THE ASTROPHYSICAL JOURNAL, 337:1–6, 1989 February 1

ON THE EXPLOSIVE ORIGIN HYPOTHESIS FOR PRESENT-DAY COSMIC VOIDS AND PECULIAR VELOCITIES

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Received 1987 November 2; accepted 1988 July 7

ABSTRACT

General similarity solutions are obtained for spherical, cooling blast waves in an $\Omega_b \ll \Omega = 1$ universe, given an energy input that is some power of time. The distortion of the microwave background is calculated in terms of the present-day peculiar velocities and the present-day voids in velocity space caused by the explosions under the assumption of sustained energy input. The predicted distortion is found to be marginally consistent with the present-day observations and could provide a test of the model in the near future. It is briefly noted that high peculiar velocities of galaxies at the epoch of their formation make dark matter capture by them questionable and also that explosions in a shadow sector would trivially overcome some of the potential problems of the model.

Subject headings: cosmic background radiation — cosmology

I. INTRODUCTION

A possible scenario to explain the apparent voids in the universe that have been recently reported (Davis *et al.* 1982) is that explosions in the early universe swept away the ambient matter, forming shells with evacuated interiors. Conjectured causes of such explosions include exploding Population III objects (Carr, Bond, and Arnett 1984), contraction of superconducting strings in the presence of a primordial magnetic field (Ostriker, Thompson, and Witten 1986), supernovae in the earliest galaxies (Ostriker and Cowie 1981), etc. "Explosive amplification" has also been invoked by various authors (e.g., Ostriker and Cowie 1981; Carr and Ikeuchi 1985) to explain the formation of galaxies as these shells expand, cool, and then undergo fragmentation.

In this paper we examine the distortion in the microwave background produced by the expanding shells. Assuming that they are thin and cool and that the universe is mostly "dark matter" we find a general class of similarity solutions. We then show that the similarity solutions are stable to the fluctuations in the initial conditions. In § II we show that most of the thermal energy stored in the shell is radiated mainly by inverse Compton scattering, for the important range of the temperatures and redshifts. We then estimate the distortion of the microwave background in terms of the peculiar velocities of present-day galaxies, the epoch of their formation and the sizes of the voids. We find that the distortion is insensitive to the type of the time evolution of the expanding shells but constrains the present-day peculiar velocities of the galaxies in terms of the epoch of their formation. Finally we consider the possible case of the "dark matter" left inside catching up with the baryonic shell.

Similar calculations concerning the distortion of the microwave background for particular assumptions have been published by Bertschinger (1983), Hogan (1984), Ostriker, Thompson, and Witten (1986). Our analytic approach makes it somewhat more general and yields a closed form result in which the dependence on various parameters, including the present-day peculiar velocities, appears explicitly. Also, the dissipation of the expanding shells is taken into account to make the calculations more relevant to the epoch of galaxy formation.

For a comprehensive review of the cosmological blast wave solutions themselves, see Ostriker and McKee (1988) and the references therein. While our approach is slightly different, it recovers the same similarity exponents in the appropriate limits.

II. EFFECTS OF COOLING

One can put a limit on z of the epoch of the formation of galaxies by the fact that in order to form galaxies a gas has to be cooled faster than the Hubble time. This implies that both the ion-electron relaxation time and the electron cooling time are shorter than the Hubble time. Assuming the cooling is due to inverse Comptonization of the blackbody background, it can be shown that $(1 + z)^4 > 248.8\,(h_0)$, where $H = 100\,h_0$ km s^{-1} Mpc^{-1}.

To check that free-free emission is less important, we must look at the rate of radiation by free-free and inverse Compton effect and the relaxation time between the electrons and protons to find out the regime in which the inverse Comptonization of the microwave background is predominant. Assuming that electrons gain energy by collisions with the protons and radiate by free-free and inverse Compton processes, we can write the energy equation

$$\frac{3}{2} n_e k \frac{T_p - T_e}{t_{\rm eq}} - \left.\frac{dE}{dv\,dt}\right|_{ff} - \left.\frac{dE}{dv\,dt}\right|_{\rm IC} \cong 0\,, \tag{2.1}$$

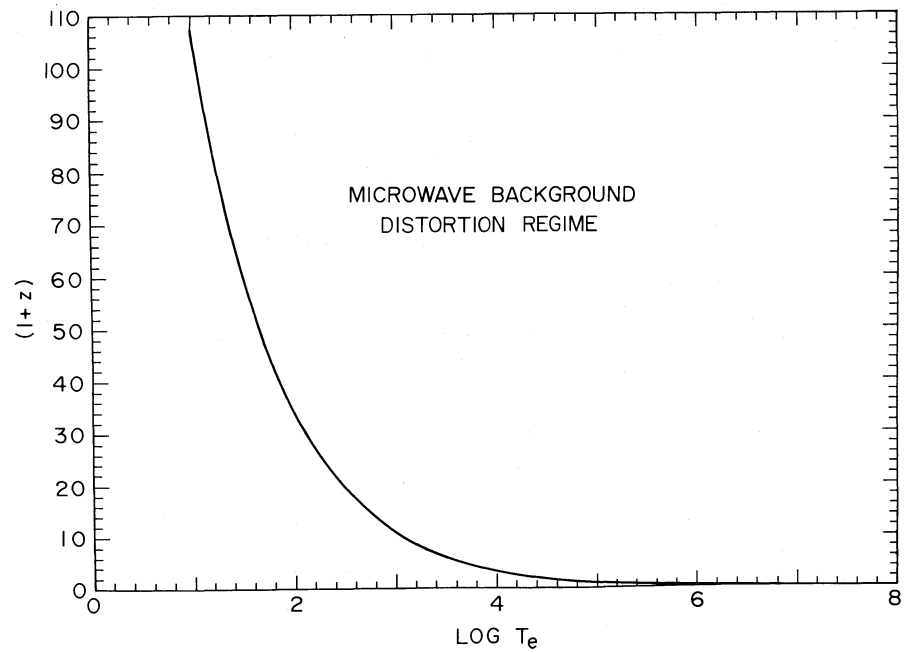

FIG. 1.—The domain in which the distortion of the microwave background by the inverse Compton scattering is dominant

where T_e and T_p are the temperatures in degrees kelvin of electrons and the protons, respectively, n_e is the number of the electrons or protons, and

$$t_{\rm eq} = \frac{3m_e m_p k^{3/2}}{8\sqrt{2\pi} n_e e^4 \ln \Lambda}\left(\frac{T_e}{m_e} + \frac{T_p}{m_p}\right)^{3/2} \tag{2.2}$$

is the relaxation time (Spitzer 1962). Assuming that baryonic density is $4\rho_b$ behind the shock front, where ρ_b is the average cosmic baryon density, and that $\Omega_b = \rho_b/\rho_c \cong 0.1$, where ρ_c is the critical density, we get

$$\frac{T_p - T_e}{T_e^{3/2}} - 1.03 \times 10^{-10} T_e^{1/2} - 3.03 \times 10^{-13} T_e(1 + z) \cong 0 , \tag{2.3}$$

where T_e is in degrees kelvin.

Figure 1 shows the domain in which the distortion of the microwave background by inverse Compton scattering is dominant. Because we are considering quite strong shocks, and, as we shall soon see, because most of the distortion occurs at relatively small z ($z \sim 10$), we can safely assume that most of the thermal energy goes into the microwave background. Also the ion-electron relaxation time is comparable to, or shorter than, the Hubble time for $z \geq 7$ for relevant peculiar velocities of the expanding shell so that we can assume an almost instant Comptonization.

III. THE SIMILARITY SOLUTION

Similarity solutions for noncooling constant energy blast waves have been obtained (e.g., Bertschinger 1983, 1985*a*) in which an ordinary differential equation is derived for the blast wave structure and then solved numerically. Here we assume that cooling keeps the shell thin and that the energy in the interior is resupplied by some sources and may even increase with time. The equation of motion for the baryonic shell can be written as

$$\frac{d}{dt}(\delta m \dot{r}) = -\left[(1 - \Omega_b)\Omega_b + \frac{1}{2}\Omega_b^2\right]\frac{G[(4\pi/3)r^3\rho_c]^2}{r^2} + \frac{2}{3}\frac{E(t)}{(4\pi/3)r^3}4\pi r^2 + \left[\frac{d}{dt}(\delta m)\right]Hr . \tag{3.1}$$

Here $\delta m = 4\pi/3r^3\rho_c\Omega_b$ is the mass of the shell, $r(t)$ is the radius, and H is the Hubble constant. The second term on the right represents the pressure due to energy E in the evacuated interior (E is not the energy of the blast wave) and the factor $\frac{2}{3}$ assumes hot nonrelativistic gas in the shell's interior (assuming the pressure to be uniform). If the pressure is in the form of hot gas at high redshift, it is necessary to assume that it is in ions rather than electrons and that the gas is sufficiently tenuous to avoid collisional energy transfer from one to the other. In any case, the scenario probably works best if the energy is released behind the dense shell.

The first term consists of two parts: (*a*) the gravitational force between the dark matter inside (which is not swept out by the expanding shell) and the shell and (*b*) the self-gravity of the shell.

If we assume that the explosion begins at $t_0 \cong 0$ and $E = E_0(t/t_0)^\Gamma$, where t is the cosmic time, then in an expanding flat universe this equation admits a self-similar solution,

$$r = (9E_0 G)^{1/5}\left[\Omega_b\left(4\alpha^2 - 5\alpha + \frac{4}{3}\right) + \frac{2}{9}\left(\Omega_b - \frac{\Omega_b^2}{2}\right)\right]^{-1/5} t^\alpha , \tag{3.2}$$

where $\alpha = \frac{1}{5}(\Gamma + 4)$. For example, if $\Gamma = 0$, then $\alpha = \frac{4}{5}$.

This form of the similarity solution was obtained previously by Schwartz, Yahil, and Ostriker (1975) and Ikeuchi, Tomisaka, and Ostriker (1983), though they assumed neither radiative losses nor energy input while we assume both.

For $\Gamma = 1$, which is expected from cosmic superconducting string dissipation, $r \propto t$, which recovers the result of Ostriker, Thompson, and Witten (1986), hereinafter OTW).

It should be noted that the pressure term represents any pressure driving the shell outward, e.g., very low frequency electromagnetic waves from cosmic strings (OTW). The case of expanding shells with no interior pressure is recovered by putting $E = 0$ in equation (3.1). Then, for $\Omega_b = 1$, $\alpha = [15 + (17)^{1/2}]/24 \approx 0.79$ and, for $\Omega_b = 0$, $\alpha = \frac{2}{3}$ as also observed by Ostriker and McKee (1988) and Bertschinger (1985*b*).

IV. STABILITY OF THE SHELLS

Writing the self-similar solution as $r = At^\alpha$ and then adding a small perturbation, i.e., letting $r = \delta r(t) + At^\alpha$ and neglecting the second and higher order terms, we can get an equation for $\delta\ddot{r}(t)$, $\delta\dot{r}(t)$, $\delta r(t)$:

$$\delta\ddot{r} + (6\alpha - 4)\frac{\delta\dot{r}}{t} + \left(13\alpha^2 - 20\alpha + \frac{70}{9} - \frac{5}{9}\Omega_b\right)\frac{\delta r}{t} = 0 . \tag{4.1}$$

If $\delta r \propto t^\beta$, then we get a quadratic equation in β, viz.,

$$\beta^2 + (5\alpha - 4)\beta + \left(13\alpha^2 - 20\alpha + \frac{70}{9} - \frac{5}{9}\Omega_b\right) = 0 , \tag{4.2}$$

for which the real part of both the solutions are either negative, for

$$\alpha > \frac{20 + [400 - 52(70/9 - 5/9\Omega_b)]^{1/2}}{26} ,$$

or positive but smaller than α, for

$$\frac{4}{5} \leq \alpha < \frac{20 + [400 - 52(70/9 - 5/9\Omega_b)]^{1/2}}{26}$$

and so $(\delta r/r)$ does not grow in time. The lower limit of α comes from the fact that equation (4.1) is true only for $\Gamma \geq 0$. Hence, if $E(t)$ has the form t^Γ for all but very early times, the similarity solution is physical.

V. EXPANSION AFTER GALAXY FORMATION

We assume that after the shell condenses into galaxies, the pressure term in the original equation does not affect its motion, and in the "coasting phase," the equation of motion can be shown to become almost that of a test particle. This is because most of the matter in the universe is assumed to be weakly interacting dark matter and is unaffected by the explosion. Thus the departures from the test particle solutions are insignificant and we can assume that after the epoch of galaxy formation $V_p \propto (1 + z)$, V_p being the peculiar velocity of the shell. The tacit assumption is made, however, that the dark matter passed by the shell has not yet overtaken it. The opposite assumption is considered below in § VIII.

VI. ENERGY STORED IN THE SHELL

Typical present-day peculiar velocities are claimed to be of the order of a few 100 km s^{-1} (Collins, Joseph, and Robertson 1986; Dressler *et al.* 1987). We define $V_p = 100\, V_{p100}$ km s^{-1}. But $V_p = \dot{r} - Hr = r/t(\alpha - \frac{2}{3})$, which gives on substitution

$$V_p = (9E_0 G)^{1/5}\left[\Omega_b\left(4\alpha^2 - 5\alpha + \frac{4}{3}\right) + \frac{2}{9}\left(\Omega_b - \frac{\Omega_b^2}{2}\right)\right]^{-1/5}\left(\alpha - \frac{2}{3}\right)\left[\frac{t_{\rm now}}{(1+z)^{3/2}}\right]^{\alpha - 1} , \tag{6.1}$$

where $t_{\rm now}$ is the age of the universe in seconds.

If we define the characteristic redshift of the epoch of the formation of galaxies as z_g and assume, as discussed in § V, that $V_p \propto (1 + z)$ for $(z \leq z_g)$, then for

$$V_p(z)|_{z=0} = 1 \times 10^7 V_{p100} \text{ cm s}^{-1} \tag{6.2}$$

we have

$$V_p|(z_g) = 1 \times 10^7 \times V_{p100} \times (1 + z_g) \text{ cm s}^{-1} , \tag{6.3}$$

which, along with equation (6.1), shows that to get a peculiar velocity $V_p(0) = 1 \times 10^7 V_{p100}(0)$ cm s^{-1} today, the needed energy of the explosion is given by

$$(E_0)^{1/5} = 1 \times 10^7 \times V_{p100} \frac{[\Omega_b(4\alpha^2 - 5\alpha + 4/3) + (\Omega_b - \Omega_b{}^2/2)]^{1/5}}{(9G)^{1/5}} \frac{1}{(\alpha - 2/3)} \frac{(1 + z_g)^{1+3/2(\alpha-1)}}{(t_{\rm now})^{\alpha-1}}, \tag{6.4}$$

and, in the era $z > z_g$,

$$r = \frac{1 \times 10^7}{(\alpha - 2/3)} (t_{\rm now}) V_{p100} \frac{(1 + z_g)^{1+3/2(\alpha-1)}}{(1 + z)^{3\alpha/2}}, \tag{6.5}$$

and

$$V_p(z) = 1 \times 10^7 (1 + z_g)^{1+3/2(\alpha-1)} (1 + z)^{-3/2(\alpha-1)} V_{p100}(0) . \tag{6.6}$$

Now if we write the thermal energy dissipated by the shock as $dE = (\frac{1}{2})(dm/dt)V_p{}^2\, dt$, where $(dm/dt) = 4\pi r^2 \rho_c \Omega_b V_p$ is the rate of sweeping of baryonic matter, then, assuming that the energy is released by Comptonization immediately as discussed in § II and redshifted by $(1 + z)$ between the time of release and the present, we can express the contribution to the microwave background as

$$\int_{z=\infty}^{z=z_g} \frac{dE}{(1 + z)} = \left(\frac{1}{2}\right) 4\pi \int \frac{r^2 \rho_c \Omega_b V_p{}^3}{(1 + z)}\, dt . \tag{6.7}$$

Using

$$\rho_c = \rho_{c,\,\rm now} \times (1 + z)^3$$

and equations (6.5) and (6.6), one gets

$$\int_{z=\infty}^{z=z_g} \frac{dE}{1 + z} = \frac{2\pi}{5} \times 10^{25} \times \rho_b \frac{(t_{\rm now})^3}{(\alpha - 2/3)^3} (1 + z_g)^{5/2} V_{p100}^2 . \tag{6.8}$$

Taking the mass of the shell as $4\pi/3 r^3 \Omega_b \rho_{z=z_g}$ at $z = z_g$, one can estimate the energy put into the present background per unit mass as

$$\frac{\int_{z=\infty}^{z=z_g} [dE/(1 + z)]}{(4\pi/3) r^3 \Omega_b \rho_{b_{z=z_g}}} = \frac{3}{10} \times 10^{14} (1 + z_g) \times V_{p100}^2 \text{ ergs g}^{-1} , \tag{6.9}$$

which is independent of Ω_b, h_0, and even α. If $z_g = 7$ then this equals $2.4 \times 10^{14} V_{p100}^2$ ergs g^{-1}.

If we now multiply this by $(0.3\eta \times \rho_{c,\,\rm now} \times \Omega_b)$, i.e., we assume that about 0.3η of the baryonic mass is in the galaxies at $z = z_g$, $\eta \approx 1$, then we can write the increase in the photon energy density of the microwave background as

$$\Delta U = 17.07 \times 10^{-17} (1 + z_g) V_{p100}^2 h_0{}^2 \Omega_b \eta \text{ ergs cm}^{-3} . \tag{6.10}$$

Taking the present energy density in the background as $U = (2.7)^4 \times 7.56 \times 10^{-15}$ ergs cm^{-3},

$$\frac{\Delta U}{U} = 4.21 \times 10^{-4} (1 + z_g) V_{p100}{}^2 h_0^2 \Omega_b \eta . \tag{6.11}$$

VII. THE RESULTING ANISOTROPY

Since the ratio $\Delta U/U$ can be written as $\Delta U/U \cong 4y$, where $y = \int (kT_e/m_e c^2) \sigma_e \eta_e\, dl$ is the Compton y parameter, then, since $\delta T/T \simeq -2y$ in the Rayleigh-Jeans part of the spectrum, $\delta T/T \simeq (\frac{1}{2})\Delta U/U$.

To obtain the variation of δT in a different direction, i.e., the anisotropy, we must determine the average number, n, of the shells crossed by the line of sight at $z = z_g$, which is $\sim N r^2/R^2$. Here N is the number of bubbles within the event horizon (which has radius R) at $z = z_g$ and r is the radius of each of the bubbles. Assuming a closely packed structure at present it can be shown that

$$n \sim \frac{10^2}{\sqrt{1 + z_g}} \frac{1}{V_{2000}}, \tag{7.1}$$

where we have defined $V = V_{2000} \times 2000$ km s^{-1} as the present radius of the holes in the velocity space (Davis *et al.* 1982). So the resulting anisotropy is $1/(N)^{1/2}(\delta T/T)^*$, where $(\delta T/T)^*$ is given by the earlier expression and

$$\delta T/T \cong 2.12 \times 10^{-5} V_{p100}^2 (1 + z_g)^{5/4} (V_{2000})^{1/2} h_0{}^2 \Omega_b \eta . \tag{7.2}$$

Given that $z_g \gtrsim 7$,

$$\delta T/T \gtrsim 2.85 \times 10^{-4} V_{p100}^2 (V_{2000})^{1/2} h_0{}^2 \Omega_b \eta .$$

The expected angular scale of δT is $\sim 1/n$.

If $V_p \cong 600$ km s^{-1}, $h_0 = \frac{1}{2}$, $\Omega_b = 0.1$, and $V_{2000} \simeq 1$, $z_g = 7$, then $\delta T/T \approx (2.5 \times 10^{-4})\eta$, which is marginally consistent with observed upper limits at an angular scale of 1°–3° (Partridge 1986) or at 4′ (Uson and Wilkinson 1985) and at 8°–10° (Davies *et al.* 1987).

The predicted anisotropy is extremely insensitive to the history and other details of the explosion, as reflected by the fact that they do not depend on $\alpha \equiv dlnr/dlnt$. One must assume that galaxies with high peculiar velocity formed not much earlier than $z \sim 7$. The explosive scenario would thus appear to be marginally viable from the point of view of microwave background constraints and the improved observations planned for the near future (Mather 1982) should be able to test the hypothesis of explosive energy input near a z of 10.

VIII. COMPARISON WITH GROWING COMPENSATED HOLES

The above calculation makes the assumption that the blast is sustained throughout sufficiently late periods and that Ω_b is sufficiently small for the dark matter which is passed by the baryonic matter shock to have not yet overtaken the latter. This puts the blast within the first stage of an $\Omega_b < \Omega = 1$ cosmological blast wave as termed by Ostriker and McKee (1988).

In the second stage, the dark matter merges with the baryonic shell (Bertschinger 1985*b*) and together they form a ridge such as discussed by Peebles (1982). The combined baryonic dark matter shell propagates approximately as an $r = \lambda t^{4/5}$ similarity solution. This is essentially a positive energy perturbation that is growing purely gravitationally but can also be considered the coasting phase of a cosmological blast wave.

The total energy to make the hole from smooth, zero-energy Hubble flow is

$$\Delta E = (47/144)MV^2 , \tag{8.1}$$

where M and V are, respectively, the mass of the shell, which now includes the dark matter, and the velocity.

The efficiency with which work can be done by the interior pressure on the shell while the hole is being initialized is readily evaluated if the energy in the interior is assumed to be a t^{Γ} power law up to a particular redshift z_* and is neglected for subsequent times. In this case the total work on the shell is

$$W(t_*) = \alpha(4\alpha^2 - 5\alpha + 14/9)MR^2/(5\alpha - 4)t_*^{\,2} \tag{8.2}$$

for $\alpha > \frac{4}{5}$. For $\alpha = \frac{4}{5}$, $E(t)$ is maintained a constant by resupplying the adiabatic losses. Hence W is logarithmically divergent if the initial radius is taken to vanish.

The inverse-redshift–weighted deposition into the microwave background by equation (6.7) is

$$dU = (3/10)(\alpha - 2/3)^2 MR^2/t_*^{\,2} . \tag{8.3}$$

The integrated radiative loss is

$$W_{\rm rad} = (3/10)(\alpha - 2/3)^3 MR^2/(\alpha - 4/5)t_*^{\,2} . \tag{8.4}$$

From equations (8.1), (8.2), and (8.4), it follows that the efficiency with which the energy source generates mechanical and potential energy is

$$\epsilon = 1 - (3/8)(\alpha - 2/3)^2/\alpha(\alpha - 7/12) \tag{8.5}$$

and that the ratio of inverse-redshift–weighted radiative loss to ϵW, the energy remaining in the matter, is

$$L = dU/\epsilon W = (3/10)(5\alpha - 4)(\alpha - 2/3)/(5\alpha^2/2 + 2\alpha - 3) . \tag{8.6}$$

Using equation (8.1), the energy per unit mass that is contributed to the present-day microwave background is

$$(47/144)LV^2/(1 + z_*) .$$

The contribution to the present-day microwave background is thus

$$\Delta U = (0.3)\eta V^2 \rho_c (47/144)L/(1 + z_*) . \tag{8.7}$$

Here we have neglected any distortion in the background in the coasting phase after the energy input phase and any internal energy of the baryonic–CDM mixture that composes the shell. For the sake of analytic tractability, we also assumed that the blast wave remains unfragmented till z_*.

If we take $\alpha = 1$ for $z > z_*$, $L = 1/15$. Choosing, as before, $V = (2000 \text{ km s}^{-1})\, V_{2000}$ and $h_0 = 1/2$ we find that

$$\Delta U/U = 2.92 \times 10^{-3} \eta V_{2000}^2/(1 + z_*)^{2/5} . \tag{8.8}$$

As before, defining n to be the average number of shells crossed by the line of sight at z_*, we have

$$n \approx \frac{10^2}{V_{2000}} (1 + z_*)^{-3/10}$$

and, the anisotropy,

$$\delta T/T = \frac{1}{2}\frac{1}{\sqrt{n}}\frac{\Delta U}{U} \approx 1.46 \times 10^{-4} (V_{2000})^{5/2} (1 + z_*)^{-1/4} \eta . \tag{8.9}$$

As $V_p = 0.16\, V$ for the self-similar phase of a growing compensated hole, the predicted distortion in the microwave background is comparable to the estimate made for the sustained blast wave (eq. [7.2]) if the z_g and z_* are taken to be about 10. Note that if the energization terminates at $z_* \gg 1/\Omega_b$ then the hole would have evolved into the second stage by the present time and would reduce the background distortion.

IX. FURTHER REMARKS

If the galaxies form before the shells evolve into the second stage, then a second issue that confronts the explosion scenario as an explanation of high peculiar velocities is the question of nonbaryonic dark matter in galaxies. In that case the peculiar velocities of galaxies at the time of their formation, $(1 + z_g)V_p$, is extremely high, and most nonbaryonic dark matter candidates are weakly interacting and nondissipative. These two points, especially when considered together, imply that galaxies should contain virtually no nonbaryonic matter. It is not yet clear that galaxies contain weakly interacting massive particles, but experiments for their detection can be designed in principle (Drukier and Stodolsky 1984; Goodman and Witten 1985; Wasserman 1986). It is interesting to note that finding the dynamically required dark matter in the form of low-mass stars (Gott 1981) would be consistent with the above picture of galaxy formation, but would raise Ω_b and, by the same factor, the estimate of $\delta T/T$ made in § VII. The difficulty may possibly be avoided if the galaxies form only after the dark matter has caught up with the baryonic shell or if the dark matter condenses into galaxy-halo–sized objects first and then cannibalizes the baryonic galaxies.

For completeness, we consider one more explosion scenario: suppose that the dark matter is shadow matter (Okun 1980; Blinnikov, Khlopov, and Yu 1983; Kolb, Seckel, and Turner 1985) that, while interacting weakly with normal matter, has enough interaction with itself to be dissipative. Explosions could take place in the shadow sector and influence the normal matter in a purely gravitational manner. The above mentioned difficulties are then removed because for gravitationally induced voids the peculiar velocities may increase with time (Peebles 1982), so galaxies' peculiar velocities at their formation can be minimal given the present-day values. Shock velocities in protogalactic material can be modest enough to permit line cooling, allowing late galaxy formation, and little energy is released in the cooling stage, so the energy can be put even into the microwave background without significantly perturbing the latter. The relative velocity between baryonic matter and shadow matter at the time of galaxy formation may easily be less than the escape velocity from galaxies, and both types are dissipative, so they may coexist in the same galaxy.

This scenario is similar from the point of view of ordinary matter to the more standard picture of gravitationally induced large-scale structure, except that the initial inhomogeneities can be manufactured much later than recombination. As such, the microwave background is unperturbed during the recombination era. It nevertheless may allow the formation of large voids at relatively late epochs.

This scenario is unattractive in that it is extremely difficult to verify or falsify. But it might be distinguished by its ability to create large voids without cleanly clearing out all of the normal matter in them, in contrast to a blast wave in normal matter which would sweep up all the free gas in its path.

X. SUMMARY

We have calculated the expected distortion of the microwave background from sustained explosions such as proposed by Ostriker, Thompson, and Witten (1986). If the explosions are sustained to a redshift of order Ω_b^{-1}, a lower limit can be expressed in terms of present-day quantities. Hence measurements of the anisotropy provide a test of such scenarios. We have been somewhat conservative in our assumptions, e.g., limb brightening has been ignored, and present-day voids were taken to be closely packed, so the dilution factor, $1/(n)^{1/2}$, that we use is probably an overestimate.

On the other hand, explosions can be hidden within the microwave background constraints if the energy input is limited to very high redshifts $z \gg \Omega_b^{-1}$, for then the growth of the hole is primarily due to its own gravitational deficit.

Our results indicate that if the energy input is sustained until $z \sim \Omega_b^{-1}$, and if galaxies form not much before then, the present limits are marginally consistent with typical peculiar velocities as large as 600 km s^{-1}.

We have noted the difficulty in explaining the presence of weakly interacting dark matter in the galaxies in case of formation of galaxies in the shells where the dark matter is yet to catch up. We have also seen that an explosion in a "shadow sector" of dark matter that interacts strongly with itself or letting the galaxies form only in the second stage of the shells should circumvent the difficulties.

This research was supported in part by NSF grant AST-86-11939. We thank J. Ostriker, P. J. E. Peebles, and J. Silk for useful discussions.

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