

Models of Universe with an inhomogeneous Big Bang singularity

II. CMBR dipole anisotropy as a byproduct of a conic Big-Bang singularity

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Abstract. The existence of stars and galaxies requires cosmological models with an inhomogeneous matter and radiation distribution. But in these models the initial singularity surface $t_0(r)$ is in general homogeneous (independent of r). In this second paper of a series devoted to an inhomogeneous Big Bang singularity, we investigate the cosmic microwave background radiation (CMBR) dipole. A special Tolman-Bondi Universe is used to study the effect of a Big-Bang singularity, depending linearly on r , on the CMBR anisotropy. It is shown that, for an observer located off the “center” of this Universe ($r = 0$), the parameters of the model can be tuned so as to reproduce, with a good approximation, the dipole and the quadrupole moments of the CMBR anisotropy observed in recent experiments. If the dipole should prove cosmological, a slight delaying of the Big-Bang over spatial coordinates would thus be a good candidate for its interpretation.

Key words: cosmology: cosmic microwave background – cosmology: theory

1. Introduction

The standard cosmological models rest mainly on a homogeneous and spherically symmetric Robertson-Walker metric $f_{\mu\nu}$, subject to the Einstein equations, with a matter-energy tensor depending only on the parameter (called the cosmic time¹) labelling the 3-surfaces.

The particular choice of one model lies in the choice of an equation of state. Inflationary cosmologies are part of this framework with peculiar choices of the equation of state imported from particle physics. Departures from this standard framework are developed for instance by the use of generalizations of the Einstein equations or by taking into account inhomogeneities in the matter distribution. Whereas there are presently no compelling reasons to abandon the Einstein equations, the introduction of inhomogeneities is unavoidable since there are large and small-scale structures in the observed Universe. These inhomogeneities induce inhomogeneities in the metric. They are used for the study of the formation of galaxies and large scale structures and of the fluctuations of the cosmic background at $3K$ (CMBR). In these models the metric $g_{\mu\nu}$ is a weakly perturbed Robertson-Walker metric: $g_{\mu\nu} = f_{\mu\nu} + h_{\mu\nu}$. In this framework the perturbation $h_{\mu\nu}$ is generally small and depends on the spatial coordinates, but the main term $f_{\mu\nu}$, characterized by a scale parameter $R(t)$ is independent of position. Therefore the cosmological singularity $R = 0$ is an equal time 3-surface; in astronomical terms, the age of the Universe is the same everywhere.

Solutions of the Einstein equations with inhomogeneous $R = 0$ hyper-surfaces have been studied analytically by Tolman (1934) and Bondi (1947). The Tolman-Bondi universes have applied to cosmological contexts such as clusters of galaxies (Tarentola, 1976) or the CMBR dipole (Paczynski & Piran, 1990).

But the parameters of the model are then set in such a way that they lead to a spatially homogeneous $R = 0$ Big Bang singular surface. There are no more reasons for this choice than for a strictly homogenous matter and radiation distribution. For instance, if the present universe were the result of the bounce from a collapse prior to the standard Big Bang (a solution which cannot be excluded), a strictly homogeneous $R = 0$ surface would result from a very unlikely fine tuning.

In a recent work (Célérier & Schneider, 1998), we have identified a class of inhomogeneous models of Universe, with a Big-Bang of “delayed” type, solving the standard horizon problem without need for an inflationary phase. In the present paper, we investigate the application of a peculiar model of this class to the CMBR anisotropy.

From a purely geometrical point of view, it is always possible to re-label the 3-surfaces so as to make the hyper-surface $R = 0$ independent of r . But such an arbitrary re-labelling is forbidden by the description of physical phenomena by the Schrödinger equation in curved space time. This equation, which gives the rate of evolution of phenomena, in particular, the thermal history of the Universe (through nucleosynthesis of light elements and matter-radiation decoupling), provides, in an inhomogeneous Universe, a clock imposing on time-like world lines a given time coordinate. Each hamiltonian used in the Schrödinger equation

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¹ Although this is an inappropriate denomination (Schneider, 1994)

gives a different time scale (e.g. the atomic transition rates) which depends on the local curvature. But there is an implicit postulate, that there is a fundamental time scale, the Planck time $\sqrt{Gh/c^3}$. Thus, the time coordinate can only be, with the choice of a given time unit, rescaled globally with a universal affine transformation $t \rightarrow at + b$ where the coefficient a and b are independent on r and t .

The dipole moment in the CMBR anisotropy is the most prominent feature in the recent observational data, as probed by the four years COBE experiments². It overcomes the quadrupole, of order 5.10^{-6} , by more than two orders of magnitude, its value being of order 10^{-3} (Smoot et al., 1992; Kogut et al., 1993).

This dipole is usually considered as resulting from a Doppler effect produced by our motion with respect to the CMBR rest-frame (Partridge, 1988). A few authors (Gunn, 1988; Paczynski & Piran, 1990; Turner, 1991; Langlois & Piran, 1996; Langlois, 1996), in the recent past, intended however to show that its origin could be in the large scale features of the Universe.

Paczynski & Piran (1990), using an ad hoc toy model, emphasized the possibility for the dipole to be generated by an entropy gradient in a Tolman-Bondi dust Universe. In the peculiar model they did study, they have assumed that the time of the Big-Bang was the same for all observers.

Hereafter we show that the dipole, and quadrupole, anisotropy, or part of it, could be considered as the outcome of a conic Big-Bang surface.

We first describe, in the following section, the special Tolman-Bondi model we use for our derivation. The calculations will be developed in Sect. 3 and the results exposed in Sect. 4. Our conclusions and a brief discussion are given in Sect. 5.

2. A flat dust spherically symmetrical model

We consider here the light cone emitted from the last scattering surface – temperature of order 4.10^3 K – towards the Earth at our present time. Since this period is matter dominated – the radiation was dynamically relevant only at times prior to a temperature of order 10^4 K – we are considering the behaviour of a photon gas immersed into an Universe satisfying

$$\rho_{\text{radiation}} \ll \rho_{\text{dust}}$$

that is, we neglect the radiation as source of gravitational field.

We then choose a Tolman-Bondi (Tolman, 1934; Bondi, 1947) model which figures out a dust (ideal non zero rest mass pressureless gas) dominated, spatially spherically symmetrical inhomogeneous Universe.

The Bondi line element, in co-moving coordinates and proper time, is:

$$ds^2 = -c^2 dt^2 + S^2(r, t) dr^2 + R^2(r, t) (d\theta^2 + \sin^2 \theta d\varphi^2) \quad (1)$$

² The three experiments aboard the COsmic Background Explorer satellite (COBE) are the Far-InfraRed Absolute Spectrophotometer (FIRAS) 60–630 GHz, the Differential Microwave Radiometers (DMR) 30–90 GHz and the Diffuse InfraRed Background Experiment (DIRBE) 1.2–240 μm . All experiments provide maps of small temperature fluctuations from an average 2.73° K for the CMBR.

It reduces to the usual Friedmann-Robertson-Walker metric for a homogeneous Universe.

Solving Einstein equations for this metric gives:

$$S^2(r, t) = \frac{R'^2(r, t)}{1 + 2E(r)/c^2} \quad (2)$$

$$\frac{1}{2} \dot{R}^2(r, t) - \frac{GM(r)}{R(r, t)} = E(r) \quad (3)$$

$$4\pi\rho(r, t) = \frac{M'(r)}{R'(r, t)R^2(r, t)} \quad (4)$$

a dot denoting differentiation with respect to t and a prime differentiation with respect to r .

$\rho(r, t)$ is the energy density ρ_{dust}

$E(r)$ and $M(r)$ are arbitrary functions of r . $E(r)$ can be interpreted as the total energy per unit mass and $M(r)$ as the baryonic mass within the sphere of co-moving radial coordinate r . $M(r)$ remaining constant with time, we use it to define the radial coordinate: $M(r) \equiv M_0 r^3$, where M_0 is a constant.

Eq. (3) can be solved and gives a parametric expression for $R(r, t)$ in case $E(r) \neq 0$ and an analytic one in case $E(r) = 0$.

We retain the flat Universe model $E(r) = 0$ and the analytic expression:

$$R(r, t) = [9GM(r)/2]^{1/3} [t - t_0(r)]^{2/3}$$

which, with the above definition for the radial coordinate, becomes:

$$R(r, t) = (9GM_0/2)^{1/3} r [t - t_0(r)]^{2/3} \quad (5)$$

The homogeneous limit of our model is the Einstein-de Sitter Universe with $\Omega = 1$.

$t_0(r)$ is another arbitrary function of r . It is the Big-Bang hyper-surface.

A class of models, identified as solving the horizon problem (C el erier & Schneider, 1998), exhibits a Big-Bang function of the form:

$$t_0(r) = br^n \quad b > 0, n > 0 \quad (6)$$

We here choose, for simplicity, to investigate the properties of the subclass:

$$t_0(r) = br \quad \text{with} \quad 1/R_H > b > 0 \quad (7)$$

This conic surface corresponds to perturbations with low ($k < 1/R_H$, R_H being the horizon radius) spatial frequencies.

Eq. (5) thus becomes:

$$R(r, t) = (9GM_0/2)^{1/3} r (t - br)^{2/3} \quad (8)$$

The specific entropy S is usually defined as the ratio of the number density of photons over the number density of baryons:

$$S \equiv \frac{k_B n_\gamma(r, t) m_b}{\rho(r, t)} \quad (9)$$

where m_b is the baryon mass and k_B , the Boltzmann constant.

In order to decouple the effect of inhomogeneous entropy distribution (as already studied by Paczynski & Piran, 1990) from an inhomogeneous $R = 0$ surface, we retain:

$$S = \text{const.}$$

The observed deviation of the CMBR from a perfect homogeneous pattern being very small, we can assume, as a reasonable approximation, thermodynamical equilibrium for the photons, so as to write, at the ultra-relativistic limit for bosons:

$$n_\gamma = a_n T^3 \quad (10)$$

T being the radiation temperature and $a_n = 2\zeta(3)k_B^3/[\pi^2(\hbar c)^3]$.

Letting, with no loss of generality, $S = \text{const.} = k_B \eta_0$, and taking the present photon to baryon density ratio η_0 to be $10^8/(2.66\Omega_b h_0^2)$, we derive the following expression for T :

$$T(r, t) = \left(\frac{10^8}{2.66h_0^2 2\pi G a_n m_b (3t - 5br)(t - br)} \right)^{1/3} \quad (11)$$

where h_0 is the Hubble constant in units 100 km/s/Mpc. Hereafter, for numerical applications, we take $h_0 = 0.75$.

3. Integration of the null geodesics and determination of the dipole and quadrupole moments

Let an observer, for example the COBE satellite, be located at (t_0, r_0) where the average temperature is T_0 , of order 2.7 K. The radial co-moving coordinate r_0 is chosen to be non zero so as to put the observer off the center of the Universe.

The light travelling from the last scattering surface to this observer follows null geodesics which we are going to numerically integrate, from the observer, until we reach this surface defined by its temperature $T_{\ell_s} = 4.10^3$ K.

In principle, one should integrate the optical depth equation along with the null geodesic equations. Here, we approximate the optical depth by a step-function. This procedure leads to integrate the null geodesics until the temperature reaches $T_{\ell_s} = 4.10^3$ K.

Our toy Universe being spherically symmetrical, an observer at a distance r_0 from the center sees an axially symmetrical Universe in the center direction. It is thus legitimate to integrate the geodesics in the meridional plane. The photons path is uniquely defined by the observer position (r_0, t_0) and the angle α between the direction from which comes the light ray as seen by the observer and the direction towards the center of the Universe.

In the following, we adopt the units:

$$c = 1, \quad 8\pi G/3 = 1 \quad \text{and} \quad M_0 = 1.$$

For the metric given by Eq. (1), the meridional plane is defined as:

$$\theta = \pi/2 \quad \sin \theta = 1 \quad k^\theta = 0$$

k^θ being the θ component of the photon wave-vector defined as:

$$k^\mu = -\frac{dx^\mu}{d\lambda}$$

which gives:

$$\frac{dt}{d\lambda} = -k^t \quad (12)$$

$$\frac{dr}{d\lambda} = \frac{k_r}{R'^2} = (16\pi/27)^{2/3} \frac{9(t-br)^{2/3}}{(3t-5br)^2} k_r \quad (13)$$

$$\frac{d\varphi}{d\lambda} = \frac{k_\varphi}{R^2} \quad (14)$$

From the geodesic equations of light:

$$dk^\mu/d\lambda + \Gamma_{\nu\lambda}^\mu k^\nu k^\lambda = 0$$

we obtain after some calculations:

$$\begin{aligned} \frac{dk^t}{d\lambda} = & -2 \left(\frac{16\pi}{27} \right)^{2/3} \left[\frac{3(3t-2br)}{(3t-5br)^3(t-br)^{1/3}} (k_r)^2 \right. \\ & \left. + \frac{1}{3r^2(t-br)^{7/3}} (k_\varphi)^2 \right] \end{aligned} \quad (15)$$

$$\begin{aligned} \frac{dk_r}{d\lambda} = & - \left(\frac{16\pi}{27} \right)^{2/3} \left[\frac{6b(2t-5br)}{(3t-5br)^3(t-br)^{1/3}} (k_r)^2 \right. \\ & \left. + \frac{3t-5br}{3r^3(t-br)^{7/3}} (k_\varphi)^2 \right] \end{aligned} \quad (16)$$

$$k_\varphi = \text{const.} \quad (17)$$

For photons: $ds^2 = 0$ coupled with Eqs. (12) to (14) gives:

$$(k^t)^2 = \left(\frac{k_r}{R'} \right)^2 + \left(\frac{k_\varphi}{R} \right)^2 \quad (18)$$

The equation for the redshift z_{ℓ_s} in co-moving coordinates is:

$$1 + z_{\ell_s} = \frac{(k^t)_{\ell_s}}{(k^t)_0}$$

$(k^t)_{\ell_s}$ and $(k^t)_0$ being the time-like component of the photons wave-vector at the last-scattering and at the observer respectively.

The former equations system can be integrated, the following initial conditions being given at the observer:

$$t = t_0 \quad r = r_0 \quad (k^t)_0 = 1 \quad (19)$$

And thus:

$$1 + z_{\ell_s} = (k^t)_{\ell_s} \quad (20)$$

At a given couple (t_0, r_0) corresponds values for R and its partial derivatives at t_0, r_0 .

We denote:

$$R_0 \equiv R(t_0, r_0) \quad \text{given by Eq. (8)}$$

$$R'_0 \equiv R'(t_0, r_0) \quad \text{and so on}$$

The observer at (t_0, r_0) seeing the photons trajectory making an angle α with the direction towards the center of the Universe, we can write:

$$(k_r)_0 = A \cos \alpha \quad (k_\varphi)_0 = B \sin \alpha$$

Substituting the former values of the coordinates of \mathbf{k}_0 into Eq. (18) written at (t_0, r_0) , we find:

$$A = R'_0 \quad B = R_0$$

And thus:

$$(k_r)_0 = R'_0 \cos \alpha \quad (k_\varphi)_0 = R_0 \sin \alpha \quad (21)$$

Eq. (17) becomes:

$$k_\varphi = R_0 \sin \alpha$$

Substituting in Eq. (18), we get:

$$(k_r)^2 = R'^2[(k^t)^2 - (R_0 \sin \alpha / R)^2]$$

which possesses two solutions:

$$k_r = \pm R'[(k^t)^2 - (R_0 \sin \alpha / R)^2]^{1/2} \quad (22)$$

From Eq. (4), with $M(r) \equiv r^3$, comes:

$$\rho_{\text{dust}} = (3/4\pi) \frac{r^2}{R'R^2}$$

As we want, for physical consistency, $\rho_{\text{dust}} \geq 0$, we get:

$$R' \geq 0$$

And because Eq. (13) implies the same sign for $dr/d\lambda$ and k_r , it follows that:

Eq. (22) with the plus sign is the solution $dr/d\lambda > 0$, where r is increasing with increasing λ parameter.

Eq. (22) with the minus sign is the solution $dr/d\lambda < 0$, where r is decreasing with increasing λ .

Substituting Eq. (22) into Eqs. (12) to (17), we get, after some calculations, the reduced system of three differential equations:

$$\frac{dt}{d\lambda} = -k^t \quad (23)$$

$$\frac{dr}{d\lambda} = \pm \left(\frac{16\pi}{27} \right)^{1/3} \frac{3(t-br)^{1/3}}{3t-5br} \left[(k^t)^2 - \frac{r_0^2(t-br_0)^{4/3} \sin^2 \alpha}{r^2(t-br)^{4/3}} \right]^{1/2} \quad (24)$$

$$\frac{dk^t}{d\lambda} = \frac{2(3t-2br)}{9(t-br)(3t-5br)} (k^t)^2 + \frac{2br_0^2(t-br_0)^{4/3} \sin^2 \alpha}{r(3t-5br)(t-br)^{7/3}} \quad (25)$$

Provided we choose the affine parameter λ increasing from $\lambda = 0$ at (t_0, r_0) to $\lambda = \lambda_{\ell_s}$ at (t_{ℓ_s}, r_{ℓ_s}) on the last scattering surface, we have to consider two cases:

- the “out-case”: the observer looks at a direction opposite to the center of the Universe ($\alpha > \pi/2$). We thus integrate the null geodesics from (t_0, r_0) to (t_{ℓ_s}, r_{ℓ_s}) with an always increasing r . We have to retain the plus sign in Eq. (24).

- the “in-case”: the observer looks at a light ray first approaching the center of the Universe, then moving away from it before reaching her eyes ($\alpha < \pi/2$). Eq. (24) with the minus sign first obtains until $dr/d\lambda = 0$, then the minus sign in Eq. (24) changes to a plus sign.

As, in the system of Eqs. (23) to (25), the dependence in α is of the form $\sin \alpha$ and as $\sin \alpha = \sin(\pi - \alpha)$, we can only discriminate between the “out” and “in” cases by the behaviour of the sign of $dr/d\lambda$.

We integrate a number of “in” and “out” null geodesics, each characterized by a value for α between zero and $\pi/2$, back in time from the observer at (t_0, r_0, T_0) until the temperature, as given by Eq. (11), reaches $T_{\ell_s} = (4/2.7)10^3 T_0$, which approximately defines the last scattering surface.

At this temperature, the redshift with respect to the observer, as given by Eq. (20), is $z_{\ell_s}^{\text{in-out}}(\alpha)$, somewhat varying, with the α angle and the “in” and “out” direction, about an average $z_{\ell_s}^{av}$.

The apparent temperature of the CMBR mesured in the α in-out direction is:

$$T_{\text{CMBR}}^{\text{in-out}}(\alpha) = \frac{T_{\ell_s}}{1 + z_{\ell_s}^{\text{in-out}}(\alpha)} = T_{\text{CMBR}}^{av} \frac{1 + z_{\ell_s}^{av}}{1 + z_{\ell_s}^{\text{in-out}}(\alpha)}$$

where the averages for T and z are calculated over the whole sky. We write with simplified notations:

$$\frac{T_{\text{CMBR}}}{T^{av}} = \frac{1 + z^{av}}{1 + z_{\ell_s}} \quad (26)$$

The CMBR temperature large scale inhomogeneities are expanded in spherical harmonics:

$$\frac{T_{\text{CMBR}}(\alpha, \varphi)}{T^{av}} = \sum_{\ell=1}^{\infty} \sum_{m=-\ell}^{+\ell} a_{\ell m} Y_{\ell m}(\alpha, \varphi)$$

α being the Euler angle usually called θ in spherical coordinates, and with:

$$a_{\ell m} = \int \frac{T_{\text{CMBR}}(\alpha, \varphi)}{T^{av}} Y_{\ell m}^*(\alpha, \varphi) \sin \alpha \, d\alpha \, d\varphi \quad (27)$$

The dipole and quadrupole moments are defined as:

$$D = (|a_{1-1}|^2 + |a_{10}|^2 + |a_{11}|^2)^{1/2}$$

$$Q = (|a_{2-2}|^2 + |a_{2-1}|^2 + |a_{20}|^2 + |a_{21}|^2 + |a_{22}|^2)^{1/2}$$

In the special case we are interested in, the large scale inhomogeneities only depend on the α angle so that all the $a_{\ell m}$ with $m \neq 0$ are zero.

The dipole and quadrupole moments thus reduce to:

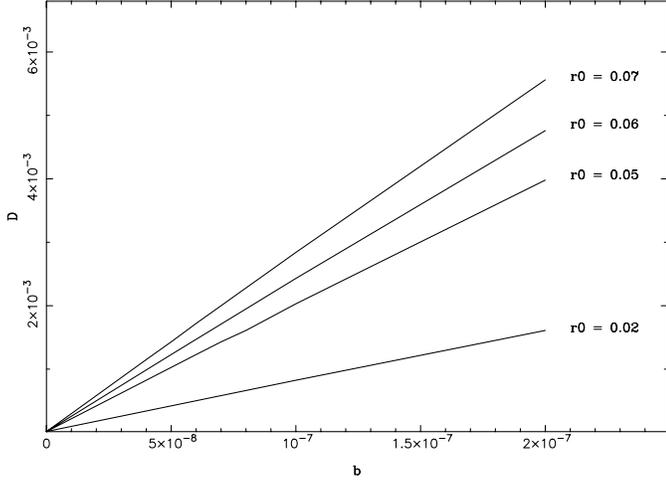
$$D = a_{10} \quad Q = a_{20}$$

a_{10} and a_{20} being given by Eq. (27) with:

$$Y_{10}(\alpha) = \sqrt{\frac{3}{4\pi}} \cos \alpha \quad Y_{20}(\alpha) = \sqrt{\frac{5}{4\pi}} \left(\frac{3}{2} \cos^2 \alpha - \frac{1}{2} \right)$$

It follows:

$$D = (1 + z^{av}) \int_0^\pi \frac{Y_{10}(\alpha)}{1 + z_{\ell_s}(\alpha)} \sin \alpha \, d\alpha \quad (28)$$


Fig. 1. The Dipole as a function of b for various r_0

$$Q = (1 + z^{av}) \int_0^\pi \frac{Y_{20}(\alpha)}{1 + z_{\ell_s}(\alpha)} \sin \alpha \, d\alpha \quad (29)$$

Taking into account Eq. (20) and the spherical symmetry of the model, we obtain:

$$D = \left| \frac{1}{2} \sqrt{\frac{3}{\pi}} k_{av}^t \left[\int_0^{\frac{\pi}{2}} \frac{\cos \alpha \sin \alpha}{k_{in}^t(\alpha)} d\alpha - \int_0^{\frac{\pi}{2}} \frac{\cos \alpha \sin \alpha}{k_{out}^t(\alpha)} d\alpha \right] \right| \quad (30)$$

$$Q = \frac{1}{4} \sqrt{\frac{5}{\pi}} k_{av}^t \left[\int_0^{\frac{\pi}{2}} \frac{(3 \cos^2 \alpha - 1) \sin \alpha}{k_{in}^t(\alpha)} d\alpha + \int_0^{\frac{\pi}{2}} \frac{(3 \cos^2 \alpha - 1) \sin \alpha}{k_{out}^t(\alpha)} d\alpha \right] \quad (31)$$

4. Results

We have first numerically integrated a number of ‘‘out’’ and ‘‘in’’ null geodesics, for various r_0 and b , and for values of α going from 0 to $\frac{\pi}{2}$, with t_0 corresponding to $T_0 = 2.7 \text{ K}$ in Eq. (11).

We have then calculated the dipole and quadrupole moments D and Q , according to Eqs. (30) and (31).

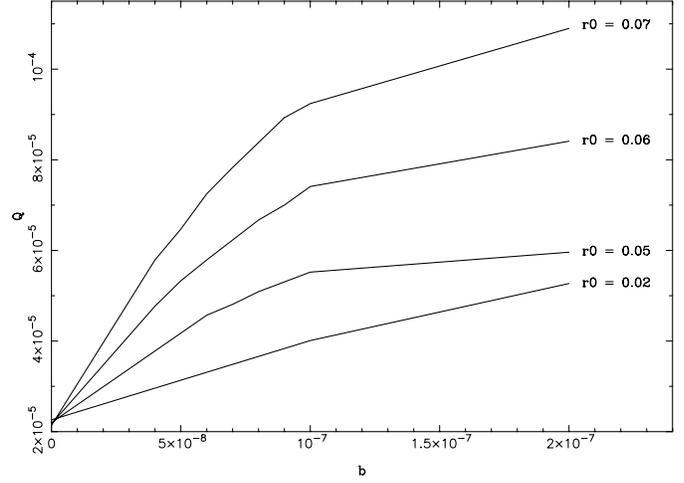
We have selected the values of the doublets leading to D and Q approaching the observed values³ $D \sim 10^{-3}$, $Q \sim 10^{-5}$. These results are given in Figs. 1 and 2.

The best fitted values of r_0 and b giving D close to 10^{-3} and Q close to 10^{-5} are shown in Table 1.

5. Conclusion and discussion

Using a toy model, chosen within the class of delayed Big-Bang models identified as solving the horizon problem without need for any inflationary phase (C el erier & Schneider, 1998), and presenting the following main features:

³ This choice will be discussed in Sect. 5.


Fig. 2. The Quadrupole as a function of b for various r_0
Table 1. Best fitted values of r_0 and b

r_0	b	D	Q
0.02	$2 \cdot 10^{-7}$	$1.61 \cdot 10^{-3}$	$5.27 \cdot 10^{-5}$
0.03	$9 \cdot 10^{-8}$	$1.11 \cdot 10^{-3}$	$3.70 \cdot 10^{-5}$
0.04	$7 \cdot 10^{-8}$	$1.15 \cdot 10^{-3}$	$3.99 \cdot 10^{-5}$
0.05	$6 \cdot 10^{-8}$	$1.23 \cdot 10^{-3}$	$4.57 \cdot 10^{-5}$
0.06	$5 \cdot 10^{-8}$	$1.23 \cdot 10^{-3}$	$5.33 \cdot 10^{-5}$
0.07	$4 \cdot 10^{-8}$	$1.15 \cdot 10^{-3}$	$5.79 \cdot 10^{-5}$

- dust dominated spherically symmetrical Tolman-Bondi Universe
- conic Big-Bang singularity
- observer located off the center of the Universe

we showed that can be found values for the parameters of the model – the location of the observer in space-time and the increasing rate of the Big-Bang function – that allow to somehow reproduce the observed dipole and quadrupole moments in the CMBR anisotropy.

This provides a new possible interpretation of the dipole (or part of it, as it is obvious that there is probably a Doppler component due to the local motion of the Galaxy with respect to the CMBR rest frame).

As has been stressed by other authors (Paczynski & Piran, 1990; Turner, 1991; Langlois & Piran, 1996; Langlois, 1996), there are various observational ways to discriminate between a local and a cosmological origin for the dipole.

From an analysis of a sparse-sampled redshift survey of IRAS Point Source Catalog $60 \mu\text{m}$ sources, performed with the tools of standard cosmology, Rowan-Robinson et al. (1990) concluded, for instance, that the peculiar velocity of the Local Group should be $579 \Omega_0^{0.6} \text{ km s}^{-1}$ towards $(l, b) = (269.5, 29.8)$.

For $\Omega_0 \sim 0.3$, this would give a velocity of order 280 km s^{-1} , to be compared to the CMBR dipole velocity: $600 \pm 50 \text{ km s}^{-1}$ (Partridge, 1988) towards $(l, b) = (124.7 \pm 0.8, 48.2 \pm 0.5)$ (Smoot et al., 1992).

In this framework, the local component of the dipole would be of order 50% of the total dipole.

If, from future analyses of observational data, part of the dipole was confirmed to appear non Doppler, other work, connected in particular with multipole moments of higher order, would be needed to discriminate between the various cosmological candidate interpretations.

It has to be stressed that, if the inhomogeneous Big Bang assumption is thus retained, a 50% shift in the dipole cosmological component would not significantly affect the results given in above Table 1.

In our formerly cited paper (Célérier & Schneider, 1998), we have shown that the horizon problem can be solved by means of a delayed Big-Bang, provided the observer is located near the center of a spherically symmetrical Universe. Work is in progress to extend these results to models with an observer arbitrarily situated off the center.

Another interesting feature of the here presented work is to show that, in a model of the above studied class, the occurrence of a cosmological component of the dipole implies a relation between the location r_0 of the observer and the slope b of the Big-Bang function.

It can be seen, from Table 1, that the larger r_0 , the smaller b , and this yields a selection within the parameters space of the conic Big-Bang models.

We conjecture that such a feature pertains to any subclass of the delayed Big-Bang models.

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