10^{-4}$ of the rest mass energy of the star (Bookbinder et al. 1980). The formula is valid for $\Delta = 1$, since the background expansion has been neglected. For $\Delta \geq \Delta_c$ the solution is obtained simply by matching the radius and velocity of the shock at $\Delta = \Delta_c$ to the snowplough solution. Equation (2) supposes that a strong shock is formed; this requires that the Mach number, $\mathcal{M}$, be very much greater than unity. I.e. $u_{\text{shock}} \gg c_s$, where $c_s = (\gamma kT/\mu m_p)^{1/2}$ is the external sound speed.
The Mach number is given by

\[ \mathcal{M} = \mathcal{M}_C \left\{ \begin{array}{ll}
\left( \frac{\Delta_c}{\Delta} \right)^{3/5} & \Delta \leq \Delta_c \\
\left( \frac{8}{5} \frac{\Delta}{\Delta_c} - \frac{3}{2} \right)^{-3/4} & \Delta > \Delta_c
\end{array} \right. \]

where

\[ \mathcal{M}_C = 1.6 \left( \frac{E_\odot}{\Omega_g h_{50}^2} \right)^{1/5} \left( \frac{(1+z)^{13/10}}{(1+z)^{23/8}} \right)^{1/2} \]

\[ T_{\text{rad}} = 2.7(1+z) \]

is the temperature of the microwave background. \( E_\odot \) is the energy in units of \( M_\odot c^2 \).

A strong shock persists until \( \mathcal{M} \) drops to \( \sim 1 \). The shock lasts for a Hubble time if

\[ E_\odot > 2.2 \times 10^{-9} \frac{\Omega_g}{h_{50}^{7/4}} \left( \frac{T_{\text{gas}}}{M_{\text{gas}} T_{\text{rad}}} \right)^{5/2} (1+z)^{23/8}, \quad z > z_c. \]

This condition is easily satisfied for the energies released by massive stars at redshifts less than 100 unless the ambient gas temperature is very much higher than the microwave background temperature.

It is important to consider the cumulative effect of the shocks from several exploding stars within a cluster. Supposing that the first objects form from a \( 10^5 M_\odot \) region of dark matter then the radius of this region when virialized is roughly

\[ r_{\text{vir}} = 1.3 \left( \frac{M_{\text{DM}}}{10^5 M_\odot} \right) \left( \frac{h_{50}^2}{1+z_c} \right)^{1/5} \text{kpc}, \]

where \( z_c \) is the redshift at which to system collapses. The shock radius when a cooling time, \( t_{\text{cc}} \), has elapsed is

\[ r_c = 690 \left( \frac{M_\odot}{100 M_\odot} \right) \left( \frac{\Omega_g h_{50}^2}{1+z} \right)^{1/5} \text{kpc}. \]

Thus even for a \( 10 M_\odot \) star \( r_c > r_{\text{vir}} \) for \( 1+z < 100 (\Omega_g = 0.05, h_{50} = 1) \). Suppose further that the stars explode within a few million years of each other. Compton cooling is then negligible during this period for \( 1+z \leq 40 \). The first objects in the CDM picture form at redshifts of \( \sim 20-30 \). To a good approximation, therefore, all explosions within a cluster can be treated as one large single explosion of the total mass: since the first explosion carries the shock front beyond the limits of the cluster, subsequent shocks quickly catching up with this front in the heated and rarefied medium before any cooling occurs.

As the shock propagates through the medium it sweeps up material into a shell moving behind the shock front. During the Sedov phase the temperature of the material is high and the shell is relatively thick. For this reason fragmentation is normally prevented during the energy conserving phase. If the shell cools, however, it may become very thin and be able to fragment into gravitationally bound objects.

It is difficult to calculate the fragmentation of the shell in detail; a simple energy condition to determine whether ‘discs’ cut from the shell are gravitationally bound or not can, however, be easily derived (Ostriker & Cowie 1981).

A disc of radius \( a \), cut from a shell of radius \( R \) with Mach number, \( \mathcal{M}_s \), is bound if

\[ a^2 \geq 16 \left( \frac{8 \Omega_g}{3\pi} \right) \mathcal{M}_s \Delta^2 q + \frac{\gamma_6}{q} < 0, \]

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where \( q = R/a \) (here \( \mathcal{M}_s \) is the ratio of the shock velocity to the sound speed in the cooled shell). For real values of \( a \), the discriminant of the quadratic must be positive. Thus, for fragmentation after \( t_{cc} \), the Mach number must satisfy

\[
\mathcal{M}_s = \frac{1.62}{\Omega_g \Delta^2}.
\]  

(8)

Fragmentation could occur before a Hubble time has elapsed only if the Mach number satisfies \( \mathcal{M}_s > 1.62/\Omega_g \). Equivalently a temperature satisfying

\[
T_{\text{shell}} < 3 \times 10^9 \frac{E_{\odot}^{2/5} \mu \Omega_g^{8/5} h_{50}^{10}}{(1+z)^{3/20}} \Delta^{5/2}
\]  

(9)

is needed for fragmentation for \( \Delta_c \leq \Delta \leq 1 \). For \( h_{50} = 1 \Delta_g = 0.05 \) and \( E = 0.5 E_{\odot} \) (equivalent to the baryons in a \( 10^5 M_{\odot} \) dark matter clump exploding with an efficiency of \( 10^{-4} \)), \( T_{\text{shell}} \approx 12 \text{ K} \). Since it seems unlikely that \( T_{\text{shell}} \) can become as low as this, it would be impossible for a \( 10^5 M_{\odot} \) condensation to trigger a detonation wave, as discussed in various related contexts by Ostriker \\& Cowie (1981), Ikeuchi (1981), Carr \\& Rees (1984) and Bertshinger (1985).

Clearly the post-shock temperature is important in determining if fragmentation could occur. The cooling is primarily due to molecular hydrogen. If \( H_2 \) has a low abundance, arising (via \( H^- \)) from the free electrons remaining after recombination (Lepp \\& Shull 1984), then the gas cannot cool much below \( \sim 1000 \text{ K} \). Maclow \\& Shull (1986) have recently shown that a higher \( H_2 \) abundance (up to \( \sim 10^{-3} \)) can arise when material cools indirectly after being shock heated to \( > 10^4 \text{ K} \) (hot enough to ionize hydrogen collisionally), and that the post-shock temperature can then drop to \( \sim 300 \text{ K} \). Taking this as a lower limit confirms that the fragmentation criterion fails by a wide margin.

![Figure 2](https://academic.oup.com/mnras/article-abstract/221/1/53/1028302/22110531028002)  

**Figure 2.** The mass swept up by a shock after a Hubble time, at a particular redshift, has elapsed. The masses are expressed in terms of the equivalent dark matter masses thus removing the dependence on \( \Omega_g \). The curves are plotted for the baryons in a dark matter clump of mass \( M = 10^5 M_{\odot} \) exploding with efficiency \( \epsilon \sim 10^{-4} \). \( M_H \) is the mass swept up after a cooling time and \( M_H \) that swept up after a Hubble time. Also shown is the mass affected by the ionizing radiation from the same cluster of baryons in two cases depending on whether it is composed of 100 \( M_{\odot} \) or 10 \( M_{\odot} \) stars. The efficiency per unit mass for supernovae explosions is assumed to be constant for stars in the mass range 10–100 \( M_{\odot} \). The shock curves scale as \((M/10^5)^{3/5}\), the ionization curves as \((M/10^5)\).
The mass swept up into the shell during the expansion of the shock is

$$M = M_c \left\{ \frac{\Delta}{\Delta_c} \right\}^{6/5} \Delta \leq \Delta_c$$

$$M = \left( \frac{8\Delta}{3\Delta_c} \right)^{3/4} \Delta \geq \Delta_c$$

(10)

where $M_c$ is the baryon mass shocked after a cooling time has elapsed.

$$M_c = 1.4 \times 10^{12} (\Omega_g h_0^2)^{2/5} E_5^{3/5} / (1+z)^{18/5} M_\odot.$$  

(11)

This formula is strictly valid only for $\Delta < 1$, since it takes no account of the cosmological expansion of the background material. For $\Delta_c < 1$ it is a simple matter to extend the momentum conserving snowplough into the regime for which $\Delta \geq 1$. In a low gas density universe (with $\Omega_{tot} = 1$ none of the less) the shock velocity so rapidly approaches the Hubble flow for $\Delta > 1$ that it is effectively a constant mass that is enclosed within the shock front at all later times. Thus the total mass shocked when a Hubble time has elapsed, roughly approximated by putting $\Delta = 1$ in the above formula, is close to the mass shocked when several Hubble times have elapsed. Indeed, for $\Omega_g = 0.05$ the mass affected were the shock to propagate to the present, for $z < 100$, is less than twice the value given by putting $\Delta = 1$ in the formula above. The mass swept up by the shock after a cooling time and that swept up after a Hubble time are shown in Fig. 2 for $z < 100$.

4 Implications for ‘Cold Dark Matter’ cosmology

4.1 Density of Initial Objects

The first bound objects may naturally be supposed to form from the highest peaks of the random density fluctuations present in the early universe. If these fluctuations can be represented by three-dimensional Gaussian noise a simple expression for the average density of maxima with height $> n \sigma$, in the high $\sigma$ limit, is given by Doroshkevich (1970):

$$n_{pk}(> n) \approx \left( \frac{j/3}{(2\pi)^2} \right)^{3/2} v^2 \exp \left( -v^2 / 2 \right)$$  

(12)

where $j$ is the second spectral moment of the power spectrum of the initial fluctuations. [This topic has recently been rigorously discussed by Bardeen et al. (1985); who confirm that the simple expression (12) is sufficiently accurate for present purposes.] Writing $v$ in terms of the redshift at which a particular overdensity condenses gives, for the number of condensed peaks at redshift $z$ of mass $M$,

$$n_{pk}(z) \approx 4.9 \times 10^{10} \eta^3 \left( \frac{1+z}{\sigma^2} \right) \exp \left[ -\left( 1+z \right)^2 / (2\sigma^2) \right] / (\text{Mpc})^3$$

(13)

where $\sigma^2 = 9\eta^2 \log (1.3 \times 10^{11} / M)$ and $\eta = 1$ for the usual normalization of $\delta M / M = 1$ at 16 Mpc ($h_0 = 1$). The formula is obtained by filtering the spectrum at some radius $R_t$ corresponding to a mass $M = 4/3\pi R_t^3 \rho_c$. The expression for $\sigma$ is valid for $M \approx 10^{10} M_\odot$.

The rough formula given above, for the density of condensed regions, allows simple estimates of the redshift at which sufficient numbers of sources exist to influence the subsequent pregalactic evolution. Suppose that some physical process filters the baryon power spectrum on a scale $R_t$. Here it is supposed that it is roughly $10^9 M_\odot$ due to the effects of molecular hydrogen cooling [equation (1)].
If a fraction $\psi$ of the baryons in a peak of mass $M$ condense into massive stars which burn for $10^8$ s and form ionized zones which last for a time $t_{\text{rec}}$ then the fraction, $F$, of the universe which will be ionized at a given redshift is given, for $h_{50} = 1$, by

$$F = 2.3 \times 10^7 \psi^3 \frac{(1+z)^{3/2}}{\sigma^2} \exp \left[-(1+z)^2/(2\sigma^2)\right].$$

(14)

For $M = 10^5 M_\odot$ and $\psi = 1$ this is unity for $1+z = 30\eta$ corresponding to $4\sigma$. Bond & Efstathiou (1984) and Couchman (1985) have given similar estimates of the mass fraction radiating at the Eddington Limit necessary to fully ionize the universe at $z<100$. Equation (14) is, in fact, relatively insensitive to the masses of the individual stars, provided that these have mass $\geq 10 M_\odot$. As discussed earlier, the mass ionized (for $t_{\text{MS}} < t_{\text{rec}}$) by a fixed mass of gas forming into lower mass stars $\leq 100 M_\odot$ is reduced only by a factor $\sim (M/100 M_\odot)^{1/2}$. We expect the first stars to form at $1+z \sim 20\sim30$ and in this case $t_{\text{MS}} < t_{\text{rec}}$ even for $10 M_\odot$ stars. For stellar masses below $10 M_\odot$ photoionization of the universe would require the condensation of a much larger fraction of the mass.

Similarly the fraction of the universe which has been shocked by a redshift $z$ is

$$F = 1.2 \times 10^8 \psi^3 \frac{(1+z)^{11/40}}{\sigma^2 M^{2/5}} \exp \left[-(1+z)^2/(2\sigma^2)\right].$$

(15)

This is unity, for $M = 10^5 M_\odot$, for $1+z = 23\eta$ corresponding to $\sim 3\sigma$.

4.2 Consequences of Photoionization

The primary effect of photoionization is to heat the baryons and thus raise the Jeans mass of the gas. A simple criterion for baryons to be prevented from settling into the dark matter potential well is that

$$\frac{3 kT}{2 \mu m_p R} \gg \frac{GM_\odot}{R}$$

(16)

where $M$ is the mass of the virialized dark matter clump and $R$ is its characteristic radius. Thus

$$M_\odot \approx 6.8 \times 10^8 \left(\frac{T}{10^4}\right)^{3/2} \left[\mu^{3/2} h_{50}(1+z_c)^{3/2}\right] M_\odot$$

(17)

where $z_c$ is the redshift at which the clump collapses (or virializes). Ionizing radiation with a $10^5$ K blackbody spectrum heats the gas to $\sim 6 \times 10^4$ K, raising the Jeans mass to $\sim 2 \times 10^8 M_\odot$ at $z \sim 30$. At large redshifts the radiating sources would be sufficiently rare that the H II regions formed around them would be isolated spheres. Once the source dies Compton cooling quickly reduces the gas temperature causing the Jeans mass to fall below $10^5 M_\odot$ for $z \geq 15$ (see Fig. 1). Since both the lifetime of the source and the Compton cooling time then are very much less than the dynamical time for collapse, the sudden formation of isolated H II regions and their rapid cooling and recombination would have a negligible effect on subsequent evolution. This conclusion fails at later times when the increasing density of sources causes the H II regions to influence one another.

A more significant and prolonged effect may occur after the ionized zones overlap and much more energy is available from the greater density of stars. Once overlap occurs further energy input will rapidly cause the medium to become highly ionized and optically thin. The reservoir of extra photons produced, in excess of the number required simply to fully ionize the medium, may be sufficient to keep the medium ionized for a substantial change in redshift. Thus a burst of star
burning at some redshift may keep the temperature raised for some time after the stars have died. This occurs because of the relatively long recombination time in the low density gas at these redshifts. For example if three times as many photons as are required to fully ionize the medium are emitted at a redshift of ~25 the temperature will remain at ~5000 K for redshifts down to ~10 (Couchman 1985). Thus the Jeans mass would be maintained at ~10^7 M_☉ during this period.

4.3 General Cosmological Implications

The most obvious and significant inference from (14) and (15) is that energy input from the first protogalactic bound systems, of characteristic masses ~10^5 M_☉, can perturb the primordial gas (raising it to a higher adiabat, inducing variations of the baryons relative to the dark matter, etc.); moreover the energy input could be so efficient that even a mass fraction of only ~10^{-5} (the ≥4σ peaks) might affect the entire baryonic content of the universe so as to 'choke off' the formation of further ~10^5 M_☉ systems by raising the Jeans mass, and maintaining it at ≥10^7 M_☉. These first bound objects (and their remnants) may therefore be very sparse – perhaps only 1/10^{10} M_☉. The idea that a small fraction of the baryons might influence subsequent evolution in this way was first discussed by Doroshkevich, Zeldovich & Novikov (1967).

The processes we have discussed in Sections 2 and 3 could obviously complicate the evolution of bound systems in the mass range 10^2–10^8 M_☉: UV or blast waves from the first ~10^5 M_☉ objects could readily suppress the formation of any further bound systems below ~10^6 M_☉. As far as galaxy formation is concerned it is necessary to consider how the effects might influence the subsequent condensation of bound systems in the galactic mass range 10^8–10^{12} M_☉.

It is interesting to consider the spatial distribution of the first bound objects which will reflect that of the high peaks in the initial density fluctuations. The significant feature of the CDM spectrum is that it has a slow fall over a wide range of mass-scales ~10^{-5}–10^{-12} M_☉. This has the result that the high σ peaks on small scales may be significantly clustered on very much larger scales (Kaiser 1984; Bardeen et al. 1985 and references cited therein).

Suppose that the first bound objects correspond to νσ peaks on a scale M_1=10^5 M_☉. The spatial distribution of these objects will be modulated by the long-wave components of the initial fluctuation spectrum: there will be more of them, per comoving volume element, in large-scale peaks (protogalaxies and protoclusters) than in the underdense regions (protovoids). Qualitatively, their concentration within a μσ peak on a scale M_2 (≫M_1) is roughly proportional to (1+ννσ_2/σ_1). In this expression σ_2/σ_1 denotes the relation rms amplitude of the initial fluctuations on scales M_2 and M_1; it is a slowly decreasing function of M_2 (see Fig. 1) being ~1/4 for M_2=10^{12} M_☉ and ~1/8 for M_2~3×10^{11} M_☉. If the first bound objects, of mass M_1, corresponded to ν=4, they would be concentrated by a factor ~2 in 1σ peaks on scales of 10^{12} M_☉, and in 2σ peaks on scales as large as 3×10^{13} M_☉.

This illustrates how pregalactic processes could in principle induce environmental variations on scales of galaxies or larger. We recall that Davis et al. (1985) showed that the data on galaxy clustering could be reconciled with an Ω = 1 cosmology if galaxies formed only at ≥2σ peaks in its mass distribution. They offered no physical model for this biasing (see however Rees 1985 and Silk 1985 for various suggestions). Our present discussion shows that even before their collapse begins the internal properties of 2σ galactic mass peaks differ from those of 1σ peaks of the same mass-scale. We cannot yet show that these differences would in themselves account for the necessary biasing, but tentatively suggest two ways in which this might come about:

(i) Photoionization (by, for instance, VMOs) could not in itself affect subsequent galaxy formation unless this were sensitive to the adiabat of the gas (or H_2 abundance) when it started collapsing. The first 10^5 M_☉ systems may, however, contain ordinary massive stars which explode.
as supernovae. If enough such systems form to produce overlapping shells [equation (15)], then the resultant heavy elements could contribute a Population II abundance. These heavy elements (unlike the UV photons from the same systems, which, after photoionization is complete, constitute an essentially uniform background) would be distributed in the same fashion as the 3σ peaks – thus the pregalactic contamination by heavy elements would be inhomogeneous on scales up to ~10^{12} M_\odot. In so far as the heavy element abundance in a protogalaxy influences its subsequent evolution this offers a possible biasing mechanism.

(ii) Although the shells from the supernovae in 10^5 M_\odot clusters do not become gravitationally unstable (and so do not initiate a detonation wave), the further density rise \sim H^2 that occurs when the two shells collide may trigger gravitational collapse. If so, rapid star formation could ensue when the shells start to overlap. A very strong galactic-scale biasing mechanism could arise if a runaway process of this kind occurred in the galactic-scale peaks with overdensities >2σ but the same process never occurred in the less overdense (more typical) peaks because photoionization quenched the condensation of 10^5 M_\odot systems there before enough formed for the shells to overlap.

The epoch when these processes occur depends on the factor η, which normalizes the amplitude in Fig. 1. This factor would, by definition, be unity if the galactic distribution traced the overall mass. If Ω=1, however, then the data can be fitted only if η=0.5, and galaxy formation is biased towards the high σ peaks of the mass distribution on scales ~10^{12} M_\odot (Davis et al. 1985). With this choice of η the feedback effects we have discussed would occur in the redshift range z~10–15, and galaxies themselves would form at z=5.

References