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Formation of the Bubbly Universe by Cumulative Explosions

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Abstract

We investigate the scale of structures of the universe which can be produced in the explosion scenario. Solving the equation of motion of an expanding shell produced by cosmic explosions based on the thin-shell approximation, we show it is possible that bubbles with the size ~ 40 Mpc are formed by successive explosions starting from the first-generation objects. However, the observed isotropy of the microwave background radiation limits the bubble size to less than ~ 30 Mpc in the simple explosion scenario. This seems to be smaller than the size of the standard bubbles. The formation of still larger bubbles may be possible, if the merging processes of bubbles and recent overlapping of the shells are considered.

Key words: Bubbly universe; Cosmic background radiation; Galaxy formation; Explosion scenario.

1. Introduction

Recent galaxy redshift surveys have revealed the three-dimensional distribution of galaxies in the universe (Einasto et al. 1980; Davis et al. 1982). In particular de Lapparent et al. (1986) showed that the universe is filled with nearly spherical voids and galaxies are distributed on the surfaces of the voids.

How was this large-scale structure of the universe formed? The explosion scenario is one of the ideas which attempt to explain it (Ikeuchi 1981; Ostriker and Cowie 1981). In this scenario a small number of seed objects are first formed and the explosive energy release from these objects produces shock waves. The shock wave sweeps up the primordial gas and a cooled shell is formed. Gravitational fragmentation of the shell results in new objects. These second-generation objects also explode and larger shells are formed and fragment into the third-generation objects. By such successive explosions, larger structures are produced and finally the presently observed bubbly structure in the universe is formed.

Carr and Ikeuchi (1985) extended the basic idea. First an object with the mass of about $10^6 M_\odot$ explodes in the Compton cooling era (redshift $z \gtrsim 10$). As a result of the detonative explosions a shell with the mass $\sim 10^{10}$–$10^{11} M_\odot$ is formed. The last explosions in the radiative cooling era ($z \lesssim 10$) result in the large-scale structure. Following the idea of Carr and Ikeuchi (1985), Bertschinger (1985) carried out a simple numerical simulation. By the use of the analytically derived expansion law of a cosmological shock wave in an expanding universe, he showed that a cosmological detonation wave can produce a large structure with the mass $\sim 10^{15} M_{\odot}$.

The explosion scenario has the following several advantages:

(1) Since galaxies are formed as a result of the fragmentation of the shocked shell, the galaxies will be distributed on the two-dimensional surface. On the other hand the space swept up by the shock wave is left as the spherical voids of galaxies. In particular, if the shells produced in various places in the universe come to overlap by the present epoch, the bubbly structure of the galactic distribution can be naturally explained. In the interior of the void, there exsist burnt-out pregalactic objects.

(2) We need not assume primordial density fluctuations on the large scale, because the first-generation objects with the relatively small mass are sufficient as the seeds of explosions.

(3) In the spectrum of a distant quasar ($z=3.8$) no absorption trough of the intervening neutral hydrogen is detected (Gunn–Peterson test). Heating and ionization of the gas by the shock waves can explain it.

However, several difficulties have been pointed out; the sizes of the observed voids are too large and considerable anisotropy and distortions of the cosmic background radiation are produced (Hogan 1984; Vishniac and Ostriker 1986; Bond et al. 1986). (Moreover, if there exists dark matter nearly uniformly in the universe, more energy is necessary in order to make big voids. This makes the situation worse.) In order to answer these problems, it is necessary to investigate the evolution of a shock wave in an expanding universe in more detail. So far the selfsimilar solutions have been derived by Ikeuchi et al. (1983), Bertschinger (1983), and Ostriker and McKee (1988). In addition the spherically symmetric numerical calculations which included the radiative cooling were performed (Ikeuchi et al. 1983; Vishniac et al. 1985). However, they did not try to connect the result of their calculations with the real structure of the universe and did not investigate the evolution of the shock wave considering successive explosions of the fragments of the shell.

In this paper we study the evolution of cosmic detonation waves numerically and give constraints on the size of the explosively produced structure. In section 2 we present the model and the method of calculation for the evolution of the shock wave in an expanding universe. The results of calculations are shown in section 3. In section 4 the resultant anisotropy of the microwave background radiation by explosions are calculated and the constraints on the bubble size are obtained. A summary and discussion of the results are presented in section 5.

2. Model and Method of Numerical Calculations

2.1. *Model*

We take a baryon-dominated universe as a cosmological model. That is to say, the density parameter of a universe Ω equals Ω_{b}. From the cosmic virial theorem and the infall of the local group into the Virgo cluster the present density parameter has been estimated as $\Omega_0=0.1$–0.4 (Davis and Peebles 1983a, b). On the other hand the calculations of primordial nucleosynthesis predict $\Omega_{b,0}=0.03$–0.1 (e.g., Yang et al. 1984). So we consider a universe with $\Omega_0=\Omega_{b,0}=0.1$.

We assume that a seed object with mass M_i releases energy explosively at the redshift z_i with efficiency ε, which is the fraction of the rest mass of the objects converted to energy. If the cause of the energy release is nuclear reactions in the seed object, $\varepsilon \lesssim 10^{-3}$ (Carr et al. 1984), and the ejectable kinetic energy is given by

$$E_i=1.8\times10^{61}(\varepsilon/10^{-4})(M_i/10^{11}\,M_\odot)\ \mathrm{erg}. \tag{1}$$

The shock wave generated by the explosion propagates in cosmic space, sweeps up the primordial gas, and forms a spherically expanding shell.

2.2. *Basic Equations*

Here we take the thin-shell approximation that the thickness of the shell is much smaller than the radius of the shell and the shell contains all the swept-up matter. Then the equation of motion of the shell is (Ostriker and McKee 1988)

$$\frac{d}{dt}(M_s V_s)=4\pi R_s^2[P+\rho_g V_h(V_s-V_h)]-\frac{GM_s^2}{2R_s^2}\,, \tag{2}$$

where R_s and V_s, respectively, denote the radius and the velocity of the shell, ρ_g is the average gas density of the universe, $M_s=(4\pi/3)R_s^3\rho_g$ is the mass of the matter swept up by the shell, $V_h=H_0(1+z)(1+\Omega_0 z)^{1/2}R_s$ is the unperturbed Hubble expansion velocity, and P is the average pressure inside the shock wave, which is given by

$$P=\frac{2}{3}\frac{E_{th}}{(4\pi/3)R_s^3}\,,$$

where E_{th} is the total thermal energy of the blast wave. We assume that the pressure outside the shock wave can be neglected compared with P. The present Hubble constant is written as $H_0=50h_{50}\mathrm{km\ s^{-1}Mpc^{-1}}$.

On the other hand the equation describing the change of the total energy of the shock wave $E_{tot}=(1/2)M_sV_s^2+E_{th}-(1/2)GM_s^2/R_s$ is

$$\frac{d}{dt}E_{tot}=-\int_V \Lambda dV=\dot{E}_{loss}\,, \tag{3}$$

where Λ is the cooling rate per unit volume. The gas is assumed to have the primordial abundance with the hydrogen to helium ratio 9: 1 by number. We consider the inverse Compton cooling of the microwave background photons and the radiative cooling as the cooling mechanism. The radiative cooling processes contain thermal bremsstrahlungs, electron collisional excitations, radiative recombinations, and dielectronic recombinations. Details of the cooling processes are given by Umemura and Ikeuchi (1984). In the calculation of the cooling rate the density in the shell is supposed to be enhanced by D times the average gas density in the universe $\rho_s=D\rho_g$. We choose $D=10$ for the subsequent discussion (Ikeuchi et al. 1983). The total cooling rate hardly depends upon the value of D in the era $z\gtrsim10$ when the Compton cooling is dominant.

With equations (2) and (3) the change in the thermal energy is described by

$$\frac{d}{dt}E_{\rm th}=4\pi R_{\rm s}^{2}\left(\frac{1}{2}\right)\rho_{\rm g}(V_{\rm s}-V_{\rm h})^{2}(V_{\rm s}-V_{\rm h})-4\pi R_{\rm s}^{2}PV_{\rm s}-\dot{E}_{\rm loss}\,. \qquad (4)$$

The first term of the right-hand side of this equation represents the effect of the dissipation of the kinetic energy into the thermal energy, the second term the work by pressure, and the last term the energy loss by cooling.

From the calculated thermal energy in the shock wave by equation (4), we can calculate the temperature of the shell by

$$T_{\rm s}=\frac{2\mu}{3k}\frac{E_{\rm th}}{M_{\rm s}}\,,$$

where μ is the mean particle mass and k is the Boltzmann constant.

We assume the simple relation between the time t and the redshift z as

$$\frac{da/dt}{a}=H=H_0(1+z)(1+\Omega_0 z)^{1/2}\,, \qquad (5)$$

where the scale factor a is related to z by

$$a/a_0=1/(1+z)\,,$$

using the value of the scale factor a_0 at $z=0$.

The above three differential equations (2), (4), and (5) are together numerically integrated.

2.3. *Gravitational Instability of the Shell*

Here we consider the gravitational instability and fragmentation of a spherically expanding shell. Ostriker and Cowie (1981) obtained a fragmentation condition using a simple energy consideration. Vishniac (1983) investigated the stability of a shell using linear analysis and obtained a qualitatively similar result. Therefore we adopt a generalized version of the Ostriker–Cowie (1981) condition, which is given by

$$Y=\left[\Omega_{\rm b,0}H_0^{2}(1+z)^{3}\frac{R_{\rm s}^{2}}{C_{\rm s}V_{\rm s}}\right]\bigg/\left[3\pi\left(\frac{3}{2}\right)^{1/2}\right]>1\,, \qquad (6)$$

where $C_{\rm s}=(kT_{\rm s}/\mu)^{1/2}$ is the sound velocity at the temperature of the shell $T_{\rm s}$. When this condition is satisfied, the shell is assumed to fragment instantly.

2.4. *Successive Explosions*

Objects fragmented out from the shell may explode again. In this case explosions are expected to occur on the shell. We treat this process as follows:

First, suppose the shell satisfies the fragmentation condition (6) at $z=z_{\rm frag}$. At this moment the shell with the mass $M_{\rm s}$ is assumed to fragment and promptly explodes with the efficiency ε again, because the lifetime of newly born stars will be small enough. The energy ejected at this time is taken to be

$$\varDelta E = \varepsilon \varDelta M_s c^2 ,$$

where $\varDelta M_s$ is the mass of the matter swept up by the shell after the last explosion till the fragmentation. $\varDelta E$ is added as an increase in the thermal energy of the shock wave. Thus the pressure P increases and accelerates the shell. The initial conditions of this motion are given by the values R_s and V_s at $z=z_{\rm frag}$. This shell may cool and fragment again as it expands into the cosmic space, and may give rise to a much larger structure.

3. Results

Here we concentrate on the size of the structure produced by each explosion. The overlapping of neighbouring shells is not considered.

Figure 1a shows an evolution of a typical shock wave, which is produced by an explosion with the total energy $E=10^{62}$ erg, which corrresponds to $M_i=5.6\times10^{11}\,M_\odot$ according to equation (1), at $z_i=20$ in a universe with $\Omega_0=\Omega_{b,0}=0.1$. This figure shows the change in the radius of the shock wave R_s, the peculiar velocity $V_p=V_s-HR_s$, the mass of the matter swept up by the shock wave M_s, and the temperature T_s. In this calculation we do not consider the fragmentation of the shell, i.e., this is only a simple explosion. From these results we see that a single explosion can produce a shell radius as large as ~ 8 Mpc at $z=0$, and the mass swept into the shell is as large as $M_s\sim 2\times10^{13}\,M_\odot$. This gives an amplification factor $M_s/M_i\sim 40$.

This result agrees well with a more exact calculation by Vishniac et al. (1985); the difference in the radius from the present result is smaller than 10%. Thus we can calculate the expansion of the shock wave with sufficient precision by using the thin-shell approximation.

Figure 1b shows the evolution of a shock wave when the fragmentation and the reexplosion of the shell are taken into consideration. The initial parameters are the same as in the case of figure 1a. The shock wave underwent two episodes of fragmentation and explosion. As a result the present radius reaches ~ 25 Mpc and the shell mass amounts to $\sim 4\times10^{14}\,M_\odot$. The amplification factor is $\sim 10^3$, which is considerably larger than the result from a single explosion as is naturally expected.

Figure 2 shows the evolution of the comoving shell radius $R_{co}=(1+z)R_s$ for various initial seed masses M_i. Here we fix the explosion redshift $z_i=50$ and take the range $M_i=10^6$ to $10^{16} M_\odot$. Small circles denote the epoch and radius at the fragmentation and explosion. Thus the curve not having a small circle shows that this shock wave cannot satisfy the fragmentation condition up to the present epoch and has kept expanding. From this figure we see that the fragmentation of the shell produced by an explosion with as small an energy as $M_i\leqq 10^{10}\,M_\odot$ is difficult. The reason for this result is understood as follows.

When the age of the shock wave t is much larger than the cosmic time at explosion τ_i, the shell radius R_s can be expressed by the power law of t, i.e., $R_s\propto t^\alpha$, where $2/3\lesssim \alpha\lesssim 1$ from selfsimilar solutions (Ostriker and McKee 1988). Since $V_s=\alpha R_s/t$ and $\tau=(2/3)H_0^{-1}(1+z)^{-3/2}\Omega_0^{-1/2}$ if $1+z\gg\Omega_0^{-1}$, the fragmentation condition (6) leads to

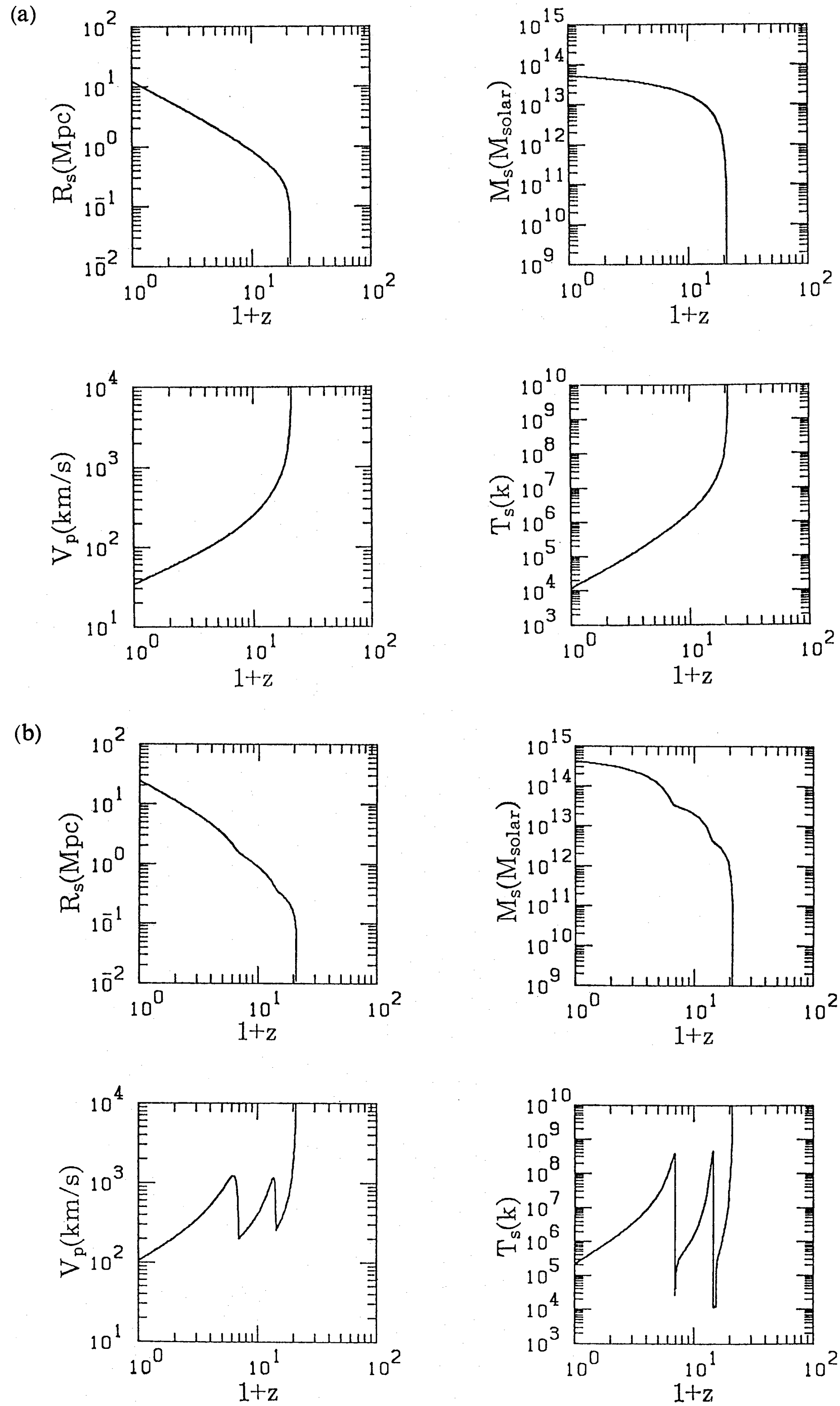


Fig. 1. The evolution of the shock wave produced by an explosion. We assume $\Omega_0 = \Omega_{b,0} = 0.1$ and $h_{50} = 1$, the efficiency of the explosion $\varepsilon = 10^{-4}$, the redshift of the explosion $z_i = 20$, the mass of the explosive object $M_i = 5.79 \times 10^{11}\ M_{\odot}$, corresponding to the initial total energy $E_i = 10^{62}$ erg. In (a) the case which does not consider fragmentation is illustrated and in (b) the case where fragmentations and reexplosions are considered is illustrated.

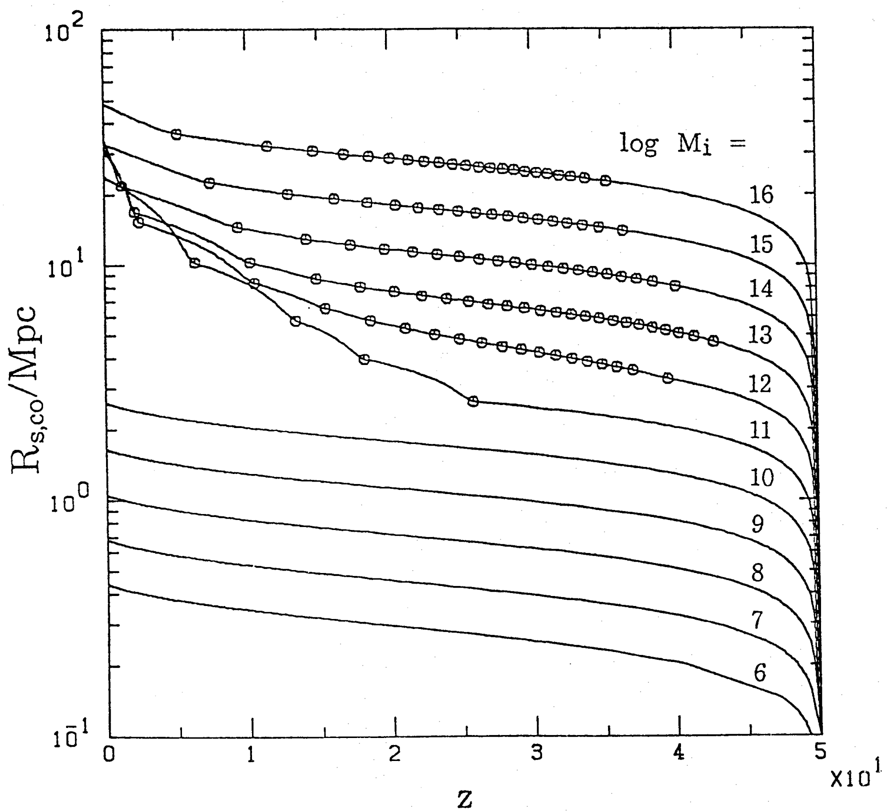


Fig. 2. The evolution of the comoving radii of shells produced by explosions at $z_i=50$ with various initial masses. The small circle represents a fragmentation and a reexplosion. We assume $\Omega_0=\Omega_{b,0}=0.1$, $h_{50}=1$, and $\varepsilon=10^{-4}$.

$$R_s > R_{crit} = 3\pi\left(\frac{3}{2}\right)^{5/2}\alpha\frac{\Omega_0}{\Omega_{b,0}}C_s t \tag{7}$$

or

$$M_s > M_{crit}\ , \tag{8}$$

where

$$\begin{aligned} M_{crit} &= \frac{4\pi}{3}R_{crit}{}^3\rho_g \\ &= \frac{1}{G}(3\pi)^3\left(\frac{3}{2}\right)^{9/2}\Omega_0{}^{3/2}\Omega_{b,0}{}^{-2}\alpha^3 H_0{}^{-1}C_s{}^3(1+z)^{-3/2} \\ &= 9.12\times10^{12}\alpha^3\Omega_0{}^{3/2}\Omega_{b,0}{}^{-2}h_{50}^{-1}\mu^{-3/2}(T_s/10^4K)^{3/2}(1+z)^{-3/2}M_\odot\ . \end{aligned}$$

Therefore, in order to satisfy the fragmentation condition the shell must sweep up a mass greater than a critical mass M_{crit} before the selfgravity of the shell becomes efficient. Since an explosion with a small energy cannot sweep up so much matter, the shell does not suffer the fragmentation.

Figure 2 shows the result for Model A ($\Omega_0=\Omega_{b,0}=0.1$, $\varepsilon=10^{-4}$). From this figure we can see that an explosion of a galactic-scale object (mass $\sim10^{11}\,M_\odot$) produces a shell with radius $\sim$20 Mpc at the present epoch.

4. Anisotropy of the Microwave Background Radiation

Here we investigate the constraints upon the explosion scenario imposed by observations of the microwave background radiation (MBR). Hogan (1984) and Vishniac and Ostriker (1986) pointed out the possibility that explosions which produce the large scale structure of the universe inevitably conflict with the observed isotropy of the MBR.

The MBR photons are scattered to higher energy by hot electrons which are heated up by the shock wave. As a result the brightness temperature T_b decreases in the Rayleigh–Jeans part of the MBR spectrum and increases in the Wien part (Zeldovich and Sunyaev 1969). If the explosions have occurred randomly in space, the number of the explosions contained in the observed beam should fluctuate. Therefore the MBR temperature is different in every direction of the observed beam. In order to produce a larger structure, a larger energy injection is required and thus a larger anisotropy of the MBR is expected to be produced.

Vishniac and Ostriker (1986) estimated the degree of the anisotropy of the MBR. They assumed that the fluctuation of the number N of explosions in a beam is equal to $N^{1/2}$ and estimated the amount of the energy required to produce a 50-Mpc-size bubble from the similarity solution. They concluded that in order to keep the anisotropy less than the upper limit of the observations, explosions should have occurred at $z<5$.

However, they examined only the following restricted cases: (1) the beam size is always larger than the shell size, (2) the thermal energy of the shock wave is constant, and (3) $\Omega_0=1$. Therefore we will relax these constraints and investigate how a large bubble-like structure can be produced by explosions without contradicting the observed isotropy of the MBR.

4.1. Method

Suppose we observe the MBR with a beam having a full width of $\Delta\theta$. Then the cross section of the beam A_b is given by

$$A_b=\pi\left(\frac{D}{2}\right)^2, \tag{9}$$

where D is related to $\Delta\theta$ as

$$\Delta\theta=\frac{D}{(c/H_0)\phi} \tag{10}$$

with

$$\phi^{-1}=\frac{(1+z)^2q_0}{q_0z+(q_0-1)[(1+2q_0z)^{1/2}-1]} .$$

Here q_0 is the present deceleration parameter and $q_0=\Omega_0/2$. The cross section of the shell is $A_s=\pi R_s{}^2$.

(1) $A_b \gg A_s$ case

Here we consider the case where the beam size is sufficiently larger than the shell size. Then the temperature change of the MBR by Compton scattering is given by (Vishniac and Ostriker 1986)

$$\frac{\varDelta T_{\mathrm{r}}}{T_{\mathrm{r}}}=-2\int\frac{\sigma_{\mathrm{T}}}{m_{\mathrm{e}}c^{2}}n_{\mathrm{e}}kT_{\mathrm{e}}dl$$
$$=-\int\frac{2}{3}\frac{\sigma_{\mathrm{T}}}{A_{\mathrm{b}}}\frac{E_{\mathrm{th}}}{m_{\mathrm{e}}c^{2}}\bar{N}_{\mathrm{s}}A_{\mathrm{b}}dl\,, \tag{11}$$

where σ_{T} is the Thomson cross section, m_{e} is the electron mass, c is the light velocity, n_{e} is the electron number density, T_{e} is the electron temperature, and E_{th} is the thermal energy of a single explosion. The y-parameter in the Kompaneets equation is essentially the same as this temperature decrement. The integration is performed along the line of sight with

$$dl=cdt=-cH_{0}^{-1}(1+z)^{-2}\,(1+\varOmega_{0}z)^{-1/2}dz\,. \tag{12}$$

In the derivation of the second equality of equation (11) we assume that the gas is fully ionized and that $\bar{N}_{\mathrm{s}}$ is the average number density of the exploding seeds. Therefore, the MBR temperature is observed to decrease on the average by

$$(\varDelta T_{\mathrm{r}})_{\mathrm{AV}}=-\left(\int\frac{2}{3}\frac{\sigma_{\mathrm{T}}}{A_{\mathrm{b}}}\frac{E_{\mathrm{th}}}{m_{\mathrm{e}}c^{2}}\bar{N}_{\mathrm{s}}A_{\mathrm{b}}dl\right)T_{\mathrm{r}}\,. \tag{13}$$

However, since the explosion sources should not be distributed homogeneously, the value of $\varDelta T_{\mathrm{r}}$ will fluctuate in each observed direction. Each observation indicated by the subscript j gives the temperature decrement,

$$(\varDelta T)_{j}=\left(\int f(z)\bar{N}_{\mathrm{s},j}A_{\mathrm{b}}dl\right)T_{\mathrm{r}}\,, \tag{14}$$

where

$$f(z)=\frac{2}{3}\frac{\sigma_{\mathrm{T}}}{A_{\mathrm{b}}}\frac{E_{\mathrm{th}}}{m_{\mathrm{e}}c^{2}}\,.$$

In fact the observations are made by the beam-switching method. Therefore, the quantity

$$(\varDelta T)_{\mathrm{obs},j}=(\varDelta T)_{j}-(\varDelta T)_{\mathrm{AV}}$$
$$=\left[\int f(z)(\bar{N}_{\mathrm{s},j}-\bar{N}_{\mathrm{s}})A_{\mathrm{b}}dl\right]T_{\mathrm{r}} \tag{15}$$

should be observed. First, we transform the integration of z into the form of summation,

$$\frac{(\varDelta T)_{\mathrm{obs},j}}{T_{\mathrm{r}}}=\lim_{n\to\infty}\sum_{k=1}^{n}f(z_{k})\delta N_{j,k}\,, \tag{16}$$

where

$$\delta N_{j,k}=[(\bar{N}_{\mathrm{s},j}-\bar{N}_{\mathrm{s}})A_{\mathrm{b}}\varDelta l]_{k}$$

and

$$\Delta l = c\Delta t = -cH_0^{-1}(1+z)^{-2}(1+\Omega_0 z)^{-1/2}\Delta z\,.$$

Dividing the line of sight into n regions with the redshift width Δz,

$$\begin{aligned}\left[\frac{(\Delta T)_{\mathrm{obs},j}}{T_{\mathrm{r}}}\right]^2 &= \lim \sum_k f(z_k)\delta N_{j,k} \lim \sum_l f(z_l)\delta N_{j,l} \\ &= \lim \sum_k [f(z_k)]^2[\delta N_{j,k}]^2 \\ &\quad + \lim \sum_{k\neq l} f(z_k)\delta N_{j,k} f(z_l)\delta N_{j,l}\,, \end{aligned} \tag{17}$$

and the averaged square temperature fluctuation is given by

$$\begin{aligned}\left\langle\left[\frac{(\Delta T)_{\mathrm{obs},j}}{T_{\mathrm{r}}}\right]^2\right\rangle_j &= \lim \sum_k [f(z_k)]^2\langle(\delta N_{j,k})^2\rangle_j \\ &\quad + \lim \sum_{k\neq l} \langle f(z_k)\delta N_{j,k} f(z_l)\delta N_{j,l}\rangle_j\,. \end{aligned} \tag{18}$$

If the spatial distribution of explosion sources is the Poisson distribution, we have

$$\begin{aligned}\langle(\delta N_{j,k})^2\rangle_j &= \langle[(\bar{N}_{\mathrm{s},j}A_{\mathrm{b}}\Delta l)_k - (\bar{N}_{\mathrm{s}}A_{\mathrm{b}}\Delta l)_k]^2{}_j \\ &= (\bar{N}_{\mathrm{s}}A_{\mathrm{b}}\Delta l)_k\,. \end{aligned} \tag{19}$$

When the regions k and l are independent of each other, it should be

$$\langle f(z_k)\delta N_{j,k} f(z_l)\delta N_{j,l}\rangle_j \to 0\,,$$

and from equations (18) and (19), we have

$$\begin{aligned}\left\langle\left[\frac{(\Delta T)_{\mathrm{obs},j}}{T_{\mathrm{r}}}\right]^2\right\rangle_j &= \lim \sum_k [f(z_k)]^2(\bar{N}_{\mathrm{s}}A_{\mathrm{b}}\Delta l)_k \\ &= \int [f(z)]^2 \bar{N}_{\mathrm{s}}A_{\mathrm{b}}dl\,. \end{aligned}$$

Thus finally

$$\left\langle\left(\frac{\Delta T_{\mathrm{r}}}{T_{\mathrm{r}}}\right)^2\right\rangle = \int\left[\frac{2}{3}\left(\frac{\sigma_{\mathrm{T}}}{A_{\mathrm{b}}}\right)\left(\frac{E_{\mathrm{th}}}{m_{\mathrm{e}}c^2}\right)\right]^2 \bar{N}_{\mathrm{s}}A_{\mathrm{b}}dl \tag{20}$$

is derived.

(2) $A_{\mathrm{b}} < A_{\mathrm{s}}$ case

In this case the size of the beam is smaller than that of the shell. The temperature change of MBR along the beam running through an explosion shell is given by

$$\frac{\Delta T_{\mathrm{r}}}{T_{\mathrm{r}}} = -\frac{2}{3}g\frac{\sigma_{\mathrm{T}}}{m_{\mathrm{e}}c^2}\frac{E_{\mathrm{th}}}{A_{\mathrm{s}}}\,, \tag{21}$$

where g is the geometrical factor $g \sim O(1)$ depending upon the spatial structure of the shell. Since the expected number for one beam to run through is $\bar{N}_{\mathrm{s}}A_{\mathrm{s}}dl$, by the similar discussion in the above the average temperature decrease and its fluctuation are given, respectively, by

$$\left(\frac{\Delta T_r}{T_r}\right)_{AV} = -\int \frac{2}{3} g \frac{\sigma_T}{A_s} \frac{E_{th}}{m_e c^2} \bar{N}_s A_s dl \tag{22}$$

and

$$\left\langle\left(\frac{\Delta T_r}{T_r}\right)^2\right\rangle = \int \left[\frac{2}{3} g \frac{\sigma_T}{A_s} \frac{E_{th}}{m_e c^2}\right]^2 \bar{N}_s A_s dl \,. \tag{23}$$

(3) Interpolation of the above two cases

Since both the sizes of the shell and the beam should change as the time goes on, one beam may observe the shells for both cases. Therefore, we must calculate the anisotropy by connecting both cases. A simple interpolation method would be

$$\left(\frac{\Delta T_r}{T_r}\right)_{AV} = -\int \frac{2}{3} \left(\frac{\sigma_T}{A_c}\right) \left(\frac{E_{th}}{m_e c^2}\right) \bar{N}_s A_c dl \tag{24}$$

and

$$\left\langle\left(\frac{\Delta T_r}{T_r}\right)^2\right\rangle = \int \left[\frac{2}{3} \left(\frac{\sigma_T}{A_c}\right) \left(\frac{E_{th}}{m_e c^2}\right)\right]^2 \bar{N}_s A_c dl \,, \tag{25}$$

where

$$A_c = \pi \left(\frac{D}{2} + R_s\right)^2 .$$

In this calculation the thermal energy of the shell E_{th} is calculated exactly, taking into account adiabatic cooling and radiative cooling as in section 2.

(4) Number of seeds

If seed objects are formed simultaneously at the redshift $z=z_i$, after which no new sources are formed, the number of seed explosions is conserved in the comoving coordinate. As a convenient method for determining $\bar{N}_{s,i}$, the number density of seeds at $z=z_i$, we take the following condition for overlapping at $z=z_{over}$,

$$\frac{4\pi}{3} R_s^3(z_{over}) \bar{N}_s(z_{over}) = 1 \,, \tag{26}$$

from which $\bar{N}_{s,i}$ can be determined as $\bar{N}_{s,i} = \bar{N}_s(z_{over})(1+z_i)^3/(1+z_{over})^3$.

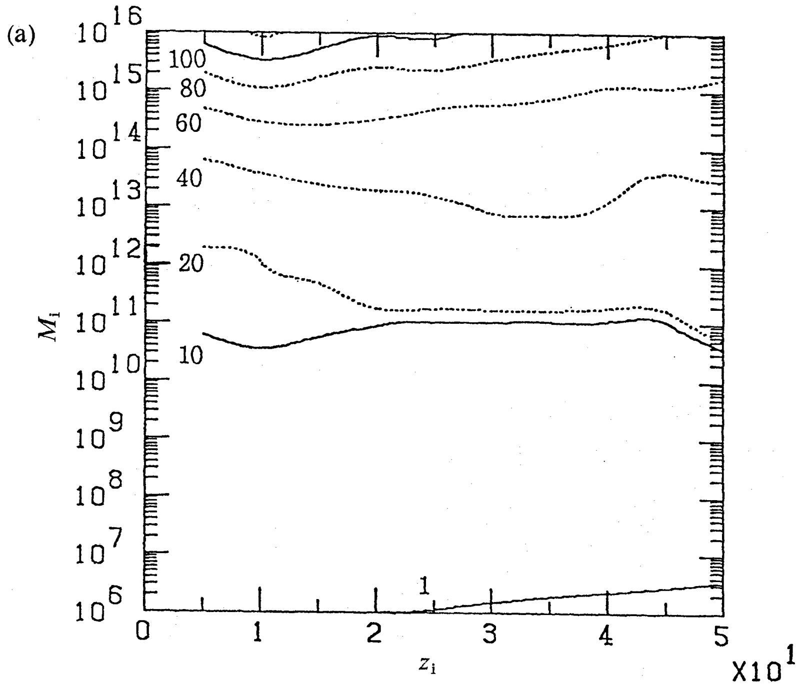
We take $z_{over}=3$ as the standard value, since (1) at present the expanding shells produced by explosions are not observed, (2) but the shells should have overlapped in order to explain the presently observed bubbly distribution of galaxies, and (3) the universe has been already reionized at $z=3.8$ from the Gunn–Peterson test. The full beam width $\Delta\theta$ is taken to be 1′.5 in order to compare with the observation by Uson and Wilkinson (1984).

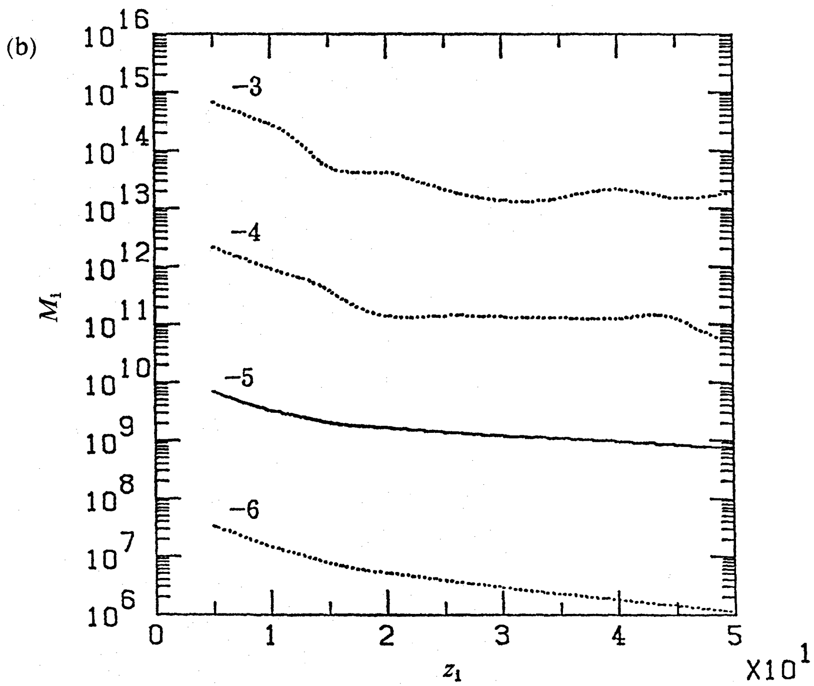
4.2. *Results*

The parameters of the models we calculated are shown in table 1. The results of the calculations for model A are shown in figure 3. Figure 3a shows the present

Table 1. Parameters of the models.

Model	Ω_0	Ω_b	h_{50}	ε	z_{over}
A	0.1	0.1	1	10^{-4}	3
B	0.1	0.1	1	10^{-4}	0.5
C	0.1	0.1	1	10^{-4}	5
D	0.1	0.1	1	10^{-5}	3
E	0.1	0.1	1	10^{-2}	3
F	0.1	0.1	1.5	10^{-4}	3
G	0.1	0.1	2.0	10^{-4}	3





Figs. 3a and b. See the legend on the next page.

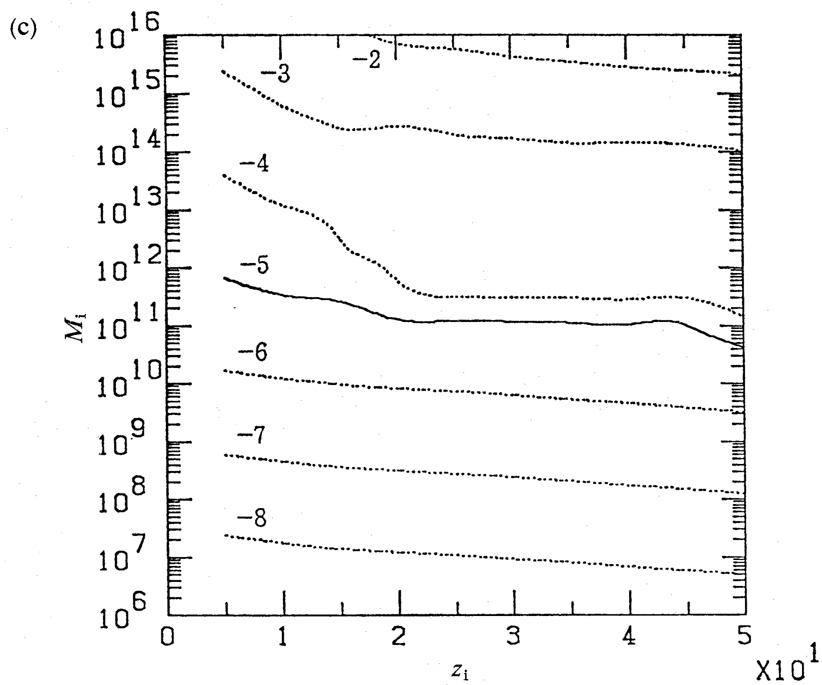


Fig. 3. The MBR anisotropy produced by the explosions. We take $\Omega_0=\Omega_{b,0}=0.1$, $h_{50}=1$, and $\varepsilon=10^{-4}$. The redshift of the overlapping of the shells is $z_{over}=3$ and the full width of the beam observing MBR is $\Delta\theta=1'.5$. (a) The comoving radius of the shell at the overlapping. (b) The temperature decrease of the MBR, $\log|\Delta T_r/T_r|$. (c) The anisotropy of the MBR produced by the explosions, $\log\delta$.

comoving radius $R_s(0)=R_s(z_{over})(1+z_{over})$ of the shell at the overlapping epoch in the z_i-M_i plane in the form of a contour map. Figures 3b and 3c show, respectively, the absolute value of the distortion of the MBR temperature $|\Delta T_r/T_r|$ in the logarithmic scale and the anisotropy of the MBR temperature $\delta=\langle(\Delta T_r/T_r)^2\rangle^{1/2}$.

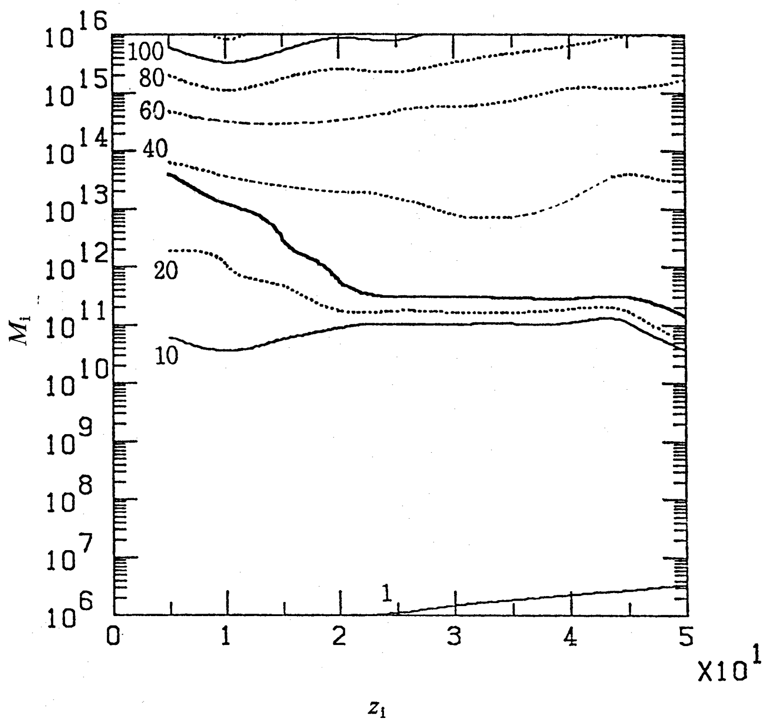
From these results we find the following:

(1) Galactic-scale explosions ($M_i\sim10^{11-12}M_\odot$) can produce $\sim$10-Mpc shells by $z_{over}=3$.

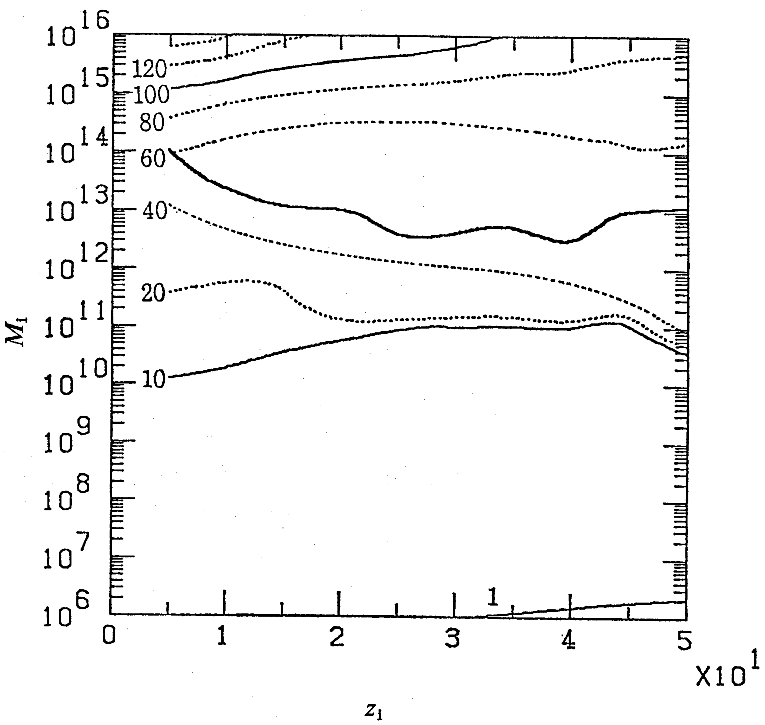
(2) The distortion of the MBR temperature $|\Delta T_r/T_r|$ is at most 10^{-3} in this model A. Therefore, it is difficult to detect it with the present accuracy of the observation of the MBR distortion. [Recently, Matsumoto et al. (1988) detected a possible distortion of the MBR in the Wien part. Their result indicates $|\Delta T_r/T_r|\sim10^{-2}$. The relation of the present model to this observation is discussed in another paper (Yoshioka and Ikeuchi 1987)].

(3) Since it may be possible that the anisotropy δ exceeds 10^{-4}, it is easily observed.

The observation of Uson and Wilkinson (1984) gives the most stringent upper limit to the MBR anisotropy. They observed the regions separated by 4'.5 from each other with the beam-switching method (full beam width $\Delta\theta=1'.5$) and obtained $\delta<4.5\times10^{-5}$. It is meaningless to compare directly our computational results with their observation, since their observational method is not a simple beam-switching one and we assume that the regions separated by the beam-switching angle are not correlated. Then, we relax the condition that our computed value of the MBR anisotropy δ should be less than 10^{-4}.

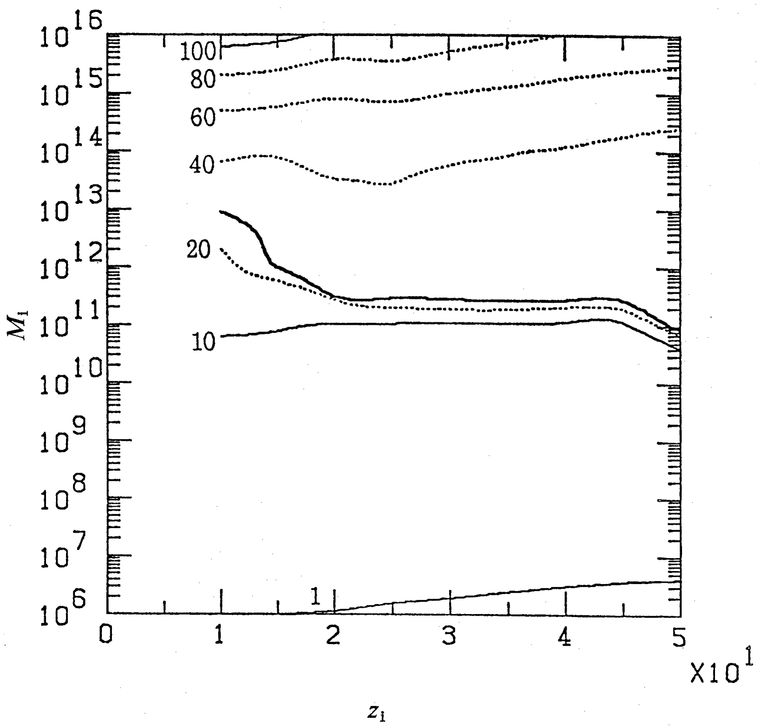


(a)

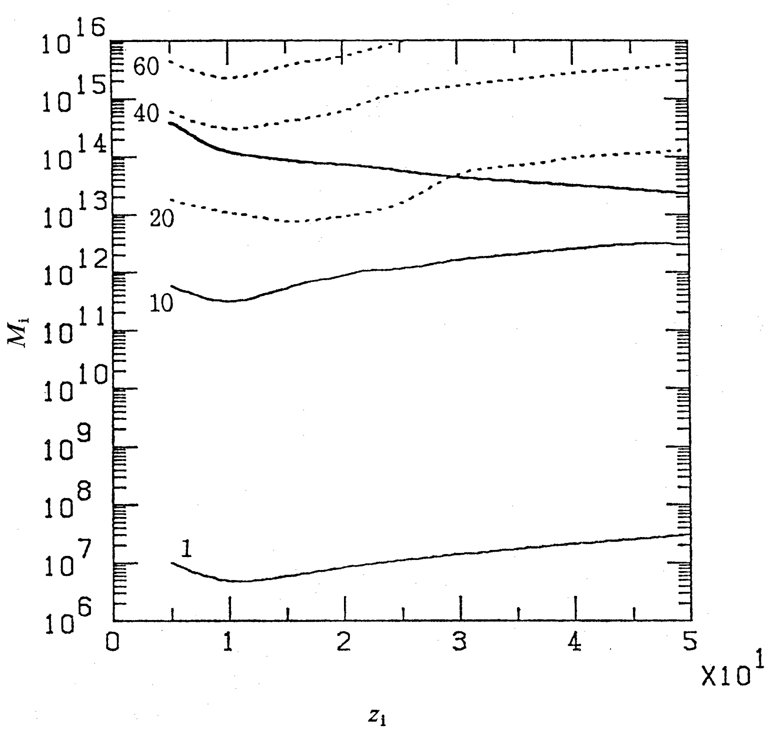


(b)

Figs. 4a and b. See the legend on page 399.

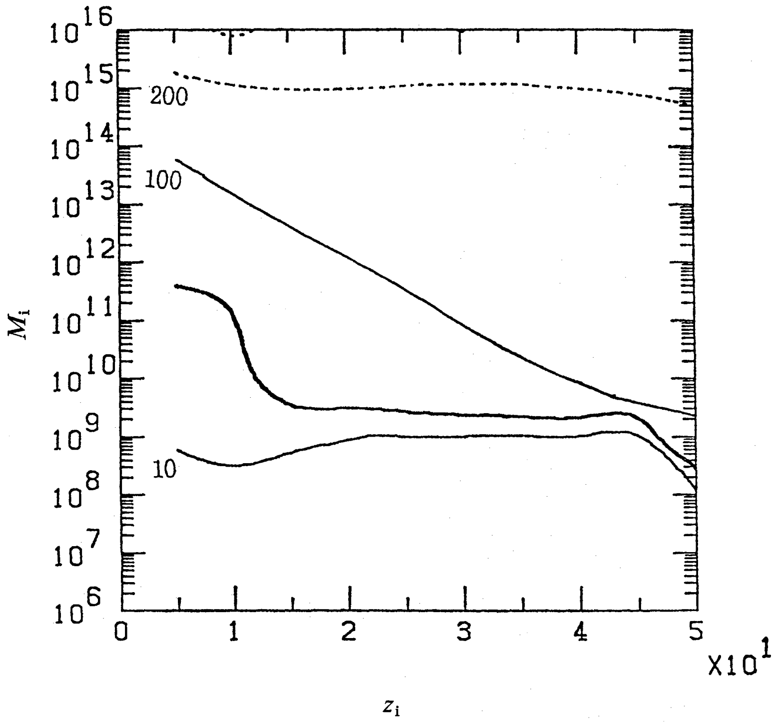


(c)

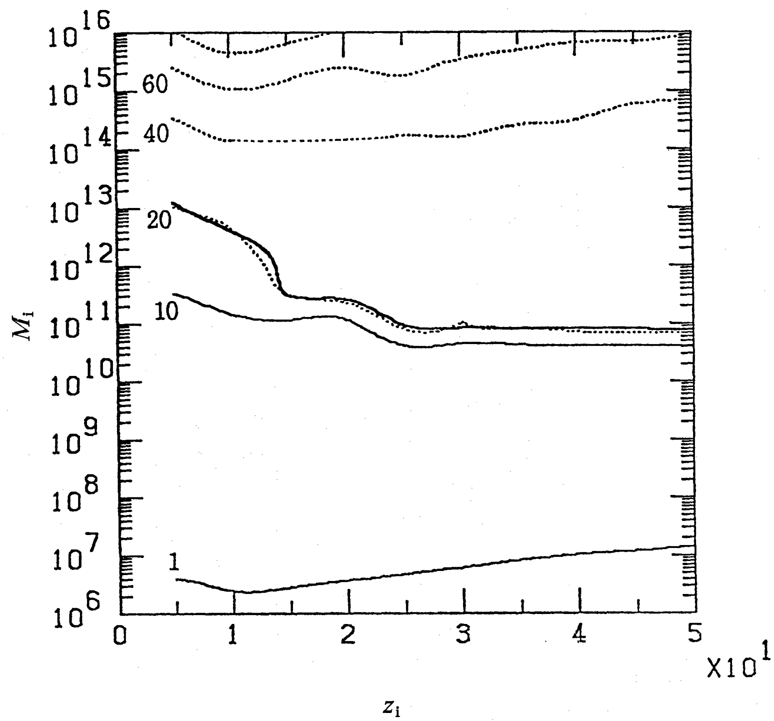


(d)

Figs. 4c and d. See the legend on page 399.

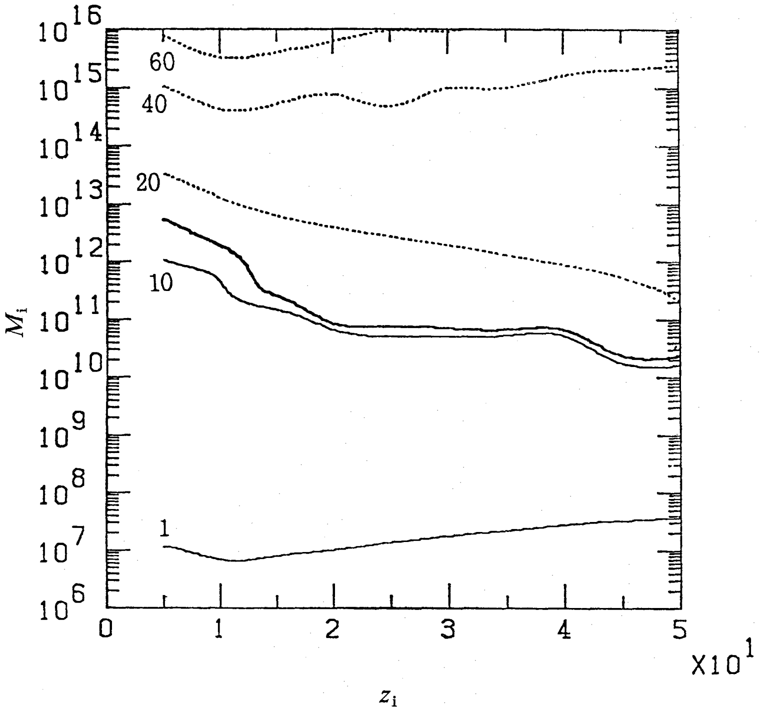


(e)



(f)

Figs. 4e and f. See the legend on page 399.



(g)

Fig. 4. The constraint from the observed MBR isotropy on the comoving diameter of the bubbles produced by the explosions. The dotted and solid lines represent the comoving diameters of the bubbles at the overlapping. The curves are labeled with the values of the diameter in units of Mpc. The predicted anisotropy of the MBR exceeds 10^{-4} above the thick solid line. The model parameters are shown in table 1. The results in models A–G are shown respectively in figures 4a–g.

Figure 4 shows what kind of large-scale structure can arise by the overlapping time, satisfying the above condition. The size of the structure is representd by the diameter of the shell. Figure 4a shows the result for model A. This indicates that the cosmic explosions can make the structure with the comoving diameter $D_{co} \sim 30$ Mpc almost independently of z_i. Figures 4b and c show, respectively, the results for models B and C which have the same parameter as model A except for z_{over}. In model B ($z_{over}=0.5$) the structure with $\sim$60-Mpc scale can form. Figures 4d and e show the results for the models having a different explosion efficiency ε for comparison as model D ($\varepsilon=10^{-5}$) in figure 4d and model E ($\varepsilon=10^{-2}$) in figure 4e. While the size of the shell for the same z_i and M_i changes in proportion to the efficiency ε, the scale of the structure which can be formed without contradicting the observed MBR isotropy hardly changes.

The larger the Hubble constant is, the smaller is the permissible size. We calculate the models with $h_{50}=1.5$ (model F) in figure 4f and $h_{50}=2$ (model G) in figure 4g. We obtain $D_{co} \lesssim 20$ Mpc for model F and $D_{co} \lesssim 14$ Mpc for model G. For $h_{50}=1$ (model A) we have already obtained $D_{co} \lesssim 30$ Mpc. Then we can approximate the comoving diameter of the bubbles which can be produced without contradicting the MBR isotropy

$$D_{co} \lesssim 30\, h_{50}{}^{-1}\ \mathrm{Mpc}$$

for $\Omega_0=\Omega_{b,0}=0.1$, $\varepsilon=10^{-4}$, $z_{over}=3$ and $\delta<10^{-4}$.

5. Summary and Discussion

5.1. Summary

1) The shocked shell produced by the explosion is cooled by radiative and Compton cooling, and then it fragments by its selfgravity if $M_i \gtrsim 10^{11} M_{\odot}$ in the case where the energy release efficiency is $\varepsilon=10^{-4}$.

2) If the fragmentation and explosions of shells repeat successively, an explosion with the mass $\sim 10^{11} M_{\odot}$ would finally produce the bubbles with the diameter ~ 40 Mpc ($h_{50}=1$).

3) The bubbles with the diameter $\sim 30 h_{50}{}^{-1}$ Mpc can form without contradicting the observed MBR isotropy if the shells overlap by $z=3$.

5.2. Discussion

We compare our results with the structure observed in the universe. Here we concentrate on the scale of the bubble structure. Vishniac and Ostriker (1986) restricted the radius of the shell-like structure to those larger than the coherence length of the galaxy distribution $R_0=18 h_{50}{}^{-1}$ Mpc. If we adopt this condition, the explosion scenario will be hardly permitted from the above summary (3). However, the coherence length of the galaxy correlation function may not be a suitable index for characterizing the scale length of bubbles.

According to de Lapparent et al. (1986), a typical diameter of galaxy voids is $50 h_{50}{}^{-1}$ Mpc and the largest void is $\sim 100 h_{50}{}^{-1}$ Mpc. The largest structure produced by the explosions is found to be $\sim 30 h_{50}{}^{-1}$ Mpc in diameter for model A. This seems considerably smaller than the typical bubble diameter. The simple explosion model is too naive to produce the bubbly structure of the universe with a scale greater than $30 h_{50}{}^{-1}$ Mpc. If the largest void is a rare event in the universe, we may consider that it is an exceptional case. A large void with the diameter $\sim 100 h_{50}{}^{-1}$ Mpc is also found in Bootes by Kirshner et al. (1981). Then, such a large void is so common that it seems difficult to produce it by a single explosion.

Some modifications of the original idea should be examined. The merging of shells is one of them. Even if the diameter of each shell is less than the MBR limit, the merging of several such shells may produce a larger structure, though the shape and size of the structure after the merging depends upon the era of the merging and the number density of the explosions.

Another possibility is that the overlapping of the shell occurred recently. In section 4 we take $z_{over}=3$, particularly because in the spectrum of the far distant quasar ($z=3.8$) no absorption trough of the intervening neutral hydrogen is detected. However, if the reionization of the intergalactic gas is due to the UV photons emitted from the first-generation objects or quasars and not due to the cosmological shock waves, we need not take $z_{over}=3$.

In model B (figure 4b) we take $z_{over}=0.5$. In this model it is possible to produce bubbles with diameter $\sim 60 h_{50}{}^{-1}$ Mpc. This assumption does not conflict with the existence of galaxies at $z\sim 1$. It is probable that galaxies were formed by the fragmentation of the shell at $z>1$, and after that these galaxies expanded keeping the nearly shell-like distribution. After that, the shells of the galaxies have overlapped at $z=0.5$.

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References

Bertschinger, E. 1983, *Astrophys. J.*, **268**, 17.
Bertschinger, E. 1985, *Astrophys, J.*, **295**, 1.
Bond, J. R., Carr, B. J., and Hogan, C. J. 1986, *Astrophys. J.*, **306**, 428.
Carr, B. J., Bond, J. R., and Arnett, W. D. 1984, *Astrophys. J.*, **277**, 445.
Carr, B. J., and Ikeuchi, S. 1985, *Monthly Notices Roy. Astron. Soc.*, **213**, 497.
Davis, M., Huchra, J. P., Latham, D. W., and Tonry, J. 1982, *Astrophys. J.*, **253**, 423.
Davis, M., and Peebles, P. J. E. 1983a, *Astrophys. J.*, **267**, 465.
Davis, M., and Peebles, P. J. E. 1983b, *Ann. Rev. Astron. Astron. Astrophys.*, **21**, 109.
de Lapparent, V., Geller, M., and Huchra, J. P. 1986, *Astrophys. J. Letters*, **302**, L1.
Einasto, J., Jôeveer, M., and Saar, E. 1980, *Monthly Notices Roy. Astron. Soc.*, **193**, 353.
Hogan, C. J. 1984, *Astrophys. J. Letters*, **284**, L1.
Ikeuchi, S. 1981, *Publ. Astron. Soc. Japan*, **33**, 211.
Ikeuchi, S., Tomisaka, K., and Ostriker, J. P. 1983, *Astrophys. J.*, **265**, 583.
Kirshner, R., Oemler, A., Jr., Schechter, P. L., and Shectman, S. A. 1981, *Astrophys. J. Letters*, **248**, L57.
Matsumoto, T., Hayakawa, S., Matsuo, H., Murakami, H., Sato, S., Lange, A. E., and Richards, P. L. 1988, *Astrophys. J.*, **329**, 567.
Ostriker, J. P., and Cowie, L. L. 1981, *Astrophys. J. Letters*, **243**, L127.
Ostriker, J. P., and McKee, C. F. 1988, *Rev. Mod. Phys.*, **60**, 1.
Umemura, M., and Ikeuchi, S. 1984, *Prog. Theor. Phys.*, **72**, 47.
Uson, J. M., and Wilkinson, D. T. 1984, *Astrophys. J.*, **283**, 471.
Vishniac, E. T. 1983, *Astrophys. J.*, **274**, 152.
Vishniac, E. T., and Ostriker, J. P. 1986, Soc. Ital. Fis., **1**, 157.
Vishniac, E. T., Ostriker, J. P., and Bertschinger, E. 1985, *Astrophys. J.*, **291**, 399.
Yang, J., Turner, M. S., Steigmann, G., Schramm, D. N., and Olive, K. A. 1984, *Astrophys. J.*, **281**, 493.
Yoshioka, S., and Ikeuchi, S. 1987, *Astrophys. J. Letters*, **323**, L7.
Zeldovich, Ya. B., and Sunyaev, R. A. 1969, *Astrophys. Space Sci.*, **4**, 301.