

ADVANCED TOPICS IN COSMOLOGY: A PEDAGOGICAL INTRODUCTION

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These lecture notes provide a concise, rapid and pedagogical introduction to several advanced topics in contemporary cosmology. The discussion of thermal history of the universe, linear perturbation theory, theory of CMBR temperature anisotropies and the inflationary generation of perturbation are presented in a manner accessible to someone who has done a first course in cosmology. The discussion of dark energy is more research oriented and reflects the personal bias of the author. Contents: (I) The cosmological paradigm and Friedmann model; (II) Thermal history of the universe; (III) Structure formation and linear perturbation theories; (IV) Perturbations in dark matter and radiation; (V) Transfer function for matter perturbations; (VI) Temperature anisotropies of CMBR; (VII) Generation of initial perturbations from inflation; (VIII) The dark energy.

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I. THE COSMOLOGICAL PARADIGM AND FRIEDMANN MODEL

Observations show that the universe is fairly homogeneous and isotropic at scales larger than about $150h^{-1}$ Mpc where $1 \text{ Mpc} \simeq 3 \times 10^{24} \text{ cm}$ and $h \approx 0.7$ is a parameter related to the expansion rate of the universe. The conventional — and highly successful — approach to cosmology separates the study of large scale ($l \gtrsim 150h^{-1} \text{ Mpc}$) dynamics of

the universe from the issue of structure formation at smaller scales. The former is modeled by a homogeneous and isotropic distribution of energy density; the latter issue is addressed in terms of gravitational instability which will amplify the small perturbations in the energy density, leading to the formation of structures like galaxies. In such an approach, the expansion of the background universe is described by the metric (We shall use units with $c = 1$ throughout, unless otherwise specified):

$$ds^2 \equiv dt^2 - a^2 d\mathbf{x}^2 \equiv dt^2 - a^2(t) [d\chi^2 + S_k^2(\chi) (d\theta^2 + \sin^2 \theta d\phi^2)] \quad (1)$$

with $S_k(\chi) = (\sin \chi, \chi, \sinh \chi)$ for the three values of the label $k = (1, 0, -1)$. The function $a(t)$ is governed by the equations:

$$\frac{\dot{a}^2 + k}{a^2} = \frac{8\pi G\rho}{3}; \quad d(\rho a^3) = -p da^3 \quad (2)$$

The first one relates expansion rate of the universe to the energy density ρ and $k = 0, \pm 1$ is a parameter which characterizes the spatial curvature of the universe. The second equation, when coupled with the equation of state $p = p(\rho)$ which relates the pressure p to the energy density, determines the evolution of energy density $\rho = \rho(a)$ in terms of the expansion factor of the universe. In particular if $p = w\rho$ with (at least, approximately) constant w then, $\rho \propto a^{-3(1+w)}$ and (if we further assume $k = 0$, which is strongly favoured by observations) the first equation in Eq.(2) gives $a \propto t^{2/[3(1+w)]}$. We will also often use the redshift $z(t)$, defined as $(1+z) = a_0/a(t)$ where the subscript zero denotes quantities evaluated at the present moment. In a $k = 0$ universe, we can set $a_0 = 1$ by rescaling the spatial coordinates.

It is convenient to measure the energy densities of different components in terms of a *critical energy density* (ρ_c) required to make $k = 0$ at the present epoch. (Of course, since k is a constant, it will remain zero at all epochs if it is zero at any given moment of time.) From Eq.(2), it is clear that $\rho_c = 3H_0^2/8\pi G$ where $H_0 \equiv (\dot{a}/a)_0$ — called the Hubble constant — is the rate of expansion of the universe at present. Numerically

$$\begin{aligned} \rho_c &= \frac{3H_0^2}{8\pi G} = 1.88h^2 \times 10^{-29} \text{ gm cm}^{-3} = 2.8 \times 10^{11} h^2 M_\odot \text{ Mpc}^{-3} \\ &= 1.1 \times 10^4 h^2 \text{ eV cm}^{-3} = 1.1 \times 10^{-5} h^2 \text{ protons cm}^{-3} \end{aligned} \quad (3)$$

The variables $\Omega_i \equiv \rho_i/\rho_c$ will give the fractional contribution of different components of the universe (i denoting baryons, dark matter, radiation, etc.) to the critical density. Observations then lead to the following results:

(1) Our universe has $0.98 \lesssim \Omega_{tot} \lesssim 1.08$. The value of Ω_{tot} can be determined from the angular anisotropy spectrum of the cosmic microwave background radiation (CMBR; see Section VI) and these observations (combined with the reasonable assumption that $h > 0.5$) show[1] that we live in a universe with critical density, so that $k = 0$.

(2) Observations of primordial deuterium produced in big bang nucleosynthesis (which took place when the universe was about few minutes in age) as well as the CMBR observations show[2] that the *total* amount of baryons in the universe contributes about $\Omega_B = (0.024 \pm 0.0012)h^{-2}$. Given the independent observations[3] which fix $h = 0.72 \pm 0.07$, we conclude that $\Omega_B \cong 0.04 - 0.06$. These observations take into account all baryons which exist in the universe today irrespective of whether they are luminous or not. *Combined with previous item we conclude that most of the universe is non-baryonic.*

(3) Host of observations related to large scale structure and dynamics (rotation curves of galaxies, estimate of cluster masses, gravitational lensing, galaxy surveys ..) all suggest[4] that the universe is populated by a non-luminous component of matter (dark matter; DM hereafter) made of weakly interacting massive particles which *does* cluster at galactic scales. This component contributes about $\Omega_{DM} \cong 0.20 - 0.35$ and has the simple equation of state $p_{DM} \approx 0$. The second equation in Eq.(2), then gives $\rho_{DM} \propto a^{-3}$ as the universe expands which arises from the evolution of number density of particles: $\rho = nmc^2 \propto n \propto a^{-3}$.

(4) Combining the last observation with the first we conclude that there must be (at least) one more component to the energy density of the universe contributing about 70% of critical density. Early analysis of several observations[5] indicated that this component is unclustered and has negative pressure. This is confirmed dramatically by the supernova observations (see Ref. [6]; for a critical look at the current data, see Ref. [7]). The observations suggest that the missing component has $w = p/\rho \lesssim -0.78$ and contributes $\Omega_{DE} \cong 0.60 - 0.75$. The simplest choice for such *dark energy* with negative pressure is the cosmological constant which is a term that can be added to Einstein's equations. This term acts like a fluid with an equation of state $p_{DE} = -\rho_{DE}$; the second equation in Eq.(2), then gives $\rho_{DE} = \text{constant}$ as universe expands.

(5) The universe also contains radiation contributing an energy density $\Omega_R h^2 = 2.56 \times 10^{-5}$ today most of which is due to photons in the CMBR. The equation of state is $p_R = (1/3)\rho_R$; the second equation in Eq.(2), then gives $\rho_R \propto a^{-4}$. Combining it with the result $\rho_R \propto T^4$ for thermal radiation, it follows that $T \propto a^{-1}$. Radiation is

dynamically irrelevant today but since $(\rho_R/\rho_{DM}) \propto a^{-1}$ it would have been the dominant component when the universe was smaller by a factor larger than $\Omega_{DM}/\Omega_R \simeq 4 \times 10^4 \Omega_{DM} h^2$.

(6) Taking all the above observations together, we conclude that our universe has (approximately) $\Omega_{DE} \simeq 0.7, \Omega_{DM} \simeq 0.26, \Omega_B \simeq 0.04, \Omega_R \simeq 5 \times 10^{-5}$. All known observations are consistent with such an — admittedly weird — composition for the universe.

Using $\rho_{NR} \propto a^{-3}, \rho_R \propto a^{-4}$ and $\rho_{DE} = \text{constant}$ we can write Eq.(2) in a convenient dimensionless form as

$$\frac{1}{2} \left(\frac{dq}{d\tau} \right)^2 + V(q) = E \quad (4)$$

where $\tau = H_0 t, a = a_0 q(\tau), \Omega_{NR} = \Omega_B + \Omega_{DM}$ and

$$V(q) = -\frac{1}{2} \left[\frac{\Omega_R}{q^2} + \frac{\Omega_{NR}}{q} + \Omega_{DE} q^2 \right]; \quad E = \frac{1}{2} (1 - \Omega_{\text{tot}}). \quad (5)$$

This equation has the structure of the first integral for motion of a particle with energy E in a potential $V(q)$. For models with $\Omega = \Omega_{NR} + \Omega_{DE} = 1$, we can take $E = 0$ so that $(dq/d\tau) = \sqrt{V(q)}$. Based on the observed composition of the universe, we can identify three distinct phases in the evolution of the universe when the temperature is less than about 100 GeV. At high redshifts (small q) the universe is radiation dominated and \dot{q} is independent of the other cosmological parameters. Then Eq.(4) can be easily integrated to give $a(t) \propto t^{1/2}$ and the temperature of the universe decreases as $T \propto t^{-1/2}$. As the universe expands, a time will come when ($t = t_{\text{eq}}, a = a_{\text{eq}}$ and $z = z_{\text{eq}}$, say) the matter energy density will be comparable to radiation energy density. For the parameters described above, $(1 + z_{\text{eq}}) = \Omega_{NR}/\Omega_R \simeq 4 \times 10^4 \Omega_{DM} h^2$. At lower redshifts, matter will dominate over radiation and we will have $a \propto t^{2/3}$ until fairly late when the dark energy density will dominate over non relativistic matter. This occurs at a redshift of z_{DE} where $(1 + z_{\text{DE}}) = (\Omega_{DE}/\Omega_{NR})^{1/3}$. For $\Omega_{DE} \approx 0.7, \Omega_{NR} \approx 0.3$, this occurs at $z_{\text{DE}} \approx 0.33$. In this phase, the velocity \dot{q} changes from being a decreasing function to an increasing function leading to an accelerating universe. In addition to these, we believe that the universe probably went through a rapidly expanding, inflationary, phase very early when $T \approx 10^{14}$ GeV; we will say more about this in Section VII. (For a textbook description of these and related issues, see e.g. Ref. [8].)

Before we conclude this section, we will briefly mention some key aspects of the background cosmology described by a Friedmann model.

(a) The metric in Eq.(1) can be rewritten using the expansion parameter a or the redshift $z = (a_0/a)^{-1} - 1$ as the time coordinate in the form

$$ds^2 = H^{-2}(a) \left(\frac{da}{a} \right)^2 - a^2 dx^2 = \frac{1}{(1+z)^2} [H^{-2}(z) dz^2 - dx^2] \quad (6)$$

This form clearly shows that the only dynamical content of the metric is encoded in the function $H(a) = (\dot{a}/a)$. An immediate consequence is that any observation which is capable of determining the geometry of the universe can only provide — at best — information about this function.

(b) Since cosmological observations usually use radiation received from distant sources, it is worth reviewing briefly the propagation of radiation in the universe. The radial light rays follow a trajectory given by

$$r_{\text{em}}(z) = S_k(\alpha); \quad \alpha \equiv \frac{1}{a_0} \int_0^z H^{-1}(z) dz \quad (7)$$

if the photon is emitted at r_{em} at the redshift z and received here today. Two other quantities closely related to $r_{\text{em}}(z)$ are the luminosity distance, d_L , and the angular diameter distance d_A . If we receive a flux F from a source of luminosity L , then the luminosity distance is defined via the relation $F \equiv L/4\pi d_L^2(z)$. If an object of transverse length l subtends a small angle θ , the angular diameter distance is defined via ($l = \theta d_A$). Simple calculation shows that:

$$d_L(z) = a_0 r_{\text{em}}(z)(1+z) = a_0(1+z)S_k(\alpha); \quad d_A(z) = a_0 r_{\text{em}}(z)(1+z)^{-1} \quad (8)$$

(c) As an example of determining the spacetime geometry of the universe from observations, let us consider how one can determine $a(t)$ from the observations of the luminosity distance. It is clear from the first equation in Eq. (8) that

$$H^{-1}(z) = \left[1 - \frac{k d_L^2(z)}{a_0^2 (1+z)^2} \right]^{-1/2} \frac{d}{dz} \left[\frac{d_L(z)}{1+z} \right] \rightarrow \frac{d}{dz} \left[\frac{d_L(z)}{1+z} \right] \quad (9)$$

where the last form is valid for a $k = 0$ universe. If we determine the form of $d_L(z)$ from observations — which can be done if we can measure the flux F from a class of sources with known value for luminosity L — then we can use this relation to determine the evolutionary history of the universe and thus the dynamics.

II. THERMAL HISTORY OF THE UNIVERSE

Let us next consider some key events in the evolutionary history of our universe [8]. The most well understood phase of the universe occurs when the temperature is less than about 10^{12} K. Above this temperature, thermal production of baryons and their strong interaction is significant and somewhat difficult to model. We can ignore such complications at lower temperatures and — as we shall see — several interesting physical phenomena did take place during the later epochs with $T \lesssim 10^{12}$.

The first thing we need to do is to determine the composition of the universe when $T \approx 10^{12}$ K. We will certainly have, at this time, copious amount of photons and all species of neutrinos and antineutrinos. In addition, neutrons and protons must exist at this time since there is no way they could be produced later on. (This implies that phenomena which took place at higher temperatures should have left a small excess of baryons over anti baryons; we do not quite understand how this happened and will just take it as an initial condition.) Since the rest mass of electrons correspond to a much lower temperature (about 0.5×10^{10} K), there will be large number of electrons and positrons at this temperature but in order to maintain charge neutrality, we need to have a slight excess of electrons over positrons (by about 1 part in 10^9) with the net negative charge compensating the positive charge contributed by protons.

An elementary calculation using the known interaction rates show that all these particles are in thermal equilibrium at this epoch. Hence standard rules of statistical mechanics allows us to determine the number density (n), energy density (ρ) and the pressure (p) in terms of the distribution function f :

$$n = \int f(\mathbf{k}) d^3\mathbf{k} = \frac{g}{2\pi^2} \int_m^\infty \frac{(E^2 - m^2)^{1/2} E dE}{\exp[(E - \mu)/T] \pm 1} \quad (10)$$

$$\rho = \int E f(\mathbf{k}) d^3\mathbf{k} = \frac{g}{2\pi^2} \int_m^\infty \frac{(E^2 - m^2)^{1/2} E^2 dE}{\exp[(E - \mu)/T] \pm 1} \quad (11)$$

$$p = \frac{1}{3} \int d^3\mathbf{k} f(\mathbf{k}) k v(\mathbf{k}) = \int \frac{1}{3} \frac{|k|^2}{E} f(\mathbf{k}) d^3\mathbf{k} = \frac{g}{6\pi^2} \int_m^\infty \frac{(E^2 - m^2)^{3/2} dE}{e^{[(E - \mu)/T]} \pm 1} \quad (12)$$

Next, we can argue that the chemical potentials for electrons, positrons and neutrinos can be taken to be zero. For example, conservation of chemical potential in the reaction $e^+e^- \rightarrow 2\gamma$ implies that the chemical potentials of electrons and positrons must differ in a sign. But since the number densities of electrons and positrons, which are determined by the chemical potential, are very close to each other, the chemical potentials of electrons and positrons must be (very closely) equal to each other. Hence both must be (very close to) zero. Similar reasoning based on lepton number shows that neutrinos should also have zero chemical potential. Given this, one can evaluate the integrals for all the relativistic species and we obtain for the total energy density

$$\rho_{\text{total}} = \sum_{i=\text{boson}} g_i \left(\frac{\pi^2}{30}\right) T_i^4 + \sum_{i=\text{fermion}} \frac{7}{8} g_i \left(\frac{\pi^2}{30}\right) T_i^4 = g_{\text{total}} \left(\frac{\pi^2}{30}\right) T^4 \quad (13)$$

where

$$g_{\text{total}} \equiv \sum_{\text{boson}} g_B + \sum_{\text{fermion}} \frac{7}{8} g_F. \quad (14)$$

The corresponding entropy density is given by

$$s \cong \frac{1}{T}(\rho + p) = \frac{2\pi^2}{45} q T^3; \quad q \equiv q_{\text{total}} = \sum_{\text{boson}} g_B + \frac{7}{8} \sum_{\text{fermion}} g_F. \quad (15)$$

A. Neutrino background

As a simple application of the above result, let us consider the fate of neutrinos in the expanding universe. From the standard weak interaction theory, one can compute the reaction rate Γ of the neutrinos with the rest of the species. When this reaction rate fall below the expansion rate H of the universe, the reactions cannot keep the neutrinos coupled to the rest of the matter. A simple calculation [8] shows that the relevant ratio is given by

$$\frac{\Gamma}{H} \simeq \left(\frac{T}{1.4\text{MeV}}\right)^3 = \left(\frac{T}{1.6 \times 10^{10}\text{K}}\right)^3 \quad (16)$$

Thus, for $T \lesssim 1.6 \times 10^{10}$ K, the neutrinos decouple from matter. At slightly lower temperature, the electrons and positrons annihilate increasing the number density of photons. Neutrinos do not get any share of this energy since they have already decoupled from the rest of the matter. As a result, the photon temperature goes up with respect to the neutrino temperature once the e^+e^- annihilation is complete. This increase in the temperature is easy to calculate. As far as the photons are concerned, the increase in the temperature is essentially due to the change in the degrees of freedom g and is given by:

$$\frac{(aT_\gamma)_{\text{after}}^3}{(aT_\gamma)_{\text{before}}^3} = \frac{g_{\text{before}}}{g_{\text{after}}} = \frac{\frac{7}{8}(2+2)+2}{2} = \frac{11}{4}. \quad (17)$$

(In the numerator, one 2 is for electron; one 2 is for positron; the 7/8 factor arises because these are fermions. The final 2 is for photons. In the denominator, there are only photons to take care of.) Therefore

$$\begin{aligned} (aT_\gamma)_{\text{after}} &= \left(\frac{11}{4}\right)^{1/3} (aT_\gamma)_{\text{before}} = \left(\frac{11}{4}\right)^{1/3} (aT_\nu)_{\text{before}} \\ &= \left(\frac{11}{4}\right)^{1/3} (aT_\nu)_{\text{after}} \simeq 1.4(aT_\nu)_{\text{after}}. \end{aligned} \quad (18)$$

The first equality is from Eq. (17); the second arises because the photons and neutrinos had the same temperature originally; the third equality is from the fact that for decoupled neutrinos aT_ν is a constant. This result leads to the prediction that, at present, the universe will contain a bath of neutrinos which has temperature that is (predictably) lower than that of CMBR. The future detection of such a cosmic neutrino background will allow us to probe the universe at its earliest epochs.

B. Primordial Nucleosynthesis

When the temperature of the universe is higher than the binding energy of the nuclei (\sim MeV), none of the heavy elements (helium and the metals) could have existed in the universe. The binding energies of the first four light nuclei, ${}^2\text{H}$, ${}^3\text{H}$, ${}^3\text{He}$ and ${}^4\text{He}$ are 2.22 MeV, 6.92 MeV, 7.72 MeV and 28.3 MeV respectively. This would suggest that these nuclei could be formed when the temperature of the universe is in the range of (1 – 30) MeV. The actual synthesis takes place only at a much lower temperature, $T_{\text{nuc}} = T_n \simeq 0.1$ MeV. The main reason for this delay is the ‘high entropy’ of our universe, i.e., the high value for the photon-to-baryon ratio, η^{-1} . Numerically,

$$\eta = \frac{n_B}{n_\gamma} = 5.5 \times 10^{-10} \left(\frac{\Omega_B h^2}{0.02}\right); \quad \Omega_B h^2 = 3.65 \times 10^{-3} \left(\frac{T_0}{2.73 \text{ K}}\right)^3 \eta_{10} \quad (19)$$

To see this, let us assume, for a moment, that the nuclear (and other) reactions are fast enough to maintain thermal equilibrium between various species of particles and nuclei. In thermal equilibrium, the number density of a nuclear species ${}^A N_Z$ with atomic mass A and charge Z will be

$$n_A = g_A \left(\frac{m_A T}{2\pi}\right)^{3/2} \exp\left[-\left(\frac{m_A - \mu_A}{T}\right)\right]. \quad (20)$$

From this one can obtain the equation for the temperature T_A at which the mass fraction of a particular species-A will be of order unity ($X_A \simeq 1$). We find that

$$T_A \simeq \frac{B_A/(A-1)}{\ln(\eta^{-1}) + 1.5 \ln(m_B/T)} \quad (21)$$

where B_A is the binding energy of the species. This temperature will be fairly lower than B_A because of the large value of η^{-1} . For ${}^2\text{H}$, ${}^3\text{He}$ and ${}^4\text{He}$ the value of T_A is 0.07 MeV, 0.11 MeV and 0.28 MeV respectively. Comparison with the binding energy of these nuclei shows that these values are lower than the corresponding binding energies B_A by a factor of about 10, at least.

Thus, even when the thermal equilibrium is maintained, significant synthesis of nuclei can occur only at $T \lesssim 0.3$ MeV and not at higher temperatures. If such is the case, then we would expect significant production ($X_A \lesssim 1$) of nuclear species-A at temperatures $T \lesssim T_A$. It turns out, however, that the rate of nuclear reactions is *not* high enough to maintain thermal equilibrium between various species. We have to determine the temperatures up to which thermal

equilibrium can be maintained and redo the calculations to find non-equilibrium mass fractions. The general procedure for studying non equilibrium abundances in an expanding universe is based on *rate equations*. Since we will require this formalism again in Section II C (for the study of recombination), we will develop it in a somewhat general context.

Consider a reaction in which two particles 1 and 2 interact to form two other particles 3 and 4. For example, $n + \nu_e \rightleftharpoons p + e$ constitutes one such reaction which converts neutrons into protons in the forward direction and protons into neutrons in the reverse direction; another example we will come across in the next section is $p + e \rightleftharpoons \text{H} + \gamma$ where the forward reaction describes recombination of electron and proton forming a neutral hydrogen atom (with the emission of a photon), while the reverse reaction is the photoionisation of a hydrogen atom. In general, we are interested in how the number density n_1 of particle species 1, say, changes due to a reaction of the form $1 + 2 \rightleftharpoons 3 + 4$.

We first note that even if there is no reaction, the number density will change as $n_1 \propto a^{-3}$ due to the expansion of the universe; so what we are really after is the change in $n_1 a^3$. Further, the forward reaction will be proportional to the product of the number densities $n_1 n_2$ while the reverse reaction will be proportional to $n_3 n_4$. Hence we can write an equation for the rate of change of particle species n_1 as

$$\frac{1}{a^3} \frac{d(n_1 a^3)}{dt} = \mu (A n_3 n_4 - n_1 n_2). \quad (22)$$

The left hand side is the relevant rate of change over and above that due to the expansion of the universe; on the right hand side, the two proportionality constants have been written as μ and $(A\mu)$, both of which, of course, will be functions of time. (The quantity μ has the dimensions of $\text{cm}^3 \text{s}^{-1}$, so that $n\mu$ has the dimensions of s^{-1} ; usually $\mu \simeq \sigma v$ where σ is the cross-section for the relevant process and v is the relative velocity.) The left hand side has to vanish when the system is in thermal equilibrium with $n_i = n_i^{\text{eq}}$, where the superscript ‘eq’ denotes the equilibrium densities for the different species labeled by $i = 1 - 4$. This condition allows us to rewrite A as $A = n_1^{\text{eq}} n_2^{\text{eq}} / (n_3^{\text{eq}} n_4^{\text{eq}})$. Hence the rate equation becomes

$$\frac{1}{a^3} \frac{d(n_1 a^3)}{dt} = \mu n_1^{\text{eq}} n_2^{\text{eq}} \left(\frac{n_3 n_4}{n_3^{\text{eq}} n_4^{\text{eq}}} - \frac{n_1 n_2}{n_1^{\text{eq}} n_2^{\text{eq}}} \right). \quad (23)$$

In the left hand side, one can write $(d/dt) = Ha(d/da)$ which shows that the relevant time scale governing the process is H^{-1} . Clearly, when $H/n\mu \gg 1$ the right hand side becomes ineffective because of the (μ/H) factor and the number of particles of species 1 does not change. We see that when the expansion rate of the universe is large compared to the reaction rate, the given reaction is ineffective in changing the number of particles. This certainly does *not* mean that the reactions have reached thermal equilibrium and $n_i = n_i^{\text{eq}}$; in fact, it means exactly the opposite: The reactions are not fast enough to drive the number densities towards equilibrium densities and the number densities “freeze out” at non equilibrium values. Of course, the right hand side will also vanish when $n_i = n_i^{\text{eq}}$ which is the other extreme limit of thermal equilibrium.

Having taken care of the general formalism, let us now apply it to the process of nucleosynthesis which requires protons and neutrons combining together to form bound nuclei of heavier elements like deuterium, helium *etc.*. The abundance of these elements are going to be determined by the relative abundance of neutrons and protons in the universe. Therefore, we need to first worry about the maintenance of thermal equilibrium between protons and the neutrons in the early universe. As long as the inter-conversion between n and p through the weak interaction processes ($\nu + n \leftrightarrow p + e$), ($\bar{\nu} + n \leftrightarrow p + \bar{\nu}$) and the ‘decay’ ($n \leftrightarrow p + e + \bar{\nu}$), is rapid (compared to the expansion rate of the universe), thermal equilibrium will be maintained. Then the equilibrium (n/p) ratio will be

$$\left(\frac{n_n}{n_p} \right) = \frac{X_n}{X_p} = \exp(-Q/T), \quad (24)$$

where $Q = m_n - m_p = 1.293$ MeV. At high ($T \gg Q$) temperatures, there will be equal number of neutrons and protons but as the temperature drops below about 1.3 MeV, the neutron fraction will start dropping exponentially provided thermal equilibrium is still maintained. To check whether thermal equilibrium is indeed maintained, we need to compare the expansion rate with the reaction rate. The expansion rate is given by $H = (8\pi G\rho/3)^{1/2}$ where $\rho = (\pi^2/30)gT^4$ with $g \approx 10.75$ representing the effective relativistic degrees of freedom present at these temperatures. At $T = Q$, this gives $H \approx 1.1 \text{ s}^{-1}$. The reaction rate needs to be computed from weak interaction theory. The neutron to proton conversion rate, for example, is well approximated by

$$\lambda_{np} \approx 0.29 \text{ s}^{-1} \left(\frac{T}{Q} \right)^5 \left[\left(\frac{Q}{T} \right)^2 + 6 \left(\frac{Q}{T} \right) + 12 \right]. \quad (25)$$

At $T = Q$, this gives $\lambda \approx 5 \text{ s}^{-1}$, slightly more rapid than the expansion rate. But as T drops below Q , this decreases rapidly and the reaction ceases to be fast enough to maintain thermal equilibrium. Hence we need to work out the neutron abundance by using Eq. (23).

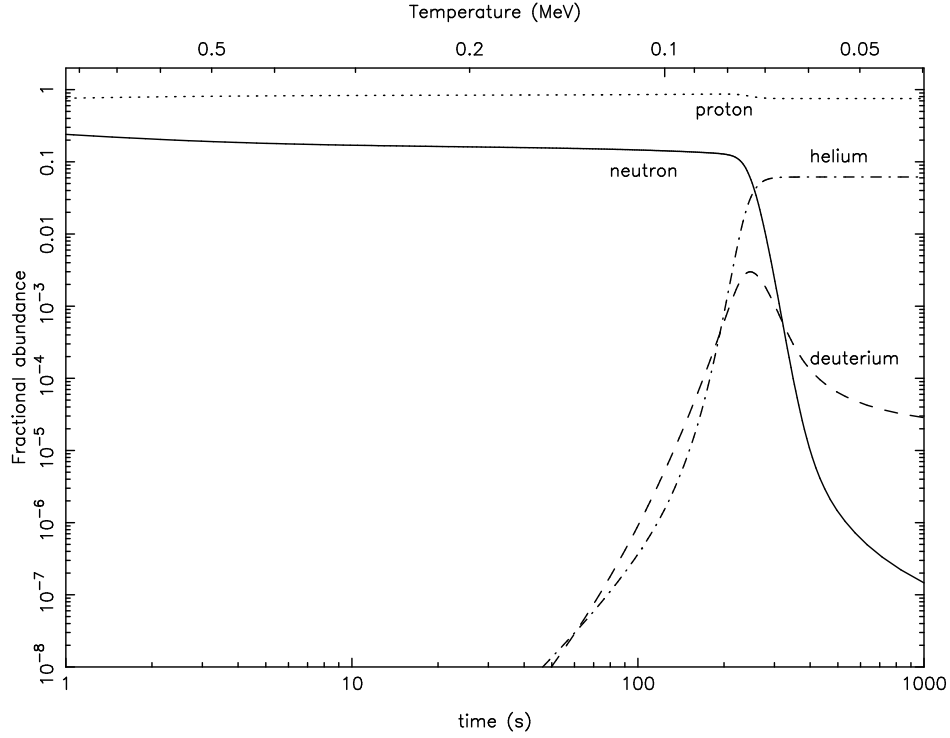


FIG. 1: The evolution of mass fraction of different species during nucleosynthesis

Using $n_1 = n_n, n_3 = n_p$ and $n_2, n_4 = n_l$ where the subscript l stands for the leptons, Eq. (23) becomes

$$\frac{1}{a^3} \frac{d(n_n a^3)}{dt} = \mu n_l^{\text{eq}} \left(\frac{n_p n_n^{\text{eq}}}{n_p^{\text{eq}}} - n_n \right). \quad (26)$$

We now use Eq. (24), write $(n_l^{\text{eq}} \mu) = \lambda_{np}$ which is the rate for neutron to proton conversion and introduce the fractional abundance $X_n = n_n / (n_n + n_p)$. Simple manipulation then leads to the equation

$$\frac{dX_n}{dt} = \lambda_{np} \left((1 - X_n) e^{-Q/T} - X_n \right). \quad (27)$$

Converting from the variable t to the variable $s = (Q/T)$ and using $(d/dt) = -HT(d/dT)$, the equations we need to solve reduce to

$$-Hs \frac{dX_n}{ds} = \lambda_{np} \left((1 - X_n) e^{-s} - X_n \right); \quad H = (1.1 \text{ sec}^{-1}) s^{-4}; \quad \lambda_{np} = \frac{0.29 \text{ s}^{-1}}{s^5} [s^2 + 6s + 12]. \quad (28)$$

It is now straightforward to integrate these equations numerically and determine how the neutron abundance changes with time. The neutron fraction falls out of equilibrium when temperatures drop below 1 MeV and it freezes to about 0.15 at temperatures below 0.5 MeV.

As the temperature decreases further, the neutron decay with a half life of $\tau_n \approx 886.7 \text{ sec}$ (which is *not* included in the above analysis) becomes important and starts depleting the neutron number density. The only way neutrons can survive is through the synthesis of light elements. As the temperature falls further to $T = T_{\text{He}} \approx 0.28 \text{ MeV}$, significant amount of He could have been produced if the nuclear reaction rates were high enough. The possible reactions which produces ${}^4\text{He}$ are $[D(D, n) {}^3\text{He}(D, p) {}^4\text{He}, D(D, p) {}^3\text{H}(D, n) {}^4\text{He}, D(D, \gamma) {}^4\text{He}]$. These are all based on D , ${}^3\text{He}$ and ${}^3\text{H}$ and do not occur rapidly enough because the mass fraction of D , ${}^3\text{He}$ and ${}^3\text{H}$ are still quite small [10^{-12} , 10^{-19} and 5×10^{-19} respectively] at $T \approx 0.3 \text{ MeV}$. The reactions $n + p \rightleftharpoons d + \gamma$ will lead to an equilibrium abundance ratio of deuterium given by

$$\frac{n_p n_n}{n_d n} = \frac{4}{3} \left(\frac{m_p m_n}{m_d} \right)^{3/2} \frac{(2\pi k_B T)^{3/2}}{(2\pi\hbar)^3 n} e^{-B/k_B T} = \exp \left[25.82 - \ln \Omega_B h^2 T_{10}^{3/2} - \left(\frac{2.58}{T_{10}} \right) \right]. \quad (29)$$

The equilibrium deuterium abundance passes through unity (for $\Omega_B h^2 = 0.02$) at the temperature of about 0.07 MeV which is when the nucleosynthesis can really begin.

So we need to determine the neutron fraction at $T = 0.07$ MeV given that it was about 0.15 at 0.5 MeV. During this epoch, the time-temperature relationship is given by $t = 130 \text{ sec } (T/0.1 \text{ MeV})^{-2}$. The neutron decay factor is $\exp(-t/\tau_n) \approx 0.74$ for $T = 0.07$ MeV. This decreases the neutron fraction to $0.15 \times 0.74 = 0.11$ at the time of nucleosynthesis. When the temperature becomes $T \lesssim 0.07$ MeV, the abundance of D and ${}^3\text{H}$ builds up and these elements further react to form ${}^4\text{He}$. A good fraction of D and ${}^3\text{H}$ is converted to ${}^4\text{He}$ (See Fig.1 which shows the growth of deuterium and its subsequent fall when helium is built up). The resultant abundance of ${}^4\text{He}$ can be easily calculated by assuming that almost all neutrons end up in ${}^4\text{He}$. Since each ${}^4\text{He}$ nucleus has two neutrons, $(n_n/2)$ helium nuclei can be formed (per unit volume) if the number density of neutrons is n_n . Thus the mass fraction of ${}^4\text{He}$ will be

$$Y = \frac{4(n_n/2)}{n_n + n_p} = \frac{2(n/p)}{1 + (n/p)} = 2x_c \quad (30)$$

where $x_c = n/(n+p)$ is the neutron abundance at the time of production of deuterium. For $\Omega_B h^2 = 0.02$, $x_c \approx 0.11$ giving $Y \approx 0.22$. Increasing baryon density to $\Omega_B h^2 = 1$ will make $Y \approx 0.25$. An accurate fitting formula for the dependence of helium abundance on various parameters is given by

$$Y = 0.226 + 0.025 \log \eta_{10} + 0.0075(g_* - 10.75) + 0.014(\tau_{1/2}(n) - 10.3 \text{ min}) \quad (31)$$

where η_{10} measures the baryon-photon ratio today via Eq. (19) and g_* is the effective number of relativistic degrees of freedom contributing to the energy density and $\tau_{1/2}(n)$ is the neutron half life. The results (of a more exact treatment) are shown in Fig. 1.

As the reactions converting D and ${}^3\text{H}$ to ${}^4\text{He}$ proceed, the number density of D and ${}^3\text{H}$ is depleted and the reaction rates - which are proportional to $\Gamma \propto X_A(\eta n_\gamma) < \sigma v >$ - become small. These reactions soon freeze-out leaving a residual fraction of D and ${}^3\text{H}$ (a fraction of about 10^{-5} to 10^{-4}). Since $\Gamma \propto \eta$ it is clear that the fraction of $(D, {}^3\text{H})$ left unreacted will decrease with η . In contrast, the ${}^4\text{He}$ synthesis - which is not limited by any reaction rate - is fairly independent of η and depends only on the (n/p) ratio at $T \simeq 0.1$ MeV. The best fits, with typical errors, to deuterium abundance calculated from the theory, for the range $\eta = (10^{-10} - 10^{-9})$ is given by

$$Y_2 \equiv \left(\frac{D}{H} \right)_p = 3.6 \times 10^{-5 \pm 0.06} \left(\frac{\eta}{5 \times 10^{-10}} \right)^{-1.6} \quad (32)$$

The production of still heavier elements - even those like ${}^{16}\text{C}$, ${}^{16}\text{O}$ which have *higher* binding energies than ${}^4\text{He}$ - is suppressed in the early universe. Two factors are responsible for this suppression: (1) For nuclear reactions to proceed, the participating nuclei must overcome their Coulomb repulsion. The probability to tunnel through the Coulomb barrier is governed by the factor $F = \exp[-2A^{1/3} (Z_1 Z_2)^{2/3} (T/1\text{MeV})^{-1/3}]$ where $A^{-1} = A_1^{-1} + A_2^{-1}$. For heavier nuclei (with larger Z), this factor suppresses the reaction rate. (2) Reaction between helium and proton would have led to an element with atomic mass 5 while the reaction of two helium nuclei would have led to an element with atomic mass 8. However, there are no stable elements in the periodic table with the atomic mass of 5 or 8! The ${}^8\text{Be}$, for example, has a half life of only 10^{-16} seconds. One can combine ${}^4\text{He}$ with ${}^8\text{Be}$ to produce ${}^{12}\text{C}$ but this can occur at significant rate only if it is a resonance reaction. That is, there should exist an excited state ${}^{12}\text{C}$ nuclei which has an energy close to the interaction energy of ${}^4\text{He} + {}^8\text{Be}$. Stars, incidentally, use this route to synthesize heavier elements. It is this triple-alpha reaction which allows the synthesis of heavier elements in stars but it is not fast enough in the early universe. (You must thank your stars that there is *no* such resonance in ${}^{16}\text{O}$ or in ${}^{20}\text{Ne}$ — which is equally important for the survival of carbon and oxygen.)

The current observations indicate, with reasonable certainty that: (i) $(D/H) \gtrsim 1 \times 10^{-5}$. (ii) $[(D + {}^3\text{He})/H] \simeq (1 - 8) \times 10^{-5}$ and (iii) $0.236 < ({}^4\text{He}/H) < 0.254$. These observations are consistent with the predictions if $10.3 \text{ min} \lesssim \tau \lesssim 10.7 \text{ min}$, and $\eta = (3 - 10) \times 10^{-10}$. Using $\eta = 2.68 \times 10^{-8} \Omega_B h^2$, this leads to the important conclusion: $0.011 \leq \Omega_B h^2 \leq 0.037$. When combined with the broad bounds on h , $0.6 \lesssim h \lesssim 0.8$, say, we can constrain the baryonic density of the universe to be: $0.01 \lesssim \Omega_B \lesssim 0.06$. These are the typical bounds on Ω_B available today. It shows that, if $\Omega_{\text{total}} \simeq 1$ then most of the matter in the universe must be non baryonic.

Since the ${}^4\text{He}$ production depends on g , the observed value of ${}^4\text{He}$ restricts the total energy density present at the time of nucleosynthesis. In particular, it constrains the number (N_ν) of light neutrinos (that is, neutrinos with $m_\nu \lesssim 1\text{MeV}$ which would have been relativistic at $T \simeq 1\text{MeV}$). The observed abundance is best explained by $N_\nu = 3$, is barely consistent with $N_\nu = 4$ and rules out $N_\nu > 4$. The laboratory bound on the total number of particles including neutrinos, which couples to the Z^0 boson is determined by measuring the decay width of the particle Z^0 ; each particle with mass less than $(m_z/2) \simeq 46 \text{ GeV}$ contributes about 180 MeV to this decay width. This bound is $N_\nu = 2.79 \pm 0.63$ which is consistent with the cosmological observations.

C. Decoupling of matter and radiation

In the early hot phase, the radiation will be in thermal equilibrium with matter; as the universe cools below $k_B T \simeq (\epsilon_a/10)$ where ϵ_a is the binding energy of atoms, the electrons and ions will combine to form neutral atoms and radiation will decouple from matter. This occurs at $T_{\text{dec}} \simeq 3 \times 10^3$ K. As the universe expands further, these photons will continue to exist without any further interaction. It will retain thermal spectrum since the redshift of the frequency $\nu \propto a^{-1}$ is equivalent to changing the temperature in the spectrum by the scaling $T \propto (1/a)$. It turns out that the major component of the extra-galactic background light (EBL) which exists today is in the microwave band and can be fitted very accurately by a thermal spectrum at a temperature of about 2.73 K. It seems reasonable to interpret this radiation as a relic arising from the early, hot, phase of the evolving universe. This relic radiation, called *cosmic microwave background radiation*, turns out to be a gold mine of cosmological information and is extensively investigated in recent times. We shall now discuss some details related to the formation of neutral atoms and the decoupling of photons.

The relevant reaction is, of course, $e + p \rightleftharpoons H + \gamma$ and if the rate of this reaction is faster than the expansion rate, then one can calculate the neutral fraction using Saha's equation. Introducing the fractional ionisation, X_i , for each of the particle species and using the facts $n_p = n_e$ and $n_p + n_H = n_B$, it follows that $X_p = X_e$ and $X_H = (n_H/n_B) = 1 - X_e$. Saha's equation now gives

$$\frac{1 - X_e}{X_e^2} \cong 3.84\eta(T/m_e)^{3/2} \exp(B/T) \quad (33)$$

where $\eta = 2.68 \times 10^{-8}(\Omega_B h^2)$ is the baryon-to-photon ratio. We may define T_{atom} as the temperature at which 90 percent of the electrons, say, have combined with protons: i.e. when $X_e = 0.1$. This leads to the condition:

$$(\Omega_B h^2)^{-1} \tau^{-\frac{3}{2}} \exp[-13.6\tau^{-1}] = 3.13 \times 10^{-18} \quad (34)$$

where $\tau = (T/1\text{eV})$. For a given value of $(\Omega_B h^2)$, this equation can be easily solved by iteration. Taking logarithms and iterating once we find $\tau^{-1} \cong 3.084 - 0.0735 \ln(\Omega_B h^2)$ with the corresponding redshift $(1+z) = (T/T_0)$ given by

$$(1+z) = 1367[1 - 0.024 \ln(\Omega_B h^2)]^{-1}. \quad (35)$$

For $\Omega_B h^2 = 1, 0.1, 0.01$ we get $T_{\text{atom}} \cong 0.324\text{eV}, 0.307\text{eV}, 0.292\text{eV}$ respectively. These values correspond to the redshifts of 1367, 1296 and 1232.

Because the preceding analysis was based on equilibrium densities, it is important to check that the rate of the reactions $p + e \leftrightarrow H + \gamma$ is fast enough to maintain equilibrium. For $\Omega_B h^2 \approx 0.02$, the equilibrium condition is only marginally satisfied, making this analysis suspect. More importantly, the direct recombination to the ground state of the hydrogen atom — which was used in deriving the Saha's equation — is not very effective in producing neutral hydrogen in the early universe. The problem is that each such recombination releases a photon of energy 13.6 eV which will end up ionizing another neutral hydrogen atom which has been formed earlier. As a result, the direct recombination to the ground state does not change the neutral hydrogen fraction at the lowest order. Recombination through the excited states of hydrogen is more effective since such a recombination ends up emitting more than one photon each of which has an energy less than 13.6 eV. Given these facts, it is necessary to once again use the rate equation developed in the previous section to track the evolution of ionisation fraction.

A simple procedure for doing this, which captures the essential physics, is as follows: We again begin with Eq. (23) and repeating the analysis done in the last section, now with $n_1 = n_e, n_2 = n_p, n_3 = n_H$ and $n_4 = n_\gamma$, and defining $X_e = n_e/(n_e + n_H) = n_p/n_H$ one can easily derive the rate equation for this case:

$$\frac{dX_e}{dt} = [\beta(1 - X_e) - \alpha n_b X_e^2] = \alpha \left(\frac{\beta}{\alpha} (1 - X_e) - n_b X_e^2 \right). \quad (36)$$

This equation is analogous to Eq. (27); the first term gives the photoionisation rate which produces the free electrons and the second term is the recombination rate which converts free electrons into hydrogen atom and we have used the fact $n_e = n_b X_e$ etc.. Since we know that direct recombination to the ground state is not effective, the recombination rate α is the rate for capture of electron by a proton forming an excited state of hydrogen. To a good approximation, this rate is given by

$$\alpha = 9.78 r_0^2 c \left(\frac{B}{T} \right)^{1/2} \ln \left(\frac{B}{T} \right) \quad (37)$$

where $r_0 = e^2/m_e c^2$ is the classical electron radius. To integrate Eq. (36) we also need to know β/α . This is easy because in thermal equilibrium the right hand side of Eq. (36) should vanish and Saha's equation tells us the value of X_e in thermal equilibrium. On using Eq. (33), this gives

$$\frac{\beta}{\alpha} = \left(\frac{m_e T}{2\pi}\right)^{3/2} \exp[-(B/T)]. \quad (38)$$

We can now integrate Eq. (36) using the variable B/T just as we used the variable Q/T in solving Eq. (27). The result shows that the actual recombination proceeds more slowly compared to that predicted by the Saha's equation. The actual fractional ionisation is higher than the value predicted by Saha's equation at temperatures below about 1300. For example, at $z = 1300$, these values differ by a factor 3; at $z \simeq 900$, they differ by a factor of 200. The value of T_{atom} , however, does not change significantly. A more rigorous analysis shows that, in the redshift range of $800 < z < 1200$, the fractional ionisation varies rapidly and is given (approximately) by the formula,

$$X_e = 2.4 \times 10^{-3} \frac{(\Omega_{\text{NR}} h^2)^{1/2}}{(\Omega_B h^2)} \left(\frac{z}{1000}\right)^{12.75}. \quad (39)$$

This is obtained by fitting a curve to the numerical solution.

The formation of neutral atoms makes the photons decouple from the matter. The redshift for decoupling can be determined as the epoch at which the optical depth for photons is unity. Using Eq. (39), we can compute the optical depth for photons to be

$$\tau = \int_0^t n(t) X_e(t) \sigma_T dt = \int_0^z n(z) X_e(z) \sigma_T \left(\frac{dt}{dz}\right) dz \simeq 0.37 \left(\frac{z}{1000}\right)^{14.25} \quad (40)$$

where we have used the relation $H_0 dt \simeq -\Omega_{\text{NR}}^{-1/2} z^{-5/2} dz$ which is valid for $z \gg 1$. This optical depth is unity at $z_{\text{dec}} = 1072$. From the optical depth, we can also compute the probability that the photon was last scattered in the interval $(z, z + dz)$. This is given by $(\exp -\tau) (d\tau/dz)$ which can be expressed as

$$P(z) = e^{-\tau} \frac{d\tau}{dz} = 5.26 \times 10^{-3} \left(\frac{z}{1000}\right)^{13.25} \exp \left[-0.37 \left(\frac{z}{1000}\right)^{14.25} \right]. \quad (41)$$

This $P(z)$ has a sharp maximum at $z \simeq 1067$ and a width of about $\Delta z \simeq 80$. It is therefore reasonable to assume that decoupling occurred at $z \simeq 1070$ in an interval of about $\Delta z \simeq 80$. We shall see later that the finite thickness of the surface of last scattering has important observational consequences.

III. STRUCTURE FORMATION AND LINEAR PERTURBATION THEORY

Having discussed the evolution of the background universe, we now turn to the study of structure formation. Before discussing the details, let us briefly summarise the broad picture and give references to some of the topics that we will *not* discuss. The key idea is that if there existed small fluctuations in the energy density in the early universe, then gravitational instability can amplify them in a well-understood manner leading to structures like galaxies etc. today. The most popular model for generating these fluctuations is based on the idea that if the very early universe went through an inflationary phase [9], then the quantum fluctuations of the field driving the inflation can lead to energy density fluctuations [10, 11]. It is possible to construct models of inflation such that these fluctuations are described by a Gaussian random field and are characterized by a power spectrum of the form $P(k) = Ak^n$ with $n \simeq 1$ (see Sec. VII). The models cannot predict the value of the amplitude A in an unambiguous manner but it can be determined from CMBR observations. The CMBR observations are consistent with the inflationary model for the generation of perturbations and gives $A \simeq (28.3h^{-1} Mpc)^4$ and $n = 0.97 \pm 0.023$ (The first results were from COBE [12] and WMAP has reconfirmed them with far greater accuracy). When the perturbation is small, one can use well defined linear perturbation theory to study its growth. But when $\delta \approx (\delta\rho/\rho)$ is comparable to unity the perturbation theory breaks down. Since there is more power at small scales, smaller scales go non-linear first and structure forms hierarchically. The non linear evolution of the *dark matter halos* (which is an example of statistical mechanics of self gravitating systems; see e.g. [13]) can be understood by simulations as well as theoretical models based on approximate ansatz [14] and nonlinear scaling relations [15]. The baryons in the halo will cool and undergo collapse in a fairly complex manner because of gas dynamical processes. It seems unlikely that the baryonic collapse and galaxy formation can be understood by analytic approximations; one needs to do high resolution computer simulations to make any progress [16]. All these results are broadly consistent with observations.

As long as these fluctuations are small, one can study their evolution by linear perturbation theory, which is what we will start with [17]. The basic idea of linear perturbation theory is well defined and simple. We perturb the background FRW metric by $g_{ik}^{FRW} \rightarrow g_{ik}^{FRW} + h_{ik}$ and also perturb the source energy momentum tensor by $T_{ik}^{FRW} \rightarrow T_{ik}^{FRW} + \delta T_{ik}$. Linearising the Einstein's equations, one can relate the perturbed quantities by a relation of the form $\mathcal{L}(g_{ik}^{FRW})h_{ik} = \delta T_{ik}$ where \mathcal{L} is second order linear differential operator depending on the back ground metric g_{ik}^{FRW} . Since the background is maximally symmetric, one can separate out time and space; for e.g, if $k = 0$, simple Fourier modes can be used for this purpose and we can write down the equation for any given mode, labelled by a wave vector \mathbf{k} as:

$$\mathcal{L}(a(t), \mathbf{k})h_{ab}(t, \mathbf{k}) = \delta T_{ab}(t, \mathbf{k}) \quad (42)$$

To every mode we can associate a wavelength normalized to today's value: $\lambda(t) = (2\pi/k)(1+z)^{-1}$ and a corresponding mass scale which is invariant under expansion:

$$M = \frac{4\pi\rho(t)}{3} \left[\frac{\lambda(t)}{2} \right]^3 = \frac{4\pi\rho_0}{3} \left(\frac{\lambda_0}{2} \right)^3 = 1.5 \times 10^{11} M_{\odot} (\Omega_m h^2) \left(\frac{\lambda_0}{1 \text{ Mpc}} \right)^3. \quad (43)$$

The behaviour of the mode depends on the relative value of $\lambda(t)$ as compared to the Hubble radius $d_H(t) \equiv (\dot{a}/a)^{-1}$. Since the Hubble radius: $d_H(t) \propto t$ while the wavelength of the mode: $\lambda(t) \propto a(t) \propto (t^{1/2}, t^{2/3})$ in the radiation dominated and matter dominated phases it follows that $\lambda(t) > d_H(t)$ at sufficiently early times. When $\lambda(t) = d_H(t)$, we say that the mode is entering the Hubble radius. Since the Hubble radius at $z = z_{eq}$ is

$$\lambda_{eq} \cong \left(\frac{H_0^{-1}}{\sqrt{2}} \right) \left(\frac{\Omega_R^{1/2}}{\Omega_{NR}} \right) \cong 14 \text{ Mpc} (\Omega_{NR} h^2)^{-1} \quad (44)$$

it follows that modes with $\lambda_0 > \lambda_{eq}$ enter Hubble radius in MD phase while the more relevant modes with $\lambda < \lambda_{eq}$ enter in the RD phase. Thus, for a given mode we can identify three distinct phases: First, very early on, when $\lambda > d_H, z > z_{eq}$ the dynamics is described by general relativity. In this stage, the universe is radiation dominated, gravity is the only relevant force and the perturbations are linear. Next, when $\lambda < d_H$ and $z > z_{eq}$ one can describe the dynamics by Newtonian considerations. The perturbations are still linear and the universe is radiation dominated. Finally, when $\lambda < d_H, z < z_{eq}$ we have a matter dominated universe in which we can use the Newtonian formalism; but at this stage — when most astrophysical structures form — we need to grapple with nonlinear astrophysical processes.

Let us now consider the metric perturbation in greater detail. When the metric is perturbed to the form: $g_{ab} \rightarrow g_{ab} + h_{ab}$ the perturbation can be split as $h_{ab} = (h_{00}, h_{0\alpha} \equiv w_{\alpha}, h_{\alpha\beta})$. We also know that any 3-vector $\mathbf{w}(\mathbf{x})$ can be split as $\mathbf{w} = \mathbf{w}^{\perp} + \mathbf{w}^{\parallel}$ in which $\mathbf{w}^{\parallel} = \nabla\Phi^{\parallel}$ is curl-free (and carries one degree of freedom) while \mathbf{w}^{\perp} is divergence-free (and has 2 degrees of freedom). This result is obvious in \mathbf{k} -space since we can write any vector $\mathbf{w}(\mathbf{k})$ as a sum of two terms, one along \mathbf{k} and one transverse to \mathbf{k} :

$$\mathbf{w}(\mathbf{k}) = \mathbf{w}^{\parallel}(\mathbf{k}) + \mathbf{w}^{\perp}(\mathbf{k}) = \mathbf{k} \underbrace{\left(\frac{\mathbf{w}(\mathbf{k}) \cdot \mathbf{k}}{k^2} \right)}_{\text{along } \mathbf{k}} + \underbrace{\left[\mathbf{w}(\mathbf{k}) - \mathbf{k} \left(\frac{\mathbf{k} \cdot \mathbf{w}(\mathbf{k})}{k^2} \right) \right]}_{\text{transverse to } \mathbf{k}}; \quad \mathbf{k} \times \mathbf{w}^{\parallel} = 0; \quad \mathbf{k} \cdot \mathbf{w}^{\perp} = 0 \quad (45)$$

Fourier transforming back, we can split \mathbf{w} into a curl-free and divergence-free parts. Similar decomposition works for $h_{\alpha\beta}$ by essentially repeating the above analysis on each index. We can write:

$$h_{\alpha\beta} = \underbrace{\psi \delta_{\alpha\beta}}_{\text{trace}} + \underbrace{(\nabla_{\alpha} u_{\beta}^{\perp} + \nabla_{\beta} u_{\alpha}^{\perp})}_{\text{traceless from vector}} + \underbrace{\left(\nabla_{\alpha} \nabla_{\beta} - \frac{1}{3} \delta_{\alpha\beta} \nabla^2 \right) \Phi_1 + h_{\alpha\beta}^{\perp\perp}}_{\text{traceless from scalar}} \Rightarrow 1 + 2 + 1 + 2 = 6 \quad (46)$$

The u_{α}^{\perp} is divergence free and $h_{\alpha\beta}^{\perp\perp}$ is traceless and divergence free. Thus the most general perturbation h_{ab} (ten degrees of freedom) can be built out of

$$h_{ab} = (h_{00}, h_{0\alpha} \equiv w_{\alpha}, h_{\alpha\beta}) = [h_{00}, (\Phi^{\parallel}, \mathbf{w}^{\perp}), (\psi, \Phi_1, u_{\alpha}^{\perp}, h_{\alpha\beta}^{\perp\perp})] \Rightarrow [1, (1, 2), (1, 1, 2, 2)] \quad (47)$$

We now use the freedom available in the choice of four coordinate transformations to set four conditions: $\Phi^{\parallel} = \Phi_1 = 0$ and $u_{\alpha}^{\perp} = 0$ thereby leaving six degrees of freedom in $(h_{00} \equiv 2\Phi, \psi, \mathbf{w}^{\perp}, h_{\alpha\beta}^{\perp\perp})$ as nonzero. Then the perturbed line element takes the form:

$$ds^2 = a^2(\eta) \{ [1 + 2\Phi(\mathbf{x}, \eta)] d\eta^2 - 2w_{\alpha}^{\perp}(\mathbf{x}, \eta) d\eta dx^{\alpha} - \{ (1 - 2\psi(\mathbf{x}, \eta)) \delta_{\alpha\beta} + 2h_{\alpha\beta}^{\perp\perp}(\mathbf{x}, \eta) \} dx^{\alpha} dx^{\beta} \} \quad (48)$$

To make further simplification we need to use two facts from Einstein's equations. It turns out that the Einstein's equations for \mathbf{w}^\perp and $h_{\alpha\beta}^\perp$ decouple from those for (Φ, ψ) . Further, in the absence of anisotropic stress, one of the equations give $\psi = \Phi$. If we use these two facts, we can simplify the structure of perturbed metric drastically. As far as the growth of matter perturbations are concerned, we can ignore w_α^\perp and $h_{\alpha\beta}^\perp$ and work with a simple metric:

$$ds^2 = a^2(\eta)[(1 + 2\Phi)d\eta^2 - (1 - 2\Phi)\delta_{\alpha\beta}dx^\alpha dx^\beta] \quad (49)$$

with just one perturbed scalar degree of freedom in Φ . This is what we will study.

Having decided on the gauge, let us consider the evolution equations for the perturbations. While one can directly work with the Einstein's equations, it turns out to be convenient to use the equations of motion for matter variables, since we are eventually interested in the matter perturbations. In what follows, we will use the over-dot to denote $(d/d\eta)$ so that the standard Hubble parameter is $H = (1/a)(da/dt) = \dot{a}/a^2$. With this notation, the continuity equation becomes:

$$\dot{\rho} + 3(aH - \dot{\Phi})(\rho + p) = -\nabla_\alpha[(\rho + p)v^\alpha] \quad (50)$$

Since the momentum flux in the relativistic case is $(\rho + p)v^\alpha$, all the terms in the above equation are intuitively obvious, except probably the $\dot{\Phi}$ term. To see the physical origin of this term, note that the perturbation in Eq. (49) changes the factor in front of the spatial metric from a^2 to $a^2(1 - 2\Phi)$ so that $\ln a \rightarrow \ln a - \Phi$; hence the effective Hubble parameter from (\dot{a}/a) to $(\dot{a}/a) - \dot{\Phi}$ which explains the extra $\dot{\Phi}$ term. This is, of course, the exact equation for matter variables in the perturbed metric given by Eq. (49); but we only need terms which are of linear order. Writing the curl-free velocity part as $v^\alpha = \nabla^\alpha v$, the *linearised* equations, for dark matter (with $p = 0$) and radiation (with $p = (1/3)\rho$) perturbations are given by:

$$\dot{\delta}_m = \frac{d}{d\eta} \left(\frac{\delta n_m}{n_m} \right) = \nabla^2 v_m + 3\dot{\Phi}; \quad \frac{3}{4}\dot{\delta}_R = \frac{d}{d\eta} \left(\frac{\delta n_R}{n_R} \right) = \nabla^2 v_R + 3\dot{\Phi} \quad (51)$$

where n_m and n_R are the number densities of dark matter particles and radiation. The same equations in Fourier space [using the same symbols for, say, $\delta(t, \mathbf{x})$ or $\delta(t, \mathbf{k})$] are simpler to handle:

$$\dot{\delta}_m = \frac{d}{d\eta} \left(\frac{\delta n_m}{n_m} \right) = -k^2 v_m + 3\dot{\Phi}; \quad \frac{3}{4}\dot{\delta}_R = \frac{d}{d\eta} \left(\frac{\delta n_R}{n_R} \right) = -k^2 v_R + 3\dot{\Phi} \quad (52)$$

Note that these equations imply

$$\frac{d}{d\eta} \left[\frac{\delta n_R}{n_R} - \frac{\delta n_m}{n_m} \right] = \frac{d}{d\eta} \left[\delta \ln \left(\frac{n_R}{n_m} \right) \right] = \frac{d}{d\eta} \left[\delta \left(\ln \left(\frac{s}{n_m} \right) \right) \right] = -k^2(v_R - v_m) \quad (53)$$

For long wavelength perturbations (in the limit of $k \rightarrow 0$), this will lead to the conservation of perturbation $\delta(s/n_m)$ in the entropy per particle.

Let us next consider the Euler equation which has the general form:

$$\partial_\eta[(\rho + p)v^\alpha] = -(\rho + p)\nabla^\alpha\Phi - \nabla^\alpha p - 4aH(\rho + p)v^\alpha \quad (54)$$

Once again each of the terms is simple to interpret. The $(\rho + p)$ arises because the pressure also contributes to inertia in a relativistic theory and the factor 4 in the last term on the right hand side arises because the term $v^\alpha \partial_\eta(\rho + p)$ on the left hand side needs to be compensated. Taking the linearised limit of this equation, for dark matter and radiation, we get:

$$\dot{v}_m = \Phi - aHv_m; \quad \dot{v}_R = \Phi + \frac{1}{4}\delta_R \quad (55)$$

Thus we now have four equations in Eqs. (52), (55) for the five variables $(\delta_m, \delta_R, v_m, v_R, \Phi)$. All we need to do is to pick one more from Einstein's equations to complete the set. The Einstein's equations for our perturbed metric are:

$${}^0_0 \text{ component} : k^2\Phi + 3\frac{\dot{a}}{a} \left(\dot{\Phi} + \frac{\dot{a}}{a}\Phi \right) = -4\pi Ga^2 \sum_A \rho_A \delta_A = -4\pi Ga^2 \rho_{bg} \delta_{total} \quad (56)$$

$${}^0_\alpha \text{ component} : \dot{\Phi} + \frac{\dot{a}}{a}\Phi = -4\pi Ga^2 \sum_A (\rho + p)_A v_A; \quad \mathbf{v} = \nabla v \quad (57)$$

$${}^\alpha_\alpha \text{ component} : 3\frac{\dot{a}}{a}\dot{\Phi} + 2\frac{\ddot{a}}{a}\Phi - \frac{\dot{a}^2}{a^2}\Phi + \ddot{\Phi} = 4\pi Ga^2 \delta p \quad (58)$$

where A denotes different components like dark matter, radiation etc. Using Eq. (57) in Eq. (56) we can get a modified Poisson equation which is purely algebraic:

$$-k^2\Phi = 4\pi G a^2 \sum_A \left(\rho_A \delta_A - 3 \left(\frac{\dot{a}}{a} \right) (\rho_A + p_A) v_A \right) \quad (59)$$

which once again emphasizes the fact that in the relativistic theory, both pressure and density act as source of gravity.

To get a feel for the solutions let us consider a flat universe dominated by a single component of matter with the equation of state $p = w\rho$. (A purely radiation dominated universe, for example, will have $w = 1/3$.) In this case the Friedmann background equation gives $\rho \propto a^{-3(1+w)}$ and

$$\frac{\dot{a}}{a} = \frac{2}{(1+3w)\eta}; \quad \frac{\ddot{a}}{a} = \frac{2(1-3w)}{(1+3w)^2\eta^2} \quad (60)$$

The equation for the potential Φ can be reduced to the form:

$$\ddot{\Phi} + \frac{6(1+w)}{1+3w} \frac{\dot{\Phi}}{\eta} + k^2 w \Phi = 0 \quad (61)$$

The second term is the damping due to the expansion while last term is the pressure support that will lead to oscillations. Clearly, the factor $k\eta$ determines which of these two terms dominates. When the pressure term dominates ($k\eta \gg 1$), we expect oscillatory behaviour while when the background expansion dominates ($k\eta \ll 1$), we expect the growth to be suppressed. This is precisely what happens. The exact solution is given in terms of the Bessel functions

$$\Phi(\eta) = \frac{C_1(\mathbf{k})J_{\nu/2}(\sqrt{w}k\eta) + C_2(\mathbf{k})Y_{\nu/2}(\sqrt{w}k\eta)}{\eta^{\nu/2}}; \quad \nu = \frac{5+3w}{1+3w} \quad (62)$$

From the theory of Bessel functions, we know that:

$$\lim_{x \rightarrow 0} J_{\nu/2}(x) \simeq \frac{x^{\nu/2}}{2^{\nu/2}\Gamma(\nu/2+1)}; \quad \lim_{x \rightarrow 0} Y_{\nu/2}(x) \propto -\frac{1}{x^{\nu/2}} \quad (63)$$

This shows that if we want a finite value for Φ as $\eta \rightarrow 0$, we can set $C_2 = 0$. This gives the gravitational potential to be

$$\Phi(\eta) = \frac{C_1(\mathbf{k})J_{\nu/2}(\sqrt{w}k\eta)}{\eta^{\nu/2}}; \quad \nu = \frac{5+3w}{1+3w} \quad (64)$$

The corresponding density perturbation will be:

$$\delta = -2\Phi - \frac{(1+3w)^2 k^2 \eta^2}{6} \frac{C_1(\mathbf{k})J_{\nu/2}(\sqrt{w}k\eta)}{\eta^{\nu/2}} + (1+3w)\sqrt{w}k\eta \frac{C_1(\mathbf{k})J_{(\nu/2)+1}(\sqrt{w}k\eta)}{\eta^{\nu/2}} \quad (65)$$

To understand the nature of the solution, note that $d_H = (\dot{a}/a)^{-1} \propto \eta$ and $kd_H \simeq d_H/\lambda \propto k\eta$. So the argument of the Bessel function is just the ratio (d_H/λ). From the theory of Bessel functions, we know that for small values of the argument $J_\nu(x) \propto x^\nu$ is a power law while for large values of the argument it oscillates with a decaying amplitude:

$$\lim_{x \rightarrow \infty} J_\nu(x) \sim \frac{\cos[x - (\nu-1)\pi/4]}{\sqrt{x}}; \quad (66)$$

Hence, for modes which are still outside the Hubble radius ($k \ll \eta^{-1}$), we have a constant amplitude for the potential and density contrast:

$$\Phi \approx \Phi_i(\mathbf{k}); \quad \delta \approx -2\Phi_i(\mathbf{k}) \quad (67)$$

That is, the perturbation is frozen (except for a decaying mode) at a constant value. On the other hand, for modes which are inside the Hubble radius ($k \gg \eta^{-1}$), the perturbation is rapidly oscillatory (if $w \neq 0$). That is the pressure is effective at small scales and leads to acoustic oscillations in the medium.

A special case of the above is the flat, matter-dominated universe with $w = 0$. In this case, we need to take the $w \rightarrow 0$ limit and the general solution is indeed a constant $\Phi = \Phi_i(\mathbf{k})$ (plus a decaying mode $\Phi_{decay} \propto \eta^{-5}$ which diverges as $\eta \rightarrow 0$). The corresponding density perturbations is:

$$\delta = -(2 + \frac{k^2 \eta^2}{6}) \Phi_i(\mathbf{k}) \quad (68)$$

which shows that density perturbation is ‘‘frozen’’ at large scales but grows at small scales:

$$\delta = \begin{cases} -2\Phi_i(\mathbf{k}) = \text{constant} & (k\eta \ll 1) \\ -\frac{1}{6}k^2\eta^2\Phi_i(\mathbf{k}) \propto \eta^2 \propto a & (k\eta \gg 1) \end{cases} \quad (69)$$

We will use these results later on.

IV. PERTURBATIONS IN DARK MATTER AND RADIATION

We shall now move on to the more realistic case of a multi-component universe consisting of radiation and collisionless dark matter. (For the moment we are ignoring the baryons, which we will study in Sec. VI). It is convenient to use $y = a/a_{eq}$ as independent variable rather than the time coordinate. The background expansion of the universe described by the function $a(t)$ can be equivalently expressed (in terms of the conformal time η) as

$$y \equiv \frac{\rho_M}{\rho_R} = \frac{a}{a_{eq}} = x^2 + 2x, \quad x \equiv \left(\frac{\Omega_M}{4a_{eq}}\right)^{1/2} H_0 \eta \quad (70)$$

It is also useful to define a critical wave number k_c by:

$$k_c^2 = \frac{H_0^2 \Omega_m}{a_{eq}} = 4(\sqrt{2} - 1)^2 \eta_{eq}^{-2} = 4(\sqrt{2} - 1)^2 k_{eq}^2; \quad k_c^{-1} = 19(\Omega_m h^2)^{-1} Mpc \quad (71)$$

which essentially sets the comoving scale corresponding to matter-radiation equality. Note that $2x = k_c \eta$ and $y \approx k_c \eta$ in the radiation dominated phase while $y = (1/4)(k_c \eta)^2$ in the matter dominated phase.

We now manipulate Eqs. (52), (55), (56), (57) governing the growth of perturbations by essentially eliminating the velocity. This leads to the three equations

$$y\Phi' + \Phi + \frac{1}{3} \frac{k^2}{k_c^2} \frac{y^2}{1+y} \Phi = -\frac{1}{2} \frac{y}{1+y} \left(\delta_m + \frac{1}{y} \delta_R \right) \quad (72)$$

$$(1+y)\delta_m'' + \frac{2+3y}{2y} \delta_m' = 3(1+y)\Phi'' + \frac{3(2+3y)}{2y} \Phi' - \frac{k^2}{k_c^2} \Phi \quad (73)$$

$$(1+y)\delta_R'' + \frac{1}{2} \delta_R' + \frac{1}{3} \frac{k^2}{k_c^2} \delta_R = 4(1+y)\Phi'' + 2\Phi' - \frac{4}{3} \frac{k^2}{k_c^2} \Phi \quad (74)$$

for the three unknowns Φ, δ_m, δ_R . Given suitable initial conditions we can solve these equations to determine the growth of perturbations. The initial conditions need to be imposed very early on when the modes are much bigger than the Hubble radius which corresponds to the $y \ll 1, k \rightarrow 0$ limit. In this limit, the equations become:

$$y\Phi' + \Phi \approx -\frac{1}{2} \delta_R; \quad \delta_m'' + \frac{1}{y} \delta_m' \approx 3\Phi'' + \frac{3}{y} \Phi'; \quad \delta_R'' + \frac{1}{2} \delta_R' \approx 4\Phi'' + 2\Phi' \quad (75)$$

We will take $\Phi(y_i, k) = \Phi_i(k)$ as given value, to be determined by the processes that generate the initial perturbations. First equation in Eq. (75) shows that we can take $\delta_R = -2\Phi_i$ for $y_i \rightarrow 0$. Further Eq. (53) shows that adiabaticity is respected at these scales and we can take $\delta_m = (3/4)\delta_R = -(3/2)\Phi_i$. The exact equation Eq. (72) determines Φ' if $(\Phi, \delta_m, \delta_R)$ are given. Finally we use the last two equations to set $\delta_m' = 3\Phi', \delta_R' = 4\Phi'$. Thus we take the initial conditions at some $y = y_i \ll 1$ to be:

$$\Phi(y_i, k) = \Phi_i(k); \quad \delta_R(y_i, k) = -2\Phi_i(k); \quad \delta_m(y_i, k) = -(3/2)\Phi_i(k) \quad (76)$$

with $\delta_m'(y_i, k) = 3\Phi'(y_i, k); \delta_R'(y_i, k) = 4\Phi'(y_i, k)$.

Given these initial conditions, it is fairly easy to integrate the equations forward in time and the numerical results are shown in Figs 2, 3, 4, 5. (In the figures k_{eq} is taken to be $a_{eq} H_{eq}$.) To understand the nature of the evolution, it is, however, useful to try out a few analytic approximations to Eqs. (72) – (74) which is what we will do now.

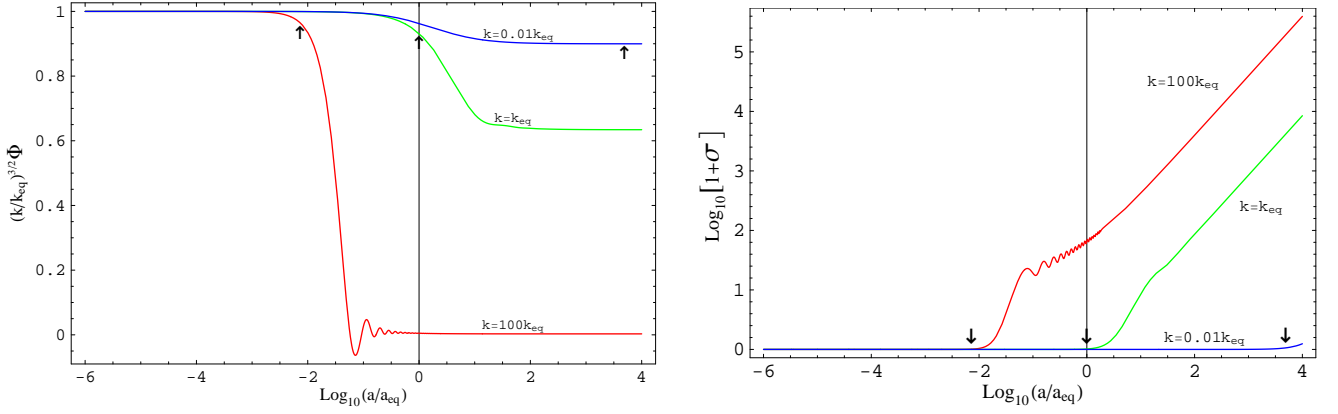


FIG. 2: Left Panel: The evolution of gravitational potential Φ for 3 different modes. The wavenumber is indicated by the label and the epoch at which the mode enters the Hubble radius is indicated by a small arrow. The top most curve is for a mode which stays outside the Hubble radius for most of its evolution and is well described by Eq. (78). The other two modes show the decay of Φ after the mode has entered the Hubble radius in the radiation dominated epoch as described by Eq. (79). Right Panel: Evolution of entropy perturbation (see Eq. (87) for the definition). The entropy perturbation is essentially zero till the mode enters Hubble radius and grows afterwards tracking the dominant energy density perturbation.

A. Evolution for $\lambda \gg d_H$

Let us begin by considering very large wavelength modes corresponding to the $k\eta \rightarrow 0$ limit. In this case adiabaticity is respected and we can set $\delta_R \approx (4/3)\delta_m$. Then Eqs. (72), (73) become

$$y\Phi' + \Phi \approx -\frac{3y+4}{8(1+y)}\delta_R; \quad \delta'_R \approx 4\Phi' \quad (77)$$

Differentiating the first equation and using the second to eliminate δ_m , we get a second order equation for Φ . Fortunately, this equation has an exact solution

$$\Phi = \Phi_i \frac{1}{10y^3} \left[16\sqrt{1+y} + 9y^3 + 2y^2 - 8y - 16 \right]; \quad \delta_R \approx 4\Phi - 6\Phi_i \quad (78)$$

[There is simple way of determining such an exact solution, which we will describe in Sec. IV D.]. The initial condition on δ_R is chosen such that it goes to $-2\Phi_i$ initially. The solution shows that, as long as the mode is bigger than the Hubble radius, the potential changes very little; it is constant initially as well as in the final matter dominated phase. At late times ($y \gg 1$) we see that $\Phi \approx (9/10)\Phi_i$ so that Φ decreases only by a factor (9/10) during the entire evolution if $k \rightarrow 0$ is a valid approximation.

B. Evolution for $\lambda \ll d_H$ in the radiation dominated phase

When the mode enters Hubble radius in the radiation dominated phase, we can no longer ignore the pressure terms. The pressure makes radiation density contrast oscillate and the gravitational potential, driven by this, also oscillates with a decay in the overall amplitude. An approximate procedure to describe this phase is to solve the coupled $\delta_R - \Phi$ system, ignoring δ_m which is sub-dominant and *then* determine δ_m using the form of Φ .

When δ_m is ignored, the problem reduces to the one solved earlier in Eqs (64), (65) with $w = 1/3$ giving $\nu = 3$. Since $J_{3/2}$ can be expressed in terms of trigonometric functions, the solution given by Eq. (64) with $\nu = 3$, simplifies to

$$\Phi = \Phi_i \frac{3}{l^3 y^3} [\sin(ly) - ly \cos(ly)]; \quad l^2 = \frac{k^2}{3k_c^2} \quad (79)$$

Note that as $y \rightarrow 0$, we have $\Phi = \Phi_i, \Phi' = 0$. This solution shows that once the mode enters the Hubble radius, the potential decays in an oscillatory manner. For $ly \gg 1$, the potential becomes $\Phi \approx -3\Phi_i (ly)^{-2} \cos(ly)$. In the same

limit, we get from Eq. (65) that

$$\delta_R \approx -\frac{2}{3}k^2\eta^2\Phi \approx -2l^2y^2\Phi \approx 6\Phi_i \cos(ly) \quad (80)$$

(This is analogous to Eq. (68) for the radiation dominated case.) This oscillation is seen clearly in Fig 3 and Fig.4 (left panel). The amplitude of oscillations is accurately captured by Eq. (80) for $k = 100k_{eq}$ mode but not for $k = k_{eq}$; this is to be expected since the mode is not entering in the radiation dominated phase.

Let us next consider matter perturbations during this phase. They grow, driven by the gravitational potential determined above. When $y \ll 1$, Eq.(73) becomes:

$$\delta_m'' + \frac{1}{y}\delta_m' = 3\Phi'' + \frac{3}{y}\Phi' - \frac{k^2}{k_c^2}\Phi \quad (81)$$

The Φ is essentially determined by radiation and satisfies Eq. (61); using this, we can rewrite Eq. (81) as

$$\frac{d}{dy}(y\delta_m') = -9(\Phi' + \frac{2}{3}l^2y\Phi) \quad (82)$$

The general solution to the homogeneous part of Eq. (82) (obtained by ignoring the right hand side) is $(c_1 + c_2 \ln y)$; hence the general solution to this equation is

$$\delta_m = (c_1 + c_2 \ln y) - 9 \int \frac{dy}{y} \int^y dy_1 [\Phi'(y_1) + \frac{2}{3}l^2y_1\Phi(y_1)] \quad (83)$$

For $y \ll 1$ the growing mode varies as $\ln y$ and dominates over the rest; hence we conclude that, matter, driven by Φ , grows logarithmically during the radiation dominated phase for modes which are inside the Hubble radius.

C. Evolution in the matter dominated phase

Finally let us consider the matter dominated phase, in which we can ignore the radiation and concentrate on Eq. (72) and Eq. (73). When $y \gg 1$ these equations become:

$$y\Phi' + \Phi \approx -\frac{1}{2}\delta_m - \frac{k^2y}{3k_c^2}\Phi; \quad y\delta_m'' + \frac{3}{2}\delta_m' = -\frac{k^2}{k_c^2}\Phi \quad (84)$$

These have a simple solution which we found earlier (see Eq. (69)):

$$\Phi = \Phi_\infty = \text{const.}; \quad \delta_m = -2\Phi_\infty - \frac{2k^2}{3k_c^2}\Phi_\infty y \sim y \quad (85)$$

In this limit, the matter perturbations grow linearly with expansion: $\delta_m \propto y \propto a$. In fact this is the most dominant growth mode in the linear perturbation theory.

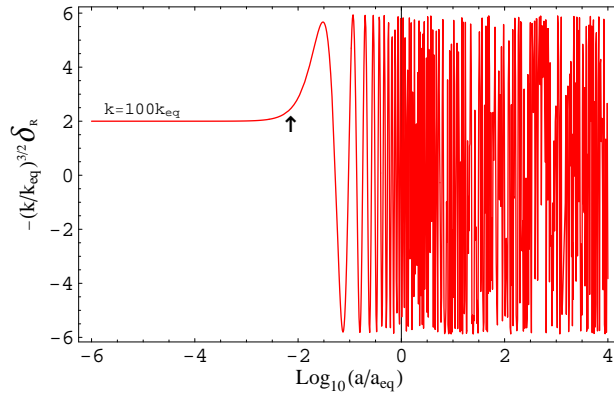


FIG. 3: Evolution of δ_R for a mode with $k = 100k_{eq}$. The mode remains frozen outside the Hubble radius at $(k/k_{eq})^{3/2}(-\delta_R) \approx (k/k_{eq})^{3/2}2\Phi = 2$ (in the normalisation used in Fig. 2) and oscillates when it enters the Hubble radius. The oscillations are well described by Eq. (80) with an amplitude of 6.

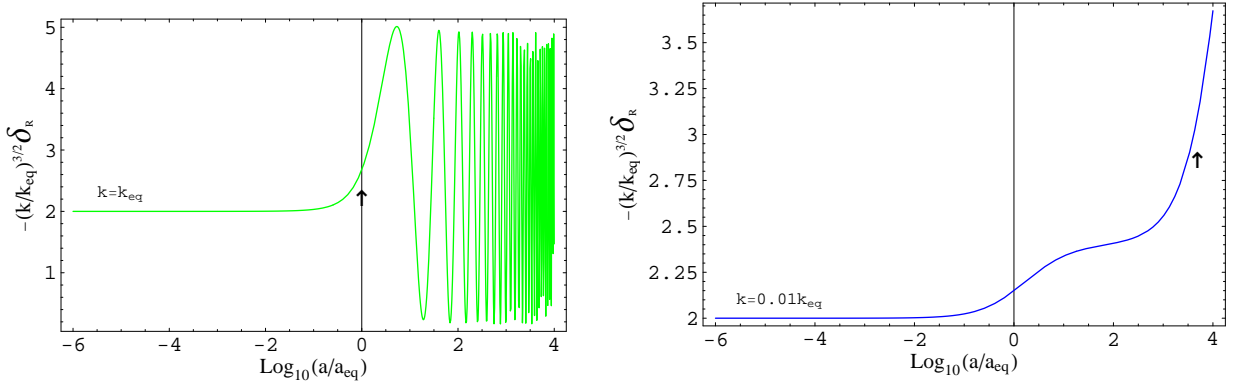


FIG. 4: Evolution of δ_R for two modes $k = k_{\text{eq}}$ and $k = 0.01 k_{\text{eq}}$. The modes remain frozen outside the Hubble radius ($-\delta_R \approx 2$) and oscillates when it enters the Hubble radius. The mode in the right panel stays outside the Hubble radius for most part of its evolution and hence changes very little.

D. An alternative description of matter-radiation system

Before proceeding further, we will describe an alternative procedure for discussing the perturbations in dark matter and radiation, which has some advantages. In the formalism we used above, we used perturbations in the energy density of radiation (δ_R) and matter (δ_m) as the dependent variables. Instead, we now use perturbations in the *total* energy density, δ and the perturbations in the entropy per particle, σ as the new dependent variables. In terms of δ_R, δ_m , these variables are defined as:

$$\delta \equiv \frac{\delta \rho_{\text{total}}}{\rho_{\text{total}}} = \frac{\rho_R \delta_R + \rho_m \delta_m}{\rho_R + \rho_m} = \frac{\delta_R + y \delta_m}{1 + y}; \quad y = \frac{\rho_m}{\rho_R} = \frac{a}{a_{\text{eq}}} \quad (86)$$

$$\sigma \equiv \left(\frac{\delta s}{s} \right) = \frac{3 \delta T_R}{T_R} - \frac{\delta \rho_m}{\rho_m} = \frac{3}{4} \delta_R - \delta_m = \frac{\delta n_R}{n_R} - \frac{\delta n_m}{n_m} \quad (87)$$

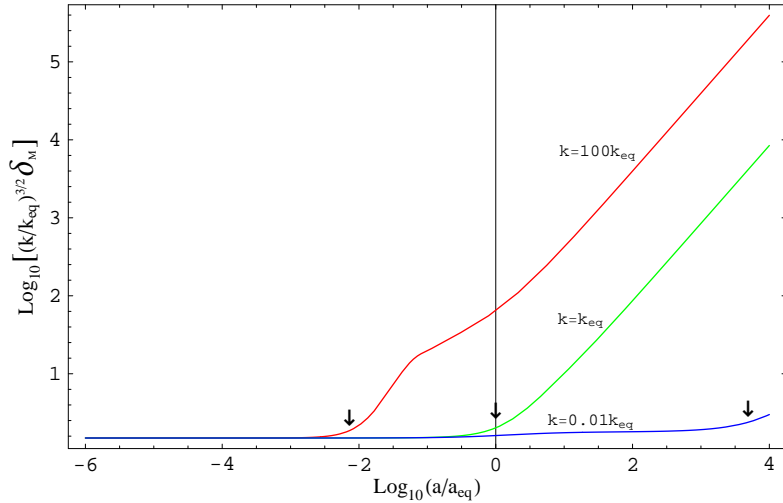


FIG. 5: Evolution of $|\delta_m|$ for 3 different modes. The modes are labelled by their wave numbers and the epochs at which they enter the Hubble radius are shown by small arrows. All the modes remain frozen when they are outside the Hubble radius and grow linearly in the matter dominated phase once they are inside the Hubble radius. The mode that enters the Hubble radius in the radiation dominated phase grows logarithmically until $y = y_{\text{eq}}$. These features are well approximated by Eqs. (83), (85).

Given the equations for δ_R, δ_m , one can obtain the corresponding equations for the new variables (δ, σ) by straight forward algebra. It is convenient to express them as two coupled equations for Φ and σ . After some direct but a bit tedious algebra, we get:

$$y\Phi'' + \frac{y\Phi'}{2(1+y)} + 3(1+c_s^2)\Phi' + \frac{3c_s^2\Phi}{4(1+y)} + c_s^2 \frac{k^2}{k_c^2} \frac{y}{1+y} \Phi = \frac{3c_s^2\sigma}{2(1+y)} \quad (88)$$

$$y\sigma'' + \frac{y\sigma'}{2(1+y)} + 3c_s^2\sigma' + \frac{3c_s^2y^2}{4(1+y)} \frac{k^2}{k_c^2} \sigma = \frac{c_s^2y^3}{2(1+y)} \left(\frac{k}{k_c}\right)^4 \Phi \quad (89)$$

where we have defined

$$c_s^2 = \frac{(4/3)\rho_R}{4\rho_R + 3\rho_m} = \frac{1}{3} \left(1 + \frac{3}{4} \frac{\rho_m}{\rho_R}\right)^{-1} = \frac{1}{3} \left(1 + \frac{3}{4}y\right)^{-1} \quad (90)$$

These equations show that the entropy perturbations and gravitational potential (which is directly related to total energy density perturbations) act as sources for each other. The coupling between the two arises through the right hand sides of Eq. (88) and Eq. (89). We also see that if we set $\sigma = 0$ as an initial condition, this is preserved to $\mathcal{O}(k^4)$ and — for long wave length modes — the Φ evolves independent of σ . The solutions to the coupled equations obtained by numerical integration is shown in Fig.(2) right panel. The entropy perturbation $\sigma \approx 0$ till the mode enters Hubble radius and grows afterwards tracking either δ_R or δ_m whichever is the dominant energy density perturbation. To illustrate the behaviour of Φ , let us consider the adiabatic perturbations at large scales with $\sigma \approx 0, k \rightarrow 0$; then the gravitational potential satisfies the equation:

$$y\Phi'' + \frac{y\Phi'}{2(1+y)} + 3(1+c_s^2)\Phi' + \frac{3c_s^2\Phi}{4(1+y)} = \frac{3c_s^2\sigma}{2(1+y)} \approx 0 \quad (91)$$

which has the two independent solutions:

$$f_1(y) = 1 + \frac{2}{9y} - \frac{8}{9y^2} - \frac{16}{9y^3}, \quad f_2(y) = \frac{\sqrt{1+y}}{y^2} \quad (92)$$

both of which diverge as $y \rightarrow 0$. We need to combine these two solutions to find the general solution, keeping in mind that the general solution should be nonsingular and become a constant (say, unity) as $y \rightarrow 0$. This fixes the linear combination uniquely:

$$f(y) = \frac{9}{10}f_1 + \frac{8}{5}f_2 = \frac{1}{10y^3} \left[16\sqrt{(1+y)} + 9y^3 + 2y^2 - 8y - 16\right] \quad (93)$$

Multiplying by Φ_i we get the solution that was found earlier (see Eq. (78)). Given the form of Φ and $\sigma \simeq 0$ we can determine all other quantities. In particular, we get:

$$\delta_R = \frac{-2(1+y)d(y\Phi)/dy + y\sigma}{1 + (3/4)y} \simeq -\frac{2(1+y)}{1 + (3/4)y} \frac{d}{dy}(y\Phi) \quad (94)$$

The corresponding velocity field, which we quote for future reference, is given by:

$$v_\alpha = -\frac{3c_s^2}{2(\dot{a}/a)}(1+y)\nabla_\alpha \frac{d(y\Phi)}{dy} \quad (95)$$

We conclude this section by mentioning another useful result related to Eq. (88). When $\sigma \approx 0$, the equation for Φ can be re-expressed as

$$a \frac{d\zeta}{da} = -\frac{2c_s^2}{3} \frac{k^2/a^2}{H^2} \frac{\rho}{\rho+p} \Phi \approx 0 \quad (\text{for } \frac{k}{aH} \ll 1) \quad (96)$$

where we have defined:

$$\zeta = \frac{2}{3} \frac{\rho}{\rho+p} \frac{a}{\dot{a}} \left(\dot{\Phi} + \frac{\dot{a}}{a}\Phi\right) + \Phi = \frac{H}{\rho+p} \frac{ik^\alpha}{k^2} \delta T_\alpha^0 + \Phi \quad (97)$$

(The i factor arises because of converting a gradient to the \mathbf{k} space; of course, when everything is done correctly, all physical quantities will be real.) Other equivalent alternative forms for ζ , which are useful are:

$$\zeta = \frac{2}{3[1+w(a)]} \frac{d}{da}(a\Phi) + \Phi = \frac{H^2}{a(\rho+p)} \frac{d}{dt} \left(\frac{a\Phi}{H} \right) \quad (98)$$

For modes which are bigger than the Hubble radius, Eq. (96) shows that ζ is conserved. When ζ =constant, we can integrate Eq. (98) easily to obtain:

$$\Phi = c_1 \frac{H}{a} + c_2 \left[1 - \frac{H}{a} \int_0^a \frac{da'}{H(a')} \right] \quad (99)$$

This is the easiest way to obtain the solution in Eq. (78).

The conservation law for ζ also allows us to understand in a simple manner our previous result that Φ only decreases by a factor (9/10) when the mode remains bigger than Hubble radius as we evolve the equations from $y \ll 1$ to $y \gg 1$. Let us compare the values of ζ early in the radiation dominated phase and late in the matter dominated phase. From the first equation in Eq. (98), [using $\Phi' \approx 0$] we find that, in the radiation dominated phase, $\zeta \approx (1/2)\Phi_i + \Phi_i = (3/2)\Phi_i$; late in the matter dominated phase, $\zeta \approx (2/3)\Phi_f + \Phi_f = (5/3)\Phi_f$. Hence the conservation of ζ gives $\Phi_f = (3/5)(3/2)\Phi_i = (9/10)\Phi_i$ which was the result obtained earlier. The expression in Eq. (99) also works at late times in the Λ dominated or curvature dominated universe.

One key feature which should be noted in the study of linear perturbation theory is the different amount of growths for Φ , δ_R and δ_m . The Φ either changes very little or decays; the δ_R grows in amplitude only by a factor of few. The physical reason, of course, is that the amplitude is frozen at super-Hubble scales and the pressure prevents the growth at sub-Hubble scales. In contrast, δ_m , which is pressureless, grows logarithmically in the radiation dominated era and linearly during the matter dominated era. Since the later phase lasts for a factor of 10^4 in expansion, we get a fair amount of growth in δ_m .

V. TRANSFER FUNCTION FOR MATTER PERTURBATIONS

We now have all the ingredients to evolve the matter perturbation from an initial value $\delta = \delta_i$ at $y = y_i \ll 1$ to the current epoch $y = y_0 = a_{\text{eq}}^{-1}$ in the matter dominated phase at $y \gg 1$. Initially, the wavelength of the perturbation will be bigger than the Hubble radius and the perturbation will essentially remain frozen. When it enters the Hubble radius in the radiation dominated phase, it begins to grow but only logarithmically (see section IV B) until the universe becomes matter dominated. In the final matter dominated phase, the perturbation grows linearly with expansion factor. The relation between final and initial perturbation can be obtained by combining these results.

Usually, one is more interested in the power spectrum $P_k(t)$ and the power per logarithmic band in k -space $\Delta_k(t)$. These quantities are defined in terms of $\delta_k(t)$ through the equations:

$$P_k(t) \equiv |\delta_k(t)|^2; \quad \Delta_k^2(t) \equiv \frac{k^3 P_k(t)}{2\pi^2} \quad (100)$$

It is therefore convenient to study the evolution of $k^{3/2}\delta_k$ since its square will immediately give the power per logarithmic band Δ_k^2 in k -space.

Let us first consider a mode which enters the Hubble radius in the radiation dominated phase at the epoch a_{enter} . From the scaling relation, $a_{\text{ent}}/k \propto t_{\text{ent}} \propto a_{\text{ent}}^2$ we find that $y_{\text{ent}} = (k_{\text{eq}}/k)$. Hence

$$k^{3/2}\delta_m(k, a=1) = \underbrace{\frac{1}{a_{\text{eq}}}}_{\text{MD}} \underbrace{\ln\left(\frac{a_{\text{eq}}}{a_{\text{ent}}}\right)}_{\text{RD}} \underbrace{\left[k^{3/2}\delta_{\text{ent}}(k) \right]}_{\text{at entry}} \propto \ln\left(\frac{k}{k_{\text{eq}}}\right) [k^{3/2}\delta_{\text{ent}}(k)] \quad (101)$$

where two factors — as indicated — gives the growth in radiation (RD) and matter dominated (MD) phases. Let us next consider the modes that enter in the matter dominated phase. In this case, $a_{\text{ent}}/k \propto t_{\text{ent}} \propto a_{\text{ent}}^{3/2}$ so that $y_{\text{ent}} = (k_{\text{eq}}/k)^2$. Hence

$$k^{3/2}\delta_m(k, a=1) = \underbrace{\frac{1}{a_{\text{ent}}}}_{\text{MD}} \underbrace{\left[k^{3/2}\delta_{\text{ent}}(k) \right]}_{\text{at entry}} \propto k^2 [k^{3/2}\delta_{\text{ent}}(k)] \quad (102)$$

To proceed further, we need to know the k -dependence of the perturbation when it enters the Hubble radius which, of course, is related to the mechanism that generates the initial power spectrum. The most natural choice will be that all the modes enter the Hubble radius with a constant amplitude at the time of entry. This would imply that the physical perturbations are scale invariant at the time of entering the Hubble radius, a possibility that was suggested by Zeldovich and Harrison [18] (years before inflation was invented!). We will see later that this is also true for perturbations generated by inflation and thus is a reasonable assumption at least in such models. Hence we shall assume

$$k^3 |\delta_{ent}(k)|^2 = k^3 P_{ent}(k) = C = \text{constant}, \quad (103)$$

Using this we find that the current value of perturbation is given by

$$P(k, a=1) \propto \left| \delta_m(k, a=1) \right|^2 \propto \begin{cases} k & (\text{for } k \ll k_{eq}) \\ k^{-3} (\ln k)^2 & (\text{for } k \gg k_{eq}) \end{cases} \quad (104)$$

The corresponding power per logarithmic band is

$$\Delta^2(k, a=1) \propto k^3 \left| \delta_m(k, a=1) \right|^2 \propto \begin{cases} k^4 & (\text{for } k \ll k_{eq}) \\ (\ln k)^2 & (\text{for } k \gg k_{eq}) \end{cases} \quad (105)$$

The form for $P(k)$ shows that the evolution imprints the scale k_{eq} on the power spectrum even though the initial power spectrum is scale invariant. For $k < k_{eq}$ (for large spatial scales), the primordial form of the spectrum is preserved and the evolution only increases the amplitude preserving the shape. For $k > k_{eq}$ (for small spatial scales), the shape is distorted and in general the power is suppressed in comparison with larger spatial scales. This arises because modes with small wavelengths enter the Hubble radius early on and have to wait till the universe becomes matter dominated in order to grow in amplitude. This is in contrast to modes with large wavelengths which continue to grow. It is this effect which suppresses the power at small wavelengths (for $k > k_{eq}$) relative to power at larger wavelengths.

VI. TEMPERATURE ANISOTROPIES OF CMBR

We shall now apply the formalism we have developed to understand the temperature anisotropies in the cosmic microwave background radiation which is probably the most useful application of *linear* perturbation theory. We shall begin by developing the general formulation and the terminology which is used to describe the temperature anisotropies.

Towards every direction in the sky, $\mathbf{n} = (\theta, \psi)$ we can define a fractional temperature fluctuation $\Delta(\mathbf{n}) \equiv (\Delta T/T)(\theta, \psi)$. Expanding this quantity in spherical harmonics on the sky plane as well as in terms of the spatial Fourier modes, we get the two relations:

$$\Delta(\mathbf{n}) \equiv \frac{\Delta T}{T}(\theta, \psi) = \sum_{l,m} a_{lm} Y_{lm}(\theta, \psi) = \int \frac{d^3 k}{(2\pi)^3} \Delta(\mathbf{k}) e^{i\mathbf{k}\cdot\mathbf{n}L} \quad (106)$$

where $L = \eta_0 - \eta_{LSS}$ is the distance to the last scattering surface (LSS) from which we are receiving the radiation. The last equality allows us to define the expansion coefficients a_{lm} in terms of the temperature fluctuation in the Fourier space $\Delta(\mathbf{k})$. Standard identities of mathematical physics now give

$$a_{lm} = \int \frac{d^3 k}{(2\pi)^3} (4\pi) i^l \Delta(\mathbf{k}) j_l(kL) Y_{lm}(\hat{\mathbf{k}}) \quad (107)$$

Next, let us consider the angular correlation function of temperature anisotropy, which is given by:

$$\mathcal{C}(\alpha) = \langle \Delta(\mathbf{n}) \Delta(\mathbf{m}) \rangle = \sum_l \sum_{m,m'} \langle a_{lm} a_{l'm'}^* \rangle Y_{lm}(\mathbf{n}) Y_{l'm'}^*(\mathbf{m}). \quad (108)$$

where the wedges denote an ensemble average. For a Gaussian random field of fluctuations we can express the ensemble average as $\langle a_{lm} a_{l'm'}^* \rangle = C_l \delta_{ll'} \delta_{mm'}$. Using Eq. (107), we get a relation between C_l and $\Delta(k)$. Given $\Delta(k)$, the C_l 's are given by:

$$C_l = \frac{2}{\pi} \int_0^\infty k^2 dk |\Delta(k)|^2 j_l^2(kL) \quad (109)$$

Further, Eq. (108) now becomes:

$$\mathcal{C}(\alpha) = \sum_l \frac{(2l+1)}{4\pi} C_l P_l(\cos \alpha) \quad (110)$$

Equation (110) shows that the mean-square value of temperature fluctuations and the quadrupole anisotropy corresponding to $l = 2$ are given by

$$\left(\frac{\Delta T}{T}\right)_{\text{rms}}^2 = \mathcal{C}(0) = \frac{1}{4\pi} \sum_{l=2}^{\infty} (2l+1) C_l, \quad \left(\frac{\Delta T}{T}\right)_Q^2 = \frac{5}{4\pi} C_2. \quad (111)$$

These can be explicitly computed if we know $\Delta(k)$ from the perturbation theory. (The motion of our local group through the CMBR leads to a large $l = 1$ dipole contribution in the temperature anisotropy. In the analysis of CMBR anisotropies, this is usually subtracted out. Hence the leading term is the quadrupole with $l = 2$.)

It should be noted that, for a given l , the C_l is the average over all $m = -l, \dots, -1, 0, 1, \dots, l$. For a Gaussian random field, one can also compute the variance around this mean value. It can be shown that this variance in C_l is $2C_l^2/(2l+1)$. In other words, there is an intrinsic root-mean-square fluctuation in the observed, mean value of C_l 's which is of the order of $\Delta C_l/C_l \approx (2l+1)^{-1/2}$. It is not possible for any CMBR observations which measures the C_l 's to reduce its uncertainty below this intrinsic variance — usually called the “cosmic variance”. For large values of l , the cosmic variance is usually sub-dominant to other observational errors but for low l this is the dominant source of uncertainty in the measurement of C_l 's. Current WMAP observations are indeed only limited by cosmic variance at low- l .

As an illustration of the formalism developed above, let us compute the C_l 's for low l which will be contributed essentially by fluctuations at large spatial scales. Since these fluctuations will be dominated by gravitational effects, we can ignore the complications arising from baryonic physics and compute these using the formalism we have developed earlier.

We begin by noting that the redshift law of photons in the unperturbed Friedmann universe, $\nu_0 = \nu(a)/a$, gets modified to the form $\nu_0 = \nu(a)/[a(1 + \Phi)]$ in a perturbed FRW universe. The argument of the Planck spectrum will thus scale as

$$\frac{\nu_0}{T_0} = \frac{\nu(a)}{aT_0(1 + \Phi)} = \frac{\nu(a)}{a\langle T_0 \rangle [1 + (\delta_R/4)](1 + \Phi)} \simeq \frac{\nu(a)}{a\langle T_0 \rangle [1 + \Phi + (\delta_R/4)]} \quad (112)$$

This is equivalent to a temperature fluctuation of the amount

$$\left(\frac{\Delta T}{T}\right)_{\text{obs}} = \frac{1}{4} \delta_R + \Phi \quad (113)$$

at large scales. (Note that the observed $\Delta T/T$ is not just $(\delta_R/4)$ as one might have naively imagined.) To proceed further, we recall our large scale solution (see Eq. (78)) for the gravitational potential:

$$\Phi = \Phi_i \frac{1}{10y^3} \left[16\sqrt{(1+y)} + 9y^3 + 2y^2 - 8y - 16 \right]; \quad \delta_R = 4\Phi - 6\Phi_i \quad (114)$$

At $y = y_{dec}$ we can take the asymptotic solution $\Phi_{dec} \approx (9/10)\Phi_i$. Hence we get

$$\left(\frac{\Delta T}{T}\right)_{\text{obs}} = \left[\frac{1}{4} \delta_R + \Phi \right]_{dec} = 2\Phi_{dec} - \frac{3}{2} \Phi_i \approx 2\Phi_{dec} - \frac{3}{2} \frac{10}{9} \Phi_{dec} = \frac{1}{3} \Phi_{dec} \quad (115)$$

We thus obtain the nice result that the observed temperature fluctuations at very large scales is simply related to the fluctuations of the gravitational potential at these scales. (For a discussion of the $1/3$ factor, see [19]). Fourier transforming this result we get $\Delta(\mathbf{k}) = (1/3)\Phi(\mathbf{k}, \eta_{LSS})$ where η_{LSS} is the conformal time at the last scattering surface. (This contribution is called Sachs-Wolfe effect.) It follows from Eq. (109) that the contribution to C_l from the gravitational potential is

$$C_l = \frac{2}{\pi} \int k^2 dk |\Delta(k)|^2 j_l^2(kL) = \frac{2}{\pi} \int_0^{\infty} \frac{dk}{k} \frac{k^3 |\Phi_k|^2}{9} j_l^2(kL) \quad (116)$$

with

$$L = \eta_0 - \eta_{LSS} \approx \eta_0 \approx 2(\Omega_m H_0^2)^{-1/2} \approx 6000 \Omega_m^{-1/2} h^{-1} \text{ Mpc} \quad (117)$$

For a scale invariant spectrum, $k^3|\Phi_k|^2$ is a constant independent of k . (Earlier on, in Eq. (103) we said that scale invariant spectrum has $k^3|\delta_k|^2 = \text{constant}$. These statements are equivalent since $\delta \approx -2\Phi$ at the large scales because of Eq. (85) with the extra correction term in Eq. (85) being about 3×10^{-4} for $k \approx L^{-1}$, $y = y_{dec}$.) As we shall see later, inflation generates such a perturbation. In this case, it is conventional to introduce a constant amplitude A and write:

$$\Delta_{\Phi}^2 \equiv \frac{k^3|\Phi_k|^2}{2\pi^2} = A^2 = \text{constant} \quad (118)$$

Substituting this form into Eq. (116) and evaluating the integral, we find that

$$\frac{l(l+1)C_l}{2\pi} = \left(\frac{A}{3}\right)^2 \quad (119)$$

As an application of this result, let us consider the observations of COBE which measured the temperature fluctuations for the first time in 1992. This satellite obtained the RMS fluctuations and the quadrupole after smoothing over an angular scale of about $\theta_c \approx 10^\circ$. Hence the observed values are slightly different from those in Eq. (111). We have, instead,

$$\left(\frac{\Delta T}{T}\right)_{\text{rms}}^2 = \frac{1}{4\pi} \sum_{l=2}^{\infty} (2l+1)C_l \exp\left(-\frac{l^2\theta_c^2}{2}\right); \quad \left(\frac{\Delta T}{T}\right)_Q^2 = \frac{5}{4\pi} C_2 e^{-2\theta_c^2}. \quad (120)$$

Using Eqs. (118), (119) we find that

$$\left(\frac{\Delta T}{T}\right)_Q \cong 0.22A; \quad \left(\frac{\Delta T}{T}\right)_{\text{rms}} \cong 0.51A. \quad (121)$$

Given these two measurements, one can verify that the fluctuations are consistent with the scale invariant spectrum by checking their ratio. Further, the numerical value of the observed $(\Delta T/T)$ can be used to determine the amplitude A . One finds that $A \approx 3 \times 10^{-5}$ which sets the scale of fluctuations in the gravitational potential at the time when the perturbation enters the Hubble radius.

Incidentally, note that the solution $\delta_R = 4\Phi - 6\Phi_i$ corresponds to $\delta_m = (3/4)\delta_R = 3\Phi - (9/2)\Phi_i$. At $y = y_{dec}$, taking $\Phi_{dec} = (9/10)\Phi_i$, we get $\delta_m = 3\Phi_{dec} - (9/2)(10/9)\Phi_{dec} = -2\Phi_{dec}$. Since $(\Delta T/T)_{obs} = (1/3)\Phi_{dec}$ we get $\delta_m = -6(\Delta T/T)_{obs}$. This shows that the amplitude of matter perturbations is a factor six larger than the amplitude of temperature anisotropy for our adiabatic initial conditions. In several other models, one gets $\delta_m = \mathcal{O}(1)(\Delta T/T)_{obs}$. So, to reach a given level of nonlinearity in the matter distribution at later times, these models will require higher values of $(\Delta T/T)_{obs}$ at decoupling. This is one reason for such models to be observationally ruled out.

There is another useful result which we can obtain from Eq. (109) along the same lines as we derived the Sachs-Wolfe effect. Whenever $k^3|\Delta(k)|^2$ is a slowly varying function of k , we can pull out this factor out of the integral and evaluate the integral over j_l^2 . This will give the result for any slowly varying $k^3|\Delta(k)|^2$

$$\frac{l(l+1)C_l}{2\pi} \approx \left(\frac{k^3|\Delta(k)|^2}{2\pi^2}\right)_{kL \approx l} \quad (122)$$

This is applicable even when different processes contribute to temperature anisotropies as long as they add in quadrature. While far from accurate, it allows one to estimate the effects rapidly.

A. CMBR Temperature Anisotropy: More detailed theory

We shall now work out a more detailed theory of temperature anisotropies of CMBR so that one can understand the effects at small scales as well. A convenient starting point is the distribution function for photons with perturbed Planckian distribution, which we can write as:

$$f(x^\alpha, \eta, E, n^\alpha) = \frac{I_\nu}{2\pi\nu^3} = f_P\left(\frac{aE}{1+\Delta}\right); \quad f_P(\epsilon) \equiv 2[\exp(\epsilon/T_0) - 1]^{-1} \quad (123)$$

The $f_P(\epsilon)$ is the standard Planck spectrum for energy ϵ and we take $\epsilon = aE(1+\Delta)^{-1}$ to take care of the perturbations. In the absence of collisions, the distribution function is conserved along the trajectories of photons so that $df/d\eta = 0$. So, in the presence of collisions, we can write the time evolution of the distribution function as

$$\frac{df}{d\eta} = \left(\frac{aE}{1+\Delta}\right) f'_P\left(\frac{aE}{1+\Delta}\right) \left[\frac{d\ln(aE)}{d\eta} - \frac{d\Delta}{d\eta}\right] = \left(\frac{df}{d\eta}\right)_{\text{coll}} \quad (124)$$

where the right hand side gives the contribution due to collisional terms. Equivalently, in terms of Δ , the same equation takes the form:

$$\frac{d\Delta}{d\eta} - \frac{d\ln(aE)}{d\eta} = - \left(\frac{1 + \Delta}{aE} \right) [f'_P]^{-1} \left(\frac{df}{d\eta} \right)_{\text{coll}} \equiv \left(\frac{d\Delta}{d\eta} \right)_{\text{coll}} \quad (125)$$

To proceed further, we need the expressions for the two terms on the left hand side. First term, on using the standard expansion for total derivative, gives:

$$\frac{d\Delta}{d\eta} = \frac{\partial\Delta}{\partial\eta} + \underbrace{\frac{\partial\Delta}{\partial x^\alpha} \frac{dx^\alpha}{d\eta}}_{n^\alpha} + \underbrace{\frac{\partial\Delta}{\partial E} \frac{dE}{d\eta}}_{\text{zero}} + \underbrace{\frac{\partial\Delta}{\partial n^\alpha} \frac{dn^\alpha}{d\eta}}_{\mathcal{O}(\Delta^2)=0} \cong \partial_\eta\Delta + n^\alpha\partial_\alpha\Delta \quad (126)$$

(Note that we are assuming $\partial\Delta/\partial E = 0$ so that the perturbations do not depend on the frequency of the photon.) To determine the second term, we note that it vanishes in the unperturbed Friedmann universe and arises essentially due to the variation of Φ . Both the intrinsic time variation of Φ as well as its variation along the photon path will contribute, giving:

$$\frac{d\ln(aE)}{d\eta} = -n^\alpha\partial_\alpha\Phi + \partial_\eta\Phi \quad (127)$$

(The minus sign arises from the fact that we have $(1 + 2\Phi)$ in g_{00} but $(1 - 2\Phi)$ in the spatial perturbations.) Putting all these together, we can bring the evolution equation Eq. (125) to the form:

$$\frac{d\Delta}{d\eta} = -n^\alpha\partial_\alpha\Phi + \partial_\eta\Phi + \left(\frac{d\Delta}{d\eta} \right)_{\text{coll}} \quad (128)$$

Let us next consider the collision term, which can be expressed in the form:

$$\begin{aligned} \left(\frac{d\Delta}{ad\eta} \right)_{\text{coll}} &= -N_e\sigma_T\Delta + N_e\sigma_T \left(\frac{1}{4}\delta_R \right) + N_e\sigma_T(\mathbf{v}\cdot\mathbf{n}) \\ &= N_e\sigma_T \left(-\Delta + \frac{1}{4}\delta_R + \mathbf{v}\cdot\mathbf{n} \right) \end{aligned} \quad (129)$$

Each of the terms in the right hand side of the first line has a simple interpretation. The first term describes the removal of photons from the beam due to Thomson scattering with the electrons while the second term gives the scattering contribution into the beam. In a static universe, we expect these two terms to cancel if $\Delta = (1/4)\delta_R$ which fixes the relative coefficients of these two terms. The third term is a correction due to the fact that the electrons which are scattering the photons are not at rest relative to the cosmic frame. This leads to a Doppler shift which is accounted for by the third term. (We denote electron number *density* by N_e rather than n_e to avoid notational conflict with n^α .)

Formally, Eq. (128) is a first order linear differential equation for Δ . To eliminate the $-N_e\sigma_T\Delta$ term which is linear in Δ in the right hand side, we use the standard integrating factor $\exp(-\tau)$ where

$$\tau(\chi) \equiv \int_0^\chi d\eta (aN_e\sigma_T) \quad (130)$$

We can then formally integrate Eq. (128) to get:

$$\Delta(\mathbf{n}) = \int_0^{\eta_0} d\chi e^{-\tau(\chi)} \left[-n^\alpha\partial_\alpha\Phi + \partial_\eta\Phi + aN_e\sigma_T \left(\frac{1}{4}\delta_R + \mathbf{v}\cdot\mathbf{n} \right) \right] \quad (131)$$

We can write

$$e^{-\tau}(-n^\alpha\partial_\alpha\Phi) = - \left(\frac{d\Phi}{d\eta} \right) e^{-\tau} + (\partial_\eta\Phi)e^{-\tau} = -\frac{d}{d\eta}(\Phi e^{-\tau}) + (aN_e\sigma_T\Phi)e^{-\tau} + (\partial_\eta\Phi)e^{-\tau} \quad (132)$$

On integration, the first term gives zero at the lower limit and an unimportant constant (which does not depend on \mathbf{n}). Using the rest of the terms, we can write Eq. (131) in the form:

$$\begin{aligned} \Delta(\mathbf{n}) &= \int_0^{\eta_0} d\chi e^{-\tau} \left[2\partial_\eta\Phi + aN_e\sigma_T \left(\Phi + \frac{1}{4}\delta_R + \mathbf{v}\cdot\mathbf{n} \right) \right] \\ &= \int_0^{\eta_0} d\chi e^{-\tau} [2\partial_\eta\Phi] + \int_0^{\eta_0} d\chi (e^{-\tau}aN_e\sigma_T) \left(\Phi + \frac{1}{4}\delta_R + \mathbf{v}\cdot\mathbf{n} \right) \end{aligned} \quad (133)$$

The first term gives the contribution due to the intrinsic time variation of the gravitational potential along the path of the photon and is called the integrated Sachs-Wolfe effect. In the second term one can make further simplifications. Note that $e^{-\tau}$ is essentially unity (optically thin) for $z < z_{\text{rec}}$ and zero (optically thick) for $z > z_{\text{rec}}$; on the other hand, $N_e \sigma_T$ is zero for $z < z_{\text{rec}}$ (all the free electrons have disappeared) and is large for $z > z_{\text{rec}}$. Hence the product ($a N_e e^{-\tau}$) is sharply peaked at $\chi = \chi_{\text{rec}}$ (i.e. at $z \simeq 10^3$ with $\Delta z \simeq 80$). Treating this sharply peaked quantity as essentially a Dirac delta function (usually called the instantaneous recombination approximation) we can approximate the second term in Eq. (133) as a contribution occurring just on the LSS:

$$\Delta(\mathbf{n}) = \left(\frac{1}{4} \delta_R + \mathbf{v} \cdot \mathbf{n} + \Phi \right)_{\text{LSS}} + 2 \int_{\eta_{\text{LSS}}}^{\eta_0} d\chi \partial_\eta \Phi \quad (134)$$

In the second term we have put $\tau = \infty$ for $\eta < \eta_{\text{LSS}}$ and $\tau = 0$ for $\eta > \eta_{\text{LSS}}$.

Once we know δ_R, Φ and \mathbf{v} on the LSS from perturbation theory, we can take a Fourier transform of this result to obtain $\Delta(k)$ and use Eq. (109) to compute C_l . At very large scales the velocity term is sub-dominant and we get back the Sachs-Wolfe effect derived earlier in Eq. (118). For understanding the small scale effects, we need to introduce baryons into the picture which is our next task.

B. Description of photon-baryon fluid

To study the interaction of photons and baryons in the fluid limit, we need to again start from the continuity equation and Euler equation. In Fourier space, the continuity equation is same as the one we had before (see Eq. (52)):

$$\left(\frac{3}{4} \right) \dot{\delta}_R = -k^2 v_R + 3\dot{\Phi}; \quad \dot{\delta}_B = -k^2 v_B + 3\dot{\Phi} \quad (135)$$

The Euler equations, however, gets modified; for photons, it becomes:

$$\dot{v}_R = \left(\frac{1}{4} \dot{\delta}_R + \dot{\Phi} \right) - \dot{\tau} (v_R - v_B); \quad \dot{\tau} = N_e \sigma_T a \quad (136)$$

The first two terms in the right hand side are exactly the same as the ones in Eq. (55). The last term is analogous to a viscous drag force between the photons and baryons which arises because of the non zero relative velocity between the two fluids. The coupling is essentially due to Thomson scattering which leads to the factor $\dot{\tau}$. (The notation, and the physics, is the same as in Eq. (130)). The corresponding Euler equation for the baryons is:

$$\dot{v}_B = -\frac{\dot{a}}{a} v_B + \dot{\Phi} + \frac{\dot{\tau} (v_R - v_B)}{R} \quad (137)$$

where

$$R \equiv \frac{p_B + \rho_B}{p_R + \rho_R} \simeq \frac{3\rho_B}{4\rho_R} \approx 30 \Omega_B h^2 \left(\frac{a}{10^{-3}} \right) \quad (138)$$

Again, the first two terms in the right hand side of Eq. (137) are the same as what we had before in Eq. (55). The last term has the same interpretation as in the case of Euler equation Eq. (136) for photons, except for the factor R . This quantity essentially takes care of the inertia of baryons relative to photons. Note that the the conserved momentum density of photon-baryon fluid has the form

$$(\rho_R + p_R)v_R + (\rho_B + p_B)v_B \approx (1 + R)(\rho_R + p_R)v_R \quad (139)$$

which accounts for the extra factor R in Eq. (137).

We can now combine the Eqs. (135), (136), (137) to obtain, to lowest order in $(k/\dot{\tau})$ the equation:

$$\ddot{\delta}_R + \frac{\dot{R}}{(1+R)} \dot{\delta}_R + k^2 c_s^2 \delta_R = F \quad (140)$$

with

$$F = 4 \left[\ddot{\Phi} + \frac{\dot{R}}{(1+R)} \dot{\Phi} - \frac{1}{3} k^2 \Phi \right]; \quad c_s^2 = \frac{1}{3(1+R)} \quad (141)$$

An exact solution to this equation is difficult to obtain. However, we can try to understand several features by an approximate method in which we treat the time variation of R to be small. In that case, we can drop the \dot{R} terms on both sides of the equation. Since we know that the physically relevant temperature fluctuation is $\Delta = (1/4)\delta_R + \Phi$, we can recast the above equation for Δ as:

$$\ddot{\Delta} + k^2 c_s^2 \Delta \approx -k^2 c_s^2 R \Phi + 2\ddot{\Phi} \quad (142)$$

Let us further ignore the time variation of all terms (especially $\ddot{\Phi}$ on the right hand side). Then, the solution is just $\Delta = -R\Phi + A \cos(kc_s \eta_{\text{LSS}}) + B \sin(kc_s \eta_{\text{LSS}})$. To fix the initial conditions which determine A and B , we recall that early on ($\eta \rightarrow 0$), we have $\Delta \rightarrow \Phi/3$ (see Eq. (115)) and corresponding velocity should vanish. This gives the solution:

$$\frac{1}{4}\delta_R + \Phi = \frac{\Phi_i}{3} (1 + 3R) \cos(kc_s \eta_{\text{LSS}}) - \Phi_i R; \quad v = -\Phi_i (1 + 3R) c_s \sin(kc_s \eta_{\text{LSS}}) \quad (143)$$

(One can do a little better by using WKB approximation in which $(kc_s \eta_{\text{LSS}})$ can be replaced by the integral of kc_s over η but it is not very important.) Given this solution, one can proceed as before and compute the C_l 's. Adding the effects of $[\Phi + (1/4)\delta_R]$ and that of $[\mathbf{v} \cdot \mathbf{n}]$ in quadrature and noticing that the angular average of $\langle (\mathbf{v} \cdot \mathbf{n})^2 \rangle = (1/3)v^2$ we can estimate the C_l for scale invariant ($k^3 |\Phi_k|^2 = 2\pi^2 A^2$) spectrum to be:

$$l(l+1)C_l = 2\pi^2 A^2 \left\{ \left[\frac{(1+3R)}{3} \cos(k^* c_s \eta_{\text{LSS}}) - R \right]^2 + \frac{(1+3R)^2}{3} c_s^2 \sin^2(k^* c_s \eta_{\text{LSS}}) \right\} \quad (144)$$

with $k^* L \approx l$ with $L = \eta_0 - \eta_{\text{LSS}} \simeq \eta_0$. The key feature is, of course, the maxima and minima which arises from the trigonometric functions. The peaks of C_l are determined by the condition $k^* c_s \eta_{\text{LSS}} = lc_s \eta_{\text{LSS}} / \eta_0 = n\pi$; that is

$$l_{\text{peak}} = \frac{n\pi}{c_s} \left(\frac{\eta_0}{\eta_{\text{LSS}}} \right) = n\pi \sqrt{3} (1 + z_{\text{dec}})^{1/2} \approx 172n \quad (145)$$

More precise work gives the first peak at $l_{\text{peak}} \simeq 200$. It is also clear that because of non zero R the peaks are larger when the cosine term is negative; that is, the odd peaks corresponding to $n = 1, 3, \dots$ have larger amplitudes than the even peaks with $n = 2, 4, \dots$

Incidentally, the above approximation is not very good for modes which enter the Hubble radius during the radiation dominated phase since Φ does evolve with time (and decays) in the radiation dominated phase. We saw that $\Phi \approx -3\Phi_i (ly)^{-2} \cos(ly)$ asymptotically in this phase (see Eq. (80)). From Eq. (80) we find that during this phase, for modes which are inside the Hubble radius, we can take $\delta_R \approx 6\Phi_i \cos(ly)$, so that $\Delta \approx \delta_R/4 \approx (3/2)\Phi_i \cos(ly)$. On the other hand, at very large scale, the amplitude was $\Delta = \Phi/3 = (1/3)(9/10)\Phi_i = (3/10)\Phi_i$. Hence the amplitude of the modes that enter the horizon during the radiation dominated phase is enhanced by a factor $(3/2)(10/9) = 5$, relative to the large scale amplitude contributed by modes which enter during matter dominated phase. This is essentially due to the driving term $\ddot{\Phi}$ which is nonzero in the radiation dominated phase but zero in the matter dominated phase. (In reality, the enhancement is smaller because the relevant modes have $k \gtrsim k_{\text{eq}}$ rather than $k \gg k_{\text{eq}}$; see Figs. 3 and 4.)

If this were the whole story, we will see a series of peaks and troughs in the temperature anisotropies as a function of angular scale. In reality, however, there are processes which damp out the anisotropies at small angular scales (large $-l$) so that only the first few peaks and troughs are really relevant. We will now discuss two key damping mechanisms which are responsible for this.

The first one is the finite width of the last scattering surface which makes it uncertain from which event we are receiving the photons. In general, if $\mathcal{P}(z)$ is the probability that the photon was last scattered at redshift z , then we can write:

$$\left(\frac{\Delta T}{T} \right)_{\text{obs}} = \int dz \left\{ \begin{array}{l} (\Delta T/T) \text{ if the last} \\ \text{scattering was at } z \end{array} \right\} \times \mathcal{P}(z). \quad (146)$$

From Eq. (41) we know that $\mathcal{P}(z)$ is a Gaussian with width $\Delta z = 80$. This corresponds to a length scale

$$\Delta l = c \left(\frac{dt}{dz} \right) \Delta z \cdot (1 + z_{\text{dec}}) \approx H_0^{-1} \frac{\Delta z}{\Omega^{1/2} z_{\text{dec}}^{3/2}} \approx 8 (\Omega h^2)^{-1/2} \text{ Mpc.} \quad (147)$$

over which the temperature fluctuations will be smoothed out.

It turns out that there is another effect, which is slightly more important. This arises from the fact that the photon-baryon fluid is not tightly coupled and the photons can diffuse through the fluid. This diffusion can be modeled as a random walk and the root mean square distance traveled by the photon during this diffusion process will smear the temperature anisotropies over that length scale. This photon diffusion length scale can be estimated as follows:

$$(\Delta x)^2 = \underbrace{N}_{\text{number of collisions}} \underbrace{\left(\frac{q}{a}\right)^2}_{\text{comoving mean free path}} = \frac{\Delta t}{q(t)} \frac{q^2}{a^2} = \frac{\Delta t}{a^2} q(t) \quad (148)$$

Integrating, we find the mean square distance traveled by the photon to be

$$x^2 \equiv \int_0^{t_{\text{dec}}} \frac{dt}{a^2(t)} q(t) = \frac{3}{5} \frac{t_{\text{dec}} q(t_{\text{dec}})}{a^2(t_{\text{dec}})} \quad (149)$$

The corresponding proper length scale below which photon diffusion will wipe out temperature anisotropies is:

$$q_{\text{diff}} = a(t_{\text{dec}}) x = \left[\frac{3}{5} t_{\text{dec}} q(t_{\text{dec}}) \right]^{\frac{1}{2}} \simeq 35 \text{ Mpc} \left(\frac{\Omega_B h^2}{0.02} \right)^{-\frac{1}{2}} (\Omega h_{50}^2)^{-\frac{1}{4}}. \quad (150)$$

It turns out that this is the dominant sources of damping of temperature anisotropies at large $l \approx 10^3$.

VII. GENERATION OF INITIAL PERTURBATIONS FROM INFLATION

In the description of linear perturbation theory given above, we assumed that some small perturbations existed in the early universe which are amplified through gravitational instability. To provide a complete picture we need a mechanism for generation of these initial perturbations. One such mechanism is provided by inflationary scenario which allows for the quantum fluctuations in the field driving the inflation to provide classical energy density perturbations at a late epoch. (Originally inflationary scenarios were suggested as pseudo-solutions to certain pseudo-problems; that is only of historical interest today and the only reason to take the possibility of an inflationary phase in the early universe seriously is because it provides a mechanism for generating these perturbations.) We shall now discuss how this can come about.

The basic assumption in inflationary scenario is that the universe underwent a rapid — nearly exponential — expansion for a brief period of time in very early universe. The simplest way of realizing such a phase is to postulate the existence of a scalar field with a nearly flat potential. The dynamics of the universe, driven by a scalar field source, is described by:

$$\frac{1}{a^2} \left(\frac{da}{dt} \right)^2 = H^2(t) = \frac{1}{3M_{\text{Pl}}^2} \left[V(\phi) + \frac{1}{2} \left(\frac{d\phi}{dt} \right)^2 \right]; \quad \frac{d^2\phi}{dt^2} + 3H \frac{d\phi}{dt} = -\frac{dV}{d\phi} \quad (151)$$

where $M_{\text{pl}} = (8\pi G)^{-1/2}$. If the potential is nearly flat for certain range of ϕ , we can introduce the “slow roll-over” approximation, under which these equations become:

$$H^2 \simeq \frac{V(\phi)}{3M_{\text{Pl}}^2}; \quad 3H \frac{d\phi}{dt} \simeq -V'(\phi) \quad (152)$$

For this slow roll-over to last for reasonable length of time, we need to assume that the terms ignored in the Eq. (151) are indeed small. This can be quantified in terms of the parameters:

$$\epsilon(\phi) = \frac{M_{\text{Pl}}^2}{2} \left(\frac{V'}{V} \right)^2; \quad \eta(\phi) = M_{\text{Pl}}^2 \frac{V''}{V} \quad (153)$$

which are taken to be small. Typically the inflation ends when this assumption breaks down. If such an inflationary phase lasts up to some time t_{end} then the universe would have undergone an expansion by a factor $\exp N(t)$ during the interval (t, t_{end}) where

$$N \equiv \ln \frac{a(t_{\text{end}})}{a(t)} = \int_t^{t_{\text{end}}} H dt \simeq \frac{1}{M_{\text{Pl}}^2} \int_{\phi_{\text{end}}}^{\phi} \frac{V}{V'} d\phi \quad (154)$$

One usually takes $N \simeq 65$ or so.

Before proceeding further, we would like to make couple of comments regarding such an inflationary phase. To begin with, it is not difficult to obtain *exact* solutions for $a(t)$ with rapid expansion by tailoring the potential for the scalar field. In fact, given any $a(t)$ and thus a $H(t) = (\dot{a}/a)$, one can determine a potential $V(\phi)$ for a scalar field such that Eq. (151) are satisfied (see the first reference in [27]). One can verify that, this is done by the choice:

$$V(t) = \frac{1}{16\pi G} H \left[6H + \frac{2}{H} \frac{dH}{dt} \right]; \quad \phi(t) = \int dt \left[\frac{-2}{8\pi G} \frac{dH}{dt} \right]^{1/2} \quad (155)$$

Given any $H(t)$, these equations give $(\phi(t), V(t))$ and thus implicitly determine the necessary $V(\phi)$. As an example, note that a power law inflation, $a(t) = a_0 t^p$ (with $p \gg 1$) is generated by:

$$V(\phi) = V_0 \exp \left(-\sqrt{\frac{2}{p}} \frac{\phi}{M_{\text{Pl}}} \right) \quad (156)$$

while an exponential of power law

$$a(t) \propto \exp(At^f), \quad f = \frac{\beta}{4 + \beta}, \quad 0 < f < 1, \quad A > 0 \quad (157)$$

can arise from

$$V(\phi) \propto \left(\frac{\phi}{M_{\text{Pl}}} \right)^{-\beta} \left(1 - \frac{\beta^2}{6} \frac{M_{\text{Pl}}^2}{\phi^2} \right) \quad (158)$$

Thus generating a rapid expansion in the early universe is trivial if we are willing to postulate scalar fields with tailor made potentials. This is often done in the literature.

The second point to note regarding any inflationary scenarios is that the modes with reasonable size today originated from sub-Planck length scales early on. A scale λ_0 today will be

$$\lambda_{end} = \lambda_0 \frac{a_{end}}{a_0} = \lambda_0 \frac{T_0}{T_{end}} \approx \lambda_0 \times 10^{-28} \quad (159)$$

at the end of inflation (if inflation took place at GUT scales) and

$$\lambda_{begin} = \lambda_{end} A^{-1} \approx \lambda_0 \times 10^{-58} (A/10^{30})^{-1} \quad (160)$$

at the beginning of inflation if the inflation changed the scale factor by $A \simeq 10^{30}$. Note that $\lambda_{begin} < L_P$ for $\lambda_0 < 3$ Mpc!! Most structures in the universe today correspond to transplanckian scales at the start of the inflation. It is not clear whether we can trust standard physics at early stages of inflation or whether transplanckian effects will lead to observable effects [20, 21].

Let us get back to conventional wisdom and consider the evolution of perturbations in a universe which underwent exponential inflation. During the inflationary phase the $a(t)$ grows exponentially and hence the wavelength of any perturbation will also grow with it. The Hubble radius, on the other hand, will remain constant. It follows that, one can have situation in which a given mode has wavelength smaller than the Hubble radius at the beginning of the inflation but grows and becomes bigger than the Hubble radius as inflation proceeds. It is conventional to say that a perturbation of comoving wavelength λ_0 ‘‘leaves the Hubble radius’’ when $\lambda_0 a = d_H$ at some time $t = t_{exit}(\lambda_0)$. For $t > t_{exit}$ the wavelength of the perturbation is bigger than the Hubble radius. Eventually the inflation ends and the universe becomes radiation dominated. Then the wavelength will grow ($\propto t^{1/2}$) slower than the Hubble radius ($\propto t$) and will enter the Hubble radius again during $t = t_{enter}(\lambda_0)$. Our first task is to relate the amplitude of the perturbation at $t = t_{exit}(\lambda_0)$ with the perturbation at $t = t_{enter}(\lambda_0)$.

We know that for modes bigger than Hubble radius, we have the conserved quantity (see Eq. (97))

$$\zeta = \frac{2}{3} \frac{\rho}{\rho + p} \frac{a}{\dot{a}} \left(\dot{\Phi} + \frac{\dot{a}}{a} \Phi \right) + \Phi = \frac{H}{\rho + p} \frac{ik^\alpha}{k^2} \delta T_\alpha^0 + \Phi \quad (161)$$

At the time of re-entry, the universe is radiation dominated and $\zeta_{entry} \approx (2/3)\Phi$. On the other hand, during inflation, we can write the scalar field as a dominant homogeneous part plus a small, spatially varying fluctuation: $\phi(t, \mathbf{x}) = \phi_0(t) + f(t, \mathbf{x})$. Perturbing the equation in Eq. (151) for the scalar field, we find that the homogeneous mode ϕ_0 satisfies Eq. (151) while the perturbation, in Fourier space satisfies:

$$\frac{d^2 f_k}{dt^2} + 3H \frac{df_k}{dt} + \frac{k^2}{a^2} f_k = 0 \quad (162)$$

Further, the energy momentum tensor for the scalar field gives [with the “dot” denoting $(d/d\eta) = a(d/dt)$]:

$$\rho = \frac{\dot{\phi}_0^2}{2a^2} + V; \quad p = \frac{\dot{\phi}_0^2}{2a^2} - V; \quad \delta T_0^\alpha = \frac{ik^\alpha}{a} \dot{\phi}_0 f \quad (163)$$

It is easy to see that Φ is negligible at $t = t_{\text{exit}}$ since

$$\Phi \sim \frac{4\pi G a^2}{k^2} \sim \frac{4\pi G}{H^2} \delta\rho \sim \frac{\delta\rho}{\rho} \sim \left(\frac{\rho+p}{\rho}\right) \left(\frac{\delta\rho}{\rho+p}\right) \ll \frac{H}{(\rho+p)} \frac{\dot{\phi}_0}{a} f_k \quad (164)$$

Therefore,

$$\zeta_{\text{exit}} \approx \frac{H}{(\dot{\phi}_0^2/a^2)} \left[-\frac{\dot{\phi}_0 f_k}{a} \right] = -aH \frac{f_k}{\dot{\phi}_0} \approx \frac{3H^2}{V'} f_k \quad (165)$$

Using the conservation law $\zeta_{\text{exit}} = \zeta_{\text{entry}}$, we get

$$\Phi_k \Big|_{\text{entry}} = \frac{9H^2}{2V'} f_k \Big|_{\text{exit}} \quad (166)$$

Thus, given a perturbation of the scalar field f_k during inflation, we can compute its value at the time of re-entry, which — in turn — can be used to compare with observations.

Equation (166) connects a *classical* energy density perturbation f_k at the time of exit with the corresponding quantity Φ_k at the time of re-entry. The next important — and conceptually difficult — question is how we can obtain a *c-number* field f_k from a quantum scalar field. There is no simple answer to this question and one possible way of doing it is as follows: Let us start with the quantum operator for a scalar field decomposed into the Fourier modes with $\hat{q}_{\mathbf{k}}(t)$ denoting an infinite set of operators:

$$\hat{\phi}(t, \mathbf{x}) = \int \frac{d^3\mathbf{k}}{(2\pi)^3} \hat{q}_{\mathbf{k}}(t) e^{i\mathbf{k}\cdot\mathbf{x}}. \quad (167)$$

We choose a quantum state $|\psi\rangle$ such the expectation value of $\hat{q}_{\mathbf{k}}(t)$ vanishes for all non-zero \mathbf{k} so that the expectation value of $\hat{\phi}(t, \mathbf{x})$ gives the homogeneous mode that drives the inflation. The quantum fluctuation around this homogeneous part in a quantum state $|\psi\rangle$ is given by

$$\sigma_{\mathbf{k}}^2(t) = \langle \psi | \hat{q}_{\mathbf{k}}^2(t) | \psi \rangle - \langle \psi | \hat{q}_{\mathbf{k}}(t) | \psi \rangle^2 = \langle \psi | \hat{q}_{\mathbf{k}}^2(t) | \psi \rangle \quad (168)$$

It is easy to verify that this fluctuation is just the Fourier transform of the two-point function in this state:

$$\sigma_{\mathbf{k}}^2(t) = \int d^3\mathbf{x} \langle \psi | \hat{\phi}(t, \mathbf{x} + \mathbf{y}) \hat{\phi}(t, \mathbf{y}) | \psi \rangle e^{i\mathbf{k}\cdot\mathbf{x}}. \quad (169)$$

Since $\sigma_{\mathbf{k}}$ characterises the quantum fluctuations, it seems reasonable to introduce a *c-number* field $f(t, \mathbf{x})$ by the definition:

$$f(t, \mathbf{x}) \equiv \int \frac{d^3\mathbf{k}}{(2\pi)^3} \sigma_{\mathbf{k}}(t) e^{i\mathbf{k}\cdot\mathbf{x}} \quad (170)$$

This *c-number* field will have same *c-number power spectrum* as the *quantum* fluctuations. Hence we may take this as our definition of an equivalent classical perturbation. (There are more sophisticated ways of getting this result but none of them are fundamentally more sound than the elementary definition given above. There is a large literature on the subject of quantum to classical transition, especially in the context of gravity; see e.g.[22]) We now have all the ingredients in place. Given the quantum state $|\psi\rangle$, one can explicitly compute $\sigma_{\mathbf{k}}$ and then — using Eq. (166) with $f_k = \sigma_{\mathbf{k}}$ — obtain the density perturbations at the time of re-entry.

The next question we need to address is what is $|\psi\rangle$. The free quantum field theory in the Friedmann background is identical to the quantum mechanics of a bunch of time dependent harmonic oscillators, each labelled by a wave vector \mathbf{k} . The action for a free scalar field in the Friedmann background

$$A = \frac{1}{2} \int d^4x \sqrt{-g} \partial_a \phi \partial^a \phi = \frac{1}{2} \int dt d^3\mathbf{x} a^3 \left[\left(\frac{\partial\phi}{\partial t} \right)^2 - \frac{1}{a^2} (\nabla\phi)^2 \right] \rightarrow \frac{1}{2} \int dt d^3\mathbf{k} a^3 \left[\left(\frac{dq_{\mathbf{k}}}{dt} \right)^2 - \frac{k^2}{a^2} q_{\mathbf{k}}^2 \right] \quad (171)$$

can be thought of as the sum over the actions for an infinite set of harmonic oscillators with mass $m = a^3$ and frequency $\omega_k^2 = k^2/a^2$. (To be precise, one needs to treat the real and imaginary parts of the Fourier transform as independent oscillators and restrict the range of \mathbf{k} ; just pretending that $q_{\mathbf{k}}$ is real amounts the same thing.) The quantum state of the field is just an infinite product of the quantum state $\psi_{\mathbf{k}}[q_{\mathbf{k}}, t]$ for each of the harmonic oscillators and satisfies the Schrodinger equation

$$i \frac{\partial \psi_{\mathbf{k}}}{\partial t} = -\frac{1}{2a^3} \frac{\partial^2 \psi_{\mathbf{k}}}{\partial q_{\mathbf{k}}^2} + \frac{1}{2} a k^2 \psi_{\mathbf{k}} \quad (172)$$

If the quantum state $\psi_{\mathbf{k}}[q_{\mathbf{k}}, t']$ of any given oscillator, labelled by \mathbf{k} , is given at some initial time, t' , we can evolve it to final time:

$$\psi[q_{\mathbf{k}}, t] = \int dq'_{\mathbf{k}} K[q_{\mathbf{k}}, t; q'_{\mathbf{k}}, t'] \psi[q'_{\mathbf{k}}, t'] \quad (173)$$

where K is known in terms of the solutions to the classical equations of motion and $\psi[q'_{\mathbf{k}}, t']$ is the initial state. There is nothing non-trivial in the mathematics, but the physics is completely unknown. The real problem is that unfortunately — in spite of confident assertions in the literature occasionally — we have no clue what $\psi[q'_{\mathbf{k}}, t']$ is. So we need to make more assumptions to proceed further.

One natural choice is the following: It turns out that, Gaussian states of the form

$$\psi_{\mathbf{k}} = A_{\mathbf{k}}(t) \exp[-B_{\mathbf{k}}(t) q_{\mathbf{k}}^2] \quad (174)$$

preserve their form under evolution governed by the Schrodinger equation in Eq. (172). Substituting Eq. (174) in Eq. (172) we can determine the ordinary differential equation which governs $B_{\mathbf{k}}(t)$. (The $A_{\mathbf{k}}(t)$ is trivially fixed by normalization.) Simple algebra shows that $B_{\mathbf{k}}(t)$ can be expressed in the form

$$B_{\mathbf{k}} = -\frac{i}{2} a^3 \left(\frac{1}{f_{\mathbf{k}}} \frac{df_{\mathbf{k}}}{dt} \right) \quad (175)$$

where f_k is the solution to the classical equation of motion:

$$\frac{d^2 f_k}{dt^2} + 3H \frac{df_k}{dt} + \frac{k^2}{a^2} f_k = 0 \quad (176)$$

For the quantum state in Eq. (174), the fluctuations are characterized by

$$\sigma_{\mathbf{k}}^2 = \frac{1}{2} (\text{Re } B_{\mathbf{k}})^{-1} = |f_{\mathbf{k}}|^2 \quad (177)$$

Since one can take different choices for the solutions of Eq. (176) one get different values for $\sigma_{\mathbf{k}}$ and different spectra for perturbations. Any prediction one makes depends on the choice of mode functions. One possibility is to choose the modes so that $\psi_{\mathbf{k}}$ represents the instantaneous vacuum state of the oscillators at some time $t = t_i$. (That is $\text{Re } B_{\mathbf{k}}(t_i) = (1/2)\omega_{\mathbf{k}}^2(t_i)$, say). The final result will then depend on the choice for t_i . One can further make an assumption that we are interested in the limit of $t_i \rightarrow -\infty$; that is the quantum state is an instantaneous ground state in the infinite past. It is easy to show that this corresponds to choosing the following solution to Eq. (176):

$$f_k = \frac{1}{a\sqrt{2k}} (1 + ix) e^{ix}; \quad x = \frac{k}{Ha} \quad (178)$$

which is usually called the Bunch-Davies vacuum. For this choice,

$$|f_k|^2 = \frac{1}{2ka^2} \left(1 + \frac{k^2}{a^2 H^2} \right); \quad |f_k|^2|_{k=aH} \approx \frac{H^2}{k^3} \quad (179)$$

where the second result is at $t = t_{\text{exit}}$ which is what we need to use in Eq. (166), (Numerical factors of order unity cannot be trusted in this computation). We can now determine the amplitude of the perturbation when it re-enters the Hubble radius. Eq. (166) gives:

$$|\Phi_k|_{\text{entry}}^2 = \left(\frac{9H^2}{2V'} \right)^2 |f_k|^2 = \frac{1}{k^3} \frac{9H^6}{4V'^2}; \quad k^3 |\Phi_k|_{\text{entry}}^2 \simeq \left(\frac{H^3}{V'} \right)_{\text{exit}}^2 \quad (180)$$

One sees that the result is scale invariant in the sense that $k^3|\Phi_k|_{\text{entry}}^2$ is independent of k .

It is sometimes claimed in the literature that scale invariant spectrum is a prediction of inflation. *This is simply wrong.* One has to make several *other* assumptions including an all important choice for the quantum state (about which we know nothing) to obtain scale invariant spectrum. In fact, one can prove that, given any power spectrum $\Phi(k)$, one can find a quantum state such that this power spectrum is generated (for an explicit construction, see the last reference in [20]). So whatever results are obtained by observations can be reconciled with inflationary generation of perturbations.

To conclude the discussion, let us work out the perturbations for one specific case. Let us consider the case of the $\lambda\phi^4$ model for which

$$\frac{d^2\phi}{dt^2} + 3H\frac{d\phi}{dt} + V'(\phi) = 0; \quad V(\phi) \approx V_0 - \frac{\lambda}{4}\phi^4 \quad (181)$$

Using

$$N \cong 8\pi G \int_{\phi}^{\phi_f} \frac{V_0}{[-V']} d\phi = \frac{3H^2}{2\lambda} \left(\frac{1}{\phi^2} - \frac{1}{\phi_f^2} \right) \approx \frac{3H^2}{2\lambda\phi^2} \quad (182)$$

we can write

$$-V'(\phi) = \lambda\phi^3 \simeq \frac{H^3}{\lambda^{1/2}N^{3/2}}. \quad (183)$$

so that the result in Eq. (180) becomes:

$$k^{3/2}\Phi_k \simeq H^3 \frac{\lambda^{1/2}N^{3/2}}{H^3} \approx \lambda^{1/2}N^{3/2}. \quad (184)$$

We do get scale invariant spectrum but the amplitude has a serious problem. If we take $N \gtrsim 50$ and note that observations require $k^{3/2}\Phi_k \sim 10^{-4}$ we need to take $\lambda \lesssim 10^{-15}$ for getting consistent values. Such a fine tuning of a dimensionless coupling constant is fairly ridiculous; but over years inflationists have learnt to successfully forget this embarrassment.

Our formalism can also be used to estimate the deviation of the power spectrum from the scale invariant form. To the lowest order we have

$$\Delta_{\Phi}^2 \sim k^3|\Phi_k|^2 \sim \frac{H^6}{(V')^2} \sim \left(\frac{V^3}{m_P^6 V'^2} \right) \quad (185)$$

Let us define the deviation from the scale invariant index by $(n-1) = (d \ln \Delta_{\Phi}^2 / d \ln k)$. Using

$$\frac{d}{d \ln k} = a \frac{d}{da} = \frac{\dot{\phi}}{H} \frac{d}{d\phi} = -\frac{m_P^2}{8\pi} \frac{V'}{V} \frac{d}{d\phi} \quad (186)$$

one finds that

$$1 - n = 6\epsilon - 2\eta \quad (187)$$

Thus, as long as ϵ and η are small we do have $n \approx 1$; what is more, given a potential one can estimate ϵ and η and thus the deviation $(n-1)$.

Finally, note that the same process can also generate spin-2 perturbations. If we take the normalised gravity wave amplitude as $h_{ab} = \sqrt{16\pi G} e_{ab}\phi$, the mode function ϕ behaves like a scalar field. (The normalisation is dictated by the fact that the action for the perturbation should reduce to that of a spin-2 field.) The corresponding power spectrum of gravity waves is

$$P_{\text{grav}}(k) \cong \frac{k^3|h_k|^2}{2\pi^2} = \frac{4}{\pi} \left(\frac{H}{m_P} \right)^2, \quad \Omega_{\text{grav}}(k)h^2 \simeq 10^{-5} \left(\frac{M}{m_P} \right)^3 \quad (188)$$

Comparing the two results

$$\Delta_{\text{scalar}}^2 \sim \frac{H^6}{(V')^2} \sim \left(\frac{V^3}{m_P^6 V'^2} \right); \quad \Delta_{\text{tensor}}^2 \sim \left(\frac{H^2}{m_P^2} \right) \sim \left(\frac{V}{m_P^4} \right) \quad (189)$$

we get $(\Delta_{\text{tensor}}/\Delta_{\text{scalar}})^2 \approx 16\pi\epsilon \ll 1$. Further, if $(1-n) \approx 4\epsilon$ (see Eq. (187) with $\epsilon \sim \eta$) we have the relation $(\Delta_{\text{tensor}}/\Delta_{\text{scalar}})^2 \approx \mathcal{O}(3)(1-n)$ which connects three quantities, all of which are independently observable in principle. If these are actually measured in future it could act as a consistency check of the inflationary paradigm.

VIII. THE DARK ENERGY

It is rather frustrating that the only component of the universe which we understand theoretically is the radiation! While understanding the baryonic and dark matter components [in particular the values of Ω_B and Ω_{DM}] is by no means trivial, the issue of dark energy is lot more perplexing, thereby justifying the attention it has received recently. In this section we will discuss several aspects of the dark energy problem.

The key observational feature of dark energy is that — treated as a fluid with a stress tensor $T_b^a = \text{dia}(\rho, -p, -p, -p)$ — it has an equation state $p = w\rho$ with $w \lesssim -0.8$ at the present epoch. The spatial part \mathbf{g} of the geodesic acceleration (which measures the relative acceleration of two geodesics in the spacetime) satisfies an *exact* equation in general relativity given by:

$$\nabla \cdot \mathbf{g} = -4\pi G(\rho + 3p) \quad (190)$$

This shows that the source of geodesic acceleration is $(\rho + 3p)$ and not ρ . As long as $(\rho + 3p) > 0$, gravity remains attractive while $(\rho + 3p) < 0$ can lead to repulsive gravitational effects. In other words, dark energy with sufficiently negative pressure will accelerate the expansion of the universe, once it starts dominating over the normal matter. This is precisely what is established from the study of high redshift supernova, which can be used to determine the expansion rate of the universe in the past [6].

The simplest model for a fluid with negative pressure is the cosmological constant (for some recent reviews, see [23]) with $w = -1, \rho = -p = \text{constant}$. If the dark energy is indeed a cosmological constant, then it introduces a fundamental length scale in the theory $L_\Lambda \equiv H_\Lambda^{-1}$, related to the constant dark energy density ρ_{DE} by $H_\Lambda^2 \equiv (8\pi G\rho_{\text{DE}}/3)$. In classical general relativity, based on the constants G, c and L_Λ , it is not possible to construct any dimensionless combination from these constants. But when one introduces the Planck constant, \hbar , it is possible to form the dimensionless combination $H_\Lambda^2(G\hbar/c^3) \equiv (L_P^2/L_\Lambda^2)$. Observations then require $(L_P^2/L_\Lambda^2) \lesssim 10^{-123}$. As has been mentioned several times in literature, this will require enormous fine tuning. What is more, in the past, the energy density of normal matter and radiation would have been higher while the energy density contributed by the cosmological constant does not change. Hence we need to adjust the energy densities of normal matter and cosmological constant in the early epoch very carefully so that $\rho_\Lambda \gtrsim \rho_{\text{NR}}$ around the current epoch. This raises the second of the two cosmological constant problems: Why is it that $(\rho_\Lambda/\rho_{\text{NR}}) = \mathcal{O}(1)$ at the *current* phase of the universe ?

Because of these conceptual problems associated with the cosmological constant, people have explored a large variety of alternative possibilities. The most popular among them uses a scalar field ϕ with a suitably chosen potential $V(\phi)$ so as to make the vacuum energy vary with time. The hope then is that, one can find a model in which the current value can be explained naturally without any fine tuning. A simple form of the source with variable w are scalar fields with Lagrangians of different forms, of which we will discuss two possibilities:

$$L_{\text{quin}} = \frac{1}{2}\partial_a\phi\partial^a\phi - V(\phi); \quad L_{\text{tach}} = -V(\phi)[1 - \partial_a\phi\partial^a\phi]^{1/2} \quad (191)$$

Both these Lagrangians involve one arbitrary function $V(\phi)$. The first one, L_{quin} , which is a natural generalization of the Lagrangian for a non-relativistic particle, $L = (1/2)\dot{q}^2 - V(q)$, is usually called quintessence (for a small sample of models, see [24]; there is an extensive and growing literature on scalar field models and more references can be found in the reviews in ref.[23]). When it acts as a source in Friedmann universe, it is characterized by a time dependent $w(t)$ with

$$\rho_q(t) = \frac{1}{2}\dot{\phi}^2 + V; \quad p_q(t) = \frac{1}{2}\dot{\phi}^2 - V; \quad w_q = \frac{1 - (2V/\dot{\phi}^2)}{1 + (2V/\dot{\phi}^2)} \quad (192)$$

The structure of the second Lagrangian (which arise in string theory [25]) in Eq. (191) can be understood by a simple analogy from special relativity. A relativistic particle with (one dimensional) position $q(t)$ and mass m is described by the Lagrangian $L = -m\sqrt{1 - \dot{q}^2}$. It has the energy $E = m/\sqrt{1 - \dot{q}^2}$ and momentum $k = m\dot{q}/\sqrt{1 - \dot{q}^2}$ which are related by $E^2 = k^2 + m^2$. As is well known, this allows the possibility of having *massless* particles with finite energy for which $E^2 = k^2$. This is achieved by taking the limit of $m \rightarrow 0$ and $\dot{q} \rightarrow 1$, while keeping the ratio in $E = m/\sqrt{1 - \dot{q}^2}$ finite. The momentum acquires a life of its own, unconnected with the velocity \dot{q} , and the energy is expressed in terms of the momentum (rather than in terms of \dot{q}) in the Hamiltonian formulation. We can now construct a field theory by upgrading $q(t)$ to a field ϕ . Relativistic invariance now requires ϕ to depend on both space and time [$\phi = \phi(t, \mathbf{x})$] and \dot{q}^2 to be replaced by $\partial_i\phi\partial^i\phi$. It is also possible now to treat the mass parameter m as a function of ϕ , say, $V(\phi)$ thereby obtaining a field theoretic Lagrangian $L = -V(\phi)\sqrt{1 - \partial^i\phi\partial_i\phi}$. The Hamiltonian structure of this theory is algebraically very similar to the special relativistic example we started with. In particular,

the theory allows solutions in which $V \rightarrow 0$, $\partial_i \phi \partial^i \phi \rightarrow 1$ simultaneously, keeping the energy (density) finite. Such solutions will have finite momentum density (analogous to a massless particle with finite momentum k) and energy density. Since the solutions can now depend on both space and time (unlike the special relativistic example in which q depended only on time), the momentum density can be an arbitrary function of the spatial coordinate. The structure of this Lagrangian is similar to those analyzed in a wide class of models called *K-essence* [26] and provides a rich gamut of possibilities in the context of cosmology [27, 28].

Since the quintessence field (or the tachyonic field) has an undetermined free function $V(\phi)$, it is possible to choose this function in order to produce a given $H(a)$. To see this explicitly, let us assume that the universe has two forms of energy density with $\rho(a) = \rho_{\text{known}}(a) + \rho_\phi(a)$ where $\rho_{\text{known}}(a)$ arises from any known forms of source (matter, radiation, ...) and $\rho_\phi(a)$ is due to a scalar field. Let us first consider quintessence. Here, the potential is given implicitly by the form [27, 29].

$$V(a) = \frac{1}{16\pi G} H(1-Q) \left[6H + 2aH' - \frac{aHQ'}{1-Q} \right]; \quad \phi(a) = \left[\frac{1}{8\pi G} \right]^{1/2} \int \frac{da}{a} \left[aQ' - (1-Q) \frac{d \ln H^2}{d \ln a} \right]^{1/2} \quad (193)$$

where $Q(a) \equiv [8\pi G \rho_{\text{known}}(a)/3H^2(a)]$ and prime denotes differentiation with respect to a . (The result used in Eq. (155) is just a special case of this when $Q = 0$) Given any $H(a)$, $Q(a)$, these equations determine $V(a)$ and $\phi(a)$ and thus the potential $V(\phi)$. *Every quintessence model studied in the literature can be obtained from these equations.*

Similar results exists for the tachyonic scalar field as well [27]. For example, given any $H(a)$, one can construct a tachyonic potential $V(\phi)$ so that the scalar field is the source for the cosmology. The equations determining $V(\phi)$ are now given by:

$$\phi(a) = \int \frac{da}{aH} \left(\frac{aQ'}{3(1-Q)} - \frac{2aH'}{3H} \right)^{1/2}; \quad V = \frac{3H^2}{8\pi G} (1-Q) \left(1 + \frac{2aH'}{3H} - \frac{aQ'}{3(1-Q)} \right)^{1/2} \quad (194)$$

Equations (194) completely solve the problem. Given any $H(a)$, these equations determine $V(a)$ and $\phi(a)$ and thus the potential $V(\phi)$. A wide variety of phenomenological models with time dependent cosmological constant have been considered in the literature all of which can be mapped to a scalar field model with a suitable $V(\phi)$.

While the scalar field models enjoy considerable popularity (one reason being they are easy to construct!) it is very doubtful whether they have helped us to understand the nature of the dark energy at any deeper level. These models, viewed objectively, suffer from several shortcomings:

- They completely lack predictive power. As explicitly demonstrated above, virtually every form of $a(t)$ can be modeled by a suitable “designer” $V(\phi)$.
- These models are degenerate in another sense. The previous discussion illustrates that even when $w(a)$ is known/specified, it is not possible to proceed further and determine the nature of the scalar field Lagrangian. The explicit examples given above show that there are *at least* two different forms of scalar field Lagrangians (corresponding to the quintessence or the tachyonic field) which could lead to the same $w(a)$. (See the second paper in ref.[7] for an explicit example of such a construction.)
- All the scalar field potentials require fine tuning of the parameters in order to be viable. This is obvious in the quintessence models in which adding a constant to the potential is the same as invoking a cosmological constant. So to make the quintessence models work, *we first need to assume the cosmological constant is zero*. These models, therefore, merely push the cosmological constant problem to another level, making it somebody else’s problem!.
- By and large, the potentials used in the literature have no natural field theoretical justification. All of them are non-renormalisable in the conventional sense and have to be interpreted as a low energy effective potential in an ad hoc manner.
- One key difference between cosmological constant and scalar field models is that the latter lead to a $w(a)$ which varies with time. If observations have demanded this, or even if observations have ruled out $w = -1$ at the present epoch, then one would have been forced to take alternative models seriously. However, all available observations are consistent with cosmological constant ($w = -1$) and — in fact — the possible variation of w is strongly constrained [30] as shown in Figure 6.
- While on the topic of observational constraints on $w(t)$, it must be stressed that: (a) There is fair amount of tension between WMAP and SN data and one should be very careful about the priors used in these analysis. (b) There is no observational evidence for $w < -1$. (c) It is likely that more homogeneous, future, data sets of SN might show better agreement with WMAP results. (For more details related to these issues, see the last reference in [30].)

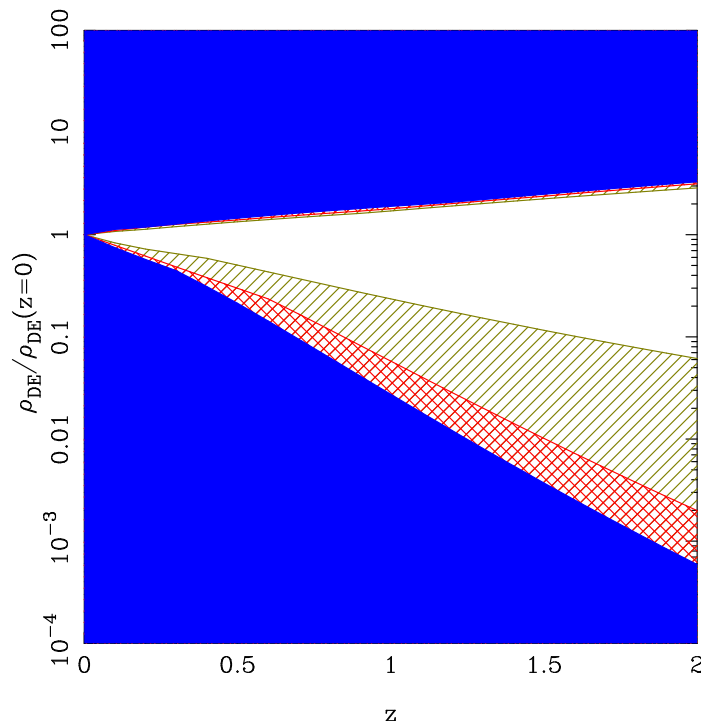


FIG. 6: The observational constraints on the variation of dark energy density as a function of redshift from WMAP and SNLS data (see [30]). The green/hatched region is excluded at 68% confidence limit, red/cross-hatched region at 95% confidence level and the blue/solid region at 99% confidence limit. The white region shows the allowed range of variation of dark energy at 68% confidence limit.

The observational and theoretical features described above suggests that one should consider cosmological constant as the most natural candidate for dark energy. Though it leads to well know fine tuning problems, it also has certain attractive features that need to kept in mind.

- Cosmological constant is the most economical [just one number] and simplest explanation for all the observations. We stress that there is absolutely *no* evidence for variation of dark energy density with redshift, which is consistent with the assumption of cosmological constant .
- Once we invoke the cosmological constant classical gravity will be described by the three constants G, c and $\Lambda \equiv L_\Lambda^{-2}$. It is *not* possible to obtain a dimensionless quantity from these; so, within classical theory, there is no fine tuning issue. Since $\Lambda(G\hbar/c^3) \equiv (L_P/L_\Lambda)^2 \approx 10^{-123}$, it is obvious that the cosmological constant is telling us something regarding *quantum gravity*, indicated by the combination $G\hbar$. *An acid test for any quantum gravity model will be its ability to explain this value*; needless to say, all the currently available models — strings, loops etc. — flunk this test.
- So, if dark energy is indeed cosmological constant this will be the greatest contribution from cosmology to fundamental physics. It will be unfortunate if we miss this chance by invoking some scalar field epicycles!

In this context, it is worth stressing another peculiar feature of cosmological constant when it is treated as a clue to quantum gravity. It is well known that, based on energy scales, the cosmological constant problem is an infra red problem *par excellence*. At the same time, it is a relic of a quantum gravitational effect or principle of unknown nature. An analogy [31] will be helpful to illustrate this point. Suppose one solves the Schrodinger equation for the Helium atom for the quantum states of the two electrons $\psi(x_1, x_2)$. When the result is compared with observations, one will find that only half the states — those in which $\psi(x_1, x_2)$ is antisymmetric under $x_1 \longleftrightarrow x_2$ interchange — are realized in nature. But the low energy Hamiltonian for electrons in the Helium atom has no information about this effect! Here is low energy (IR) effect which is a relic of relativistic quantum field theory (spin-statistics theorem) that is totally non perturbative, in the sense that writing corrections to the Helium atom Hamiltonian in some $(1/c)$ expansion will *not* reproduce this result. The current value of cosmological constant could very well be related to quantum gravity in a similar way. There must exist a deep principle in quantum gravity which leaves its non perturbative trace even in the low energy limit that appears as the cosmological constant .

Let us now turn our attention to few of the many attempts to understand the cosmological constant. The choice is, of course, dictated by personal bias and is definitely a non-representative sample. A host of other approaches exist in literature, some of which can be found in [32].

A. Gravitational Holography

One possible way of addressing this issue is to simply eliminate from the gravitational theory those modes which couple to cosmological constant. If, for example, we have a theory in which the source of gravity is $(\rho + p)$ rather than $(\rho + 3p)$ in Eq. (190), then cosmological constant will not couple to gravity at all. (The non linear coupling of matter with gravity has several subtleties; see eg. [33].) Unfortunately it is not possible to develop a covariant theory of gravity using $(\rho + p)$ as the source. But we can probably gain some insight from the following considerations. Any metric g_{ab} can be expressed in the form $g_{ab} = f^2(x)q_{ab}$ such that $\det q = 1$ so that $\det g = f^4$. From the action functional for gravity

$$A = \frac{1}{16\pi G} \int d^4x (R - 2\Lambda) \sqrt{-g} = \frac{1}{16\pi G} \int d^4x R \sqrt{-g} - \frac{\Lambda}{8\pi G} \int d^4x f^4(x) \quad (195)$$

it is obvious that the cosmological constant couples *only* to the conformal factor f . So if we consider a theory of gravity in which $f^4 = \sqrt{-g}$ is kept constant and only q_{ab} is varied, then such a model will be oblivious of direct coupling to cosmological constant. If the action (without the Λ term) is varied, keeping $\det g = -1$, say, then one is lead to a *unimodular theory of gravity* that has the equations of motion $R_{ab} - (1/4)g_{ab}R = \kappa(T_{ab} - (1/4)g_{ab}T)$ with zero trace on both sides. Using the Bianchi identity, it is now easy to show that this is equivalent to the usual theory with an *arbitrary* cosmological constant. That is, cosmological constant arises as an undetermined integration constant in this model [34].

The same result arises in another, completely different approach to gravity. In the standard approach to gravity one uses the Einstein-Hilbert Lagrangian $L_{EH} \propto R$ which has a formal structure $L_{EH} \sim R \sim (\partial g)^2 + \partial^2 g$. If the surface term obtained by integrating $L_{sur} \propto \partial^2 g$ is ignored (or, more formally, canceled by an extrinsic curvature term) then the Einstein's equations arise from the variation of the bulk term $L_{bulk} \propto (\partial g)^2$ which is the non-covariant Γ^2 Lagrangian. There is, however, a remarkable relation between L_{bulk} and L_{sur} :

$$\sqrt{-g}L_{sur} = -\partial_a \left(g_{ij} \frac{\partial \sqrt{-g} L_{bulk}}{\partial (\partial_a g_{ij})} \right) \quad (196)$$

which allows a dual description of gravity using either L_{bulk} or L_{sur} ! It is possible to obtain the dynamics of gravity [35] from an approach which uses *only* the surface term of the Hilbert action; *we do not need the bulk term at all!* This suggests that *the true degrees of freedom of gravity for a volume \mathcal{V} reside in its boundary $\partial\mathcal{V}$* — a point of view that is strongly supported by the study of horizon entropy, which shows that the degrees of freedom hidden by a horizon scales as the area and not as the volume. The resulting equations can be cast in a thermodynamic form $TdS = dE + PdV$ and the continuum spacetime is like an elastic solid (see e.g. [36]) with Einstein's equations providing the macroscopic description. Interestingly, the *cosmological constant arises again in this approach as a undetermined integration constant* but closely related to the ‘bulk expansion’ of the solid.

While this is all very interesting, we still need an extra physical principle to fix the value (even the sign) of cosmological constant. One possible way of doing this is to interpret the Λ term in the action as a Lagrange multiplier for the proper volume of the spacetime. Then it is reasonable to choose the cosmological constant such that the total proper volume of the universe is equal to a specified number. While this will lead to a cosmological constant which has the correct order of magnitude, it has several obvious problems. First, the proper four volume of the universe is infinite unless we make the spatial sections compact and restrict the range of time integration. Second, this will lead to a dark energy density which varies as t^{-2} (corresponding to $w = -1/3$) which is ruled out by observations.

B. Cosmic Lenz law

Another possibility which has been attempted in the literature tries to “cancel out” the cosmological constant by some process, usually quantum mechanical in origin. One of the simplest ideas will be to ask whether switching on a cosmological constant will lead to a vacuum polarization with an effective energy momentum tensor that will tend to cancel out the cosmological constant. A less subtle way of doing this is to invoke another scalar field (here we go again!) such that it can couple to cosmological constant and reduce its effective value [37]. Unfortunately, none of this could be made to work properly. By and large, these approaches lead to an energy density which is either

$\rho_{UV} \propto L_P^{-4}$ (where L_P is the Planck length) or to $\rho_{IR} \propto L_\Lambda^{-4}$ (where $L_\Lambda = H_\Lambda^{-1}$ is the Hubble radius associated with the cosmological constant). The first one is too large while the second one is too small!

C. Geometrical Duality in our Universe

While the above ideas do not work, it gives us a clue. A universe with two length scales L_Λ and L_P will be asymptotically De Sitter with $a(t) \propto \exp(t/L_\Lambda)$ at late times. There are some curious features in such a universe which we will now describe. Given the two length scales L_P and L_Λ , one can construct two energy scales $\rho_{UV} = 1/L_P^4$ and $\rho_{IR} = 1/L_\Lambda^4$ in natural units ($c = \hbar = 1$). There is sufficient amount of justification from different theoretical perspectives to treat L_P as the zero point length of spacetime [38], giving a natural interpretation to ρ_{UV} . The second one, ρ_{IR} also has a natural interpretation. The universe which is asymptotically De Sitter has a horizon and associated thermodynamics [39] with a temperature $T = H_\Lambda/2\pi$ and the corresponding thermal energy density $\rho_{thermal} \propto T^4 \propto 1/L_\Lambda^4 = \rho_{IR}$. Thus L_P determines the *highest* possible energy density in the universe while L_Λ determines the *lowest* possible energy density in this universe. As the energy density of normal matter drops below this value, the thermal ambience of the De Sitter phase will remain constant and provide the irreducible ‘vacuum noise’. *Note that the dark energy density is the the geometric mean $\rho_{DE} = \sqrt{\rho_{IR}\rho_{UV}}$ between the two energy densities.* If we define a dark energy length scale L_{DE} such that $\rho_{DE} = 1/L_{DE}^4$ then $L_{DE} = \sqrt{L_P L_\Lambda}$ is the geometric mean of the two length scales in the universe. (Incidentally, $L_{DE} \approx 0.04$ mm is macroscopic; it is also pretty close to the length scale associated with a neutrino mass of 10^{-2} eV; another intriguing coincidence ?!)

Using the characteristic length scale of expansion, the Hubble radius $d_H \equiv (\dot{a}/a)^{-1}$, we can distinguish between three different phases of such a universe. The first phase is when the universe went through a inflationary expansion with $d_H = \text{constant}$; the second phase is the radiation/matter dominated phase in which most of the standard cosmology operates and d_H increases monotonically; the third phase is that of re-inflation (or accelerated expansion) governed by the cosmological constant in which d_H is again a constant. The first and last phases are time translation invariant; that is, $t \rightarrow t + \text{constant}$ is an (approximate) invariance for the universe in these two phases. The universe satisfies the perfect cosmological principle and is in steady state during these phases!

In fact, one can easily imagine a scenario in which the two De Sitter phases (first and last) are of arbitrarily long duration [40]. If $\Omega_\Lambda \approx 0.7$, $\Omega_{DM} \approx 0.3$ the final De Sitter phase *does* last forever; as regards the inflationary phase, nothing prevents it from lasting for arbitrarily long duration. Viewed from this perspective, the in between phase — in which most of the ‘interesting’ cosmological phenomena occur — is of negligible measure in the span of time. It merely connects two steady state phases of the universe.

While the two De Sitter phases can last forever in principle, there is a natural cut off length scale in both of them which makes the region of physical relevance to be finite [40]. Let us first discuss the case of re-inflation in the late universe. As the universe grows exponentially in the phase 3, the wavelength of CMBR photons are being redshifted rapidly. When the temperature of the CMBR radiation drops below the De Sitter temperature (which happens when the wavelength of the typical CMBR photon is stretched to the L_Λ .) the universe will be essentially dominated by the vacuum thermal noise [39] due to the horizon in the De Sitter phase. This happens when the expansion factor is $a = a_F$ determined by the equation $T_0(a_0/a_F) = (1/2\pi L_\Lambda)$. Let $a = a_\Lambda$ be the epoch at which cosmological constant started dominating over matter, so that $(a_\Lambda/a_0)^3 = (\Omega_{DM}/\Omega_\Lambda)$. Then we find that the dynamic range of the phase 3 is

$$\frac{a_F}{a_\Lambda} = 2\pi T_0 L_\Lambda \left(\frac{\Omega_\Lambda}{\Omega_{DM}} \right)^{1/3} \approx 3 \times 10^{30} \quad (197)$$

Interestingly enough, one can also impose a similar bound on the physically relevant duration of inflation. We know that the quantum fluctuations, generated during this inflationary phase, could act as seeds of structure formation in the universe. Consider a perturbation at some given wavelength scale which is stretched with the expansion of the universe as $\lambda \propto a(t)$. During the inflationary phase, the Hubble radius remains constant while the wavelength increases, so that the perturbation will ‘exit’ the Hubble radius at some time. In the radiation dominated phase, the Hubble radius $d_H \propto t \propto a^2$ grows faster than the wavelength $\lambda \propto a(t)$. Hence, normally, the perturbation will ‘re-enter’ the Hubble radius at some time. If there was no re-inflation, *all* wavelengths will re-enter the Hubble radius sooner or later. But if the universe undergoes re-inflation, then the Hubble radius ‘flattens out’ at late times and some of the perturbations will *never* reenter the Hubble radius! If we use the criterion that we need the perturbation to reenter the Hubble radius, we get a natural bound on the duration of inflation which is of direct astrophysical relevance. Consider a perturbation which leaves the Hubble radius (H_{in}^{-1}) during the inflationary epoch at $a = a_i$. It will grow to the size $H_{in}^{-1}(a/a_i)$ at a later epoch. We want to determine a_i such that this length scale grows to L_Λ just when the dark energy starts dominating over matter; that is at the epoch $a = a_\Lambda = a_0(\Omega_{DM}/\Omega_\Lambda)^{1/3}$. This gives

$H_{in}^{-1}(a_\Lambda/a_i) = L_\Lambda$ so that $a_i = (H_{in}^{-1}/L_\Lambda)(\Omega_{DM}/\Omega_\Lambda)^{1/3}a_0$. On the other hand, the inflation ends at $a = a_{end}$ where $a_{end}/a_0 = T_0/T_{reheat}$ where T_{reheat} is the temperature to which the universe has been reheated at the end of inflation. Using these two results we can determine the dynamic range of this phase 1 to be

$$\frac{a_{end}}{a_i} = \left(\frac{T_0 L_\Lambda}{T_{reheat} H_{in}^{-1}} \right) \left(\frac{\Omega_\Lambda}{\Omega_{DM}} \right)^{1/3} = \frac{(a_F/a_\Lambda)}{2\pi T_{reheat} H_{in}^{-1}} \cong 10^{25} \quad (198)$$

where we have used the fact that, for a GUTs scale inflation with $E_{GUT} = 10^{14} GeV$, $T_{reheat} = E_{GUT}$, $\rho_{in} = E_{GUT}^4$ we have $2\pi H_{in}^{-1} T_{reheat} = (3\pi/2)^{1/2}(E_P/E_{GUT}) \approx 10^5$. If we consider a quantum gravitational, Planck scale, inflation with $2\pi H_{in}^{-1} T_{reheat} = \mathcal{O}(1)$, the ranges in Eq. (197) and Eq. (198) are approximately equal.

This fact is definitely telling us something regarding the duality between Planck scale and Hubble scale or between the infrared and ultraviolet limits of the theory. The mystery is compounded by the fact the asymptotic De Sitter phase has an observer dependent horizon and related thermal properties [39]. Recently, it has been shown — in a series of papers, see ref.[35] — that it is possible to obtain classical relativity from purely thermodynamic considerations. It is difficult to imagine that these features are unconnected and accidental; at the same time, it is difficult to prove a definite connection between these ideas and the cosmological constant.

D. Gravity as detector of the vacuum energy

Finally, we will describe an idea which *does* lead to the correct value of cosmological constant. The conventional discussion of the relation between cosmological constant and vacuum energy density is based on evaluating the zero point energy of quantum fields with an ultraviolet cutoff and using the result as a source of gravity. Any reasonable cutoff will lead to a vacuum energy density ρ_{vac} which is unacceptably high. This argument, however, is too simplistic since the zero point energy — obtained by summing over the $(1/2)\hbar\omega_k$ — has no observable consequence in any other phenomena and can be subtracted out by redefining the Hamiltonian. The observed non trivial features of the vacuum state of QED, for example, arise from the *fluctuations* (or modifications) of this vacuum energy rather than the vacuum energy itself. This was, in fact, known fairly early in the history of cosmological constant problem and is stressed by Zeldovich [41] who explicitly calculated one possible contribution to *fluctuations* after subtracting away the mean value. This suggests that we should consider the fluctuations in the vacuum energy density in addressing the cosmological constant problem.

If the vacuum probed by the gravity can readjust to take away the bulk energy density $\rho_{UV} \simeq L_P^{-4}$, quantum *fluctuations* can generate the observed value ρ_{DE} . One of the simplest models [42] which achieves this uses the fact that, in the semi-classical limit, the wave function describing the universe of proper four-volume \mathcal{V} will vary as $\Psi \propto \exp(-iA_0) \propto \exp[-i(\Lambda_{eff}\mathcal{V}/L_P^2)]$. If we treat $(\Lambda/L_P^2, \mathcal{V})$ as conjugate variables then uncertainty principle suggests $\Delta\Lambda \approx L_P^2/\Delta\mathcal{V}$. If the four volume is built out of Planck scale substructures, giving $\mathcal{V} = NL_P^4$, then the Poisson fluctuations will lead to $\Delta\mathcal{V} \approx \sqrt{\mathcal{V}}L_P^2$ giving $\Delta\Lambda = L_P^2/\Delta\mathcal{V} \approx 1/\sqrt{\mathcal{V}} \approx H_0^2$. (This idea can be a more quantitative; see [42]).

Similar viewpoint arises, more rigorously, when we study the question of *detecting* the energy density using gravitational field as a probe. Recall that an Unruh-DeWitt detector with a local coupling $L_I = M(\tau)\phi[x(\tau)]$ to the *field* ϕ actually responds to $\langle 0|\phi(x)\phi(y)|0\rangle$ rather than to the field itself [43]. Similarly, one can use the gravitational field as a natural “detector” of energy momentum tensor T_{ab} with the standard coupling $L = \kappa h_{ab}T^{ab}$. Such a model was analysed in detail in ref. [44] and it was shown that the gravitational field responds to the two point function $\langle 0|T_{ab}(x)T_{cd}(y)|0\rangle$. In fact, it is essentially this fluctuations in the energy density which is computed in the inflationary models (see Eq. (170)) as the seed *source* for gravitational field, as stressed in ref. [11]. All these suggest treating the energy fluctuations as the physical quantity “detected” by gravity, when one needs to incorporate quantum effects. If the cosmological constant arises due to the energy density of the vacuum, then one needs to understand the structure of the quantum vacuum at cosmological scales. Quantum theory, especially the paradigm of renormalization group has taught us that the energy density — and even the concept of the vacuum state — depends on the scale at which it is probed. The vacuum state which we use to study the lattice vibrations in a solid, say, is not the same as vacuum state of the QED.

In fact, it seems *inevitable* that in a universe with two length scale L_Λ, L_P , the vacuum fluctuations will contribute an energy density of the correct order of magnitude $\rho_{DE} = \sqrt{\rho_{IR}\rho_{UV}}$. The hierarchy of energy scales in such a universe, as detected by the gravitational field has [40, 45] the pattern

$$\rho_{vac} = \frac{1}{L_P^4} + \frac{1}{L_P^4} \left(\frac{L_P}{L_\Lambda} \right)^2 + \frac{1}{L_P^4} \left(\frac{L_P}{L_\Lambda} \right)^4 + \dots \quad (199)$$

The first term is the bulk energy density which needs to be renormalized away (by a process which we do not understand at present); the third term is just the thermal energy density of the De Sitter vacuum state; what is interesting is that quantum fluctuations in the matter fields *inevitably generate* the second term.

The key new ingredient arises from the fact that the properties of the vacuum state depends on the scale at which it is probed and it is not appropriate to ask questions without specifying this scale. If the spacetime has a cosmological horizon which blocks information, the natural scale is provided by the size of the horizon, L_Λ , and we should use observables defined within the accessible region. The operator $H(< L_\Lambda)$, corresponding to the total energy inside a region bounded by a cosmological horizon, will exhibit fluctuations ΔE since vacuum state is not an eigenstate of *this* operator. The corresponding fluctuations in the energy density, $\Delta\rho \propto (\Delta E)/L_\Lambda^3 = f(L_P, L_\Lambda)$ will now depend on both the ultraviolet cutoff L_P as well as L_Λ . To obtain $\Delta\rho_{\text{vac}} \propto \Delta E/L_\Lambda^3$ which scales as $(L_P L_\Lambda)^{-2}$ we need to have $(\Delta E)^2 \propto L_P^{-4} L_\Lambda^2$; that is, the square of the energy fluctuations should scale as the surface area of the bounding surface which is provided by the cosmic horizon. Remarkably enough, a rigorous calculation [45] of the dispersion in the energy shows that for $L_\Lambda \gg L_P$, the final result indeed has the scaling

$$(\Delta E)^2 = c_1 \frac{L_\Lambda^2}{L_P^4} \quad (200)$$

where the constant c_1 depends on the manner in which ultra violet cutoff is imposed. Similar calculations have been done (with a completely different motivation, in the context of entanglement entropy) by several people and it is known that the area scaling found in Eq. (200), proportional to L_Λ^2 , is a generic feature [46]. For a simple exponential UV-cutoff, $c_1 = (1/30\pi^2)$ but cannot be computed reliably without knowing the full theory. We thus find that the fluctuations in the energy density of the vacuum in a sphere of radius L_Λ is given by

$$\Delta\rho_{\text{vac}} = \frac{\Delta E}{L_\Lambda^3} \propto L_P^{-2} L_\Lambda^{-2} \propto \frac{H_\Lambda^2}{G} \quad (201)$$

The numerical coefficient will depend on c_1 as well as the precise nature of infrared cutoff radius (like whether it is L_Λ or $L_\Lambda/2\pi$ etc.). It would be pretentious to cook up the factors to obtain the observed value for dark energy density. But it is a fact of life that a fluctuation of magnitude $\Delta\rho_{\text{vac}} \simeq H_\Lambda^2/G$ will exist in the energy density inside a sphere of radius H_Λ^{-1} if Planck length is the UV cut off. *One cannot get away from it.* On the other hand, observations suggest that there is a ρ_{vac} of similar magnitude in the universe. It seems natural to identify the two, after subtracting out the mean value by hand. Our approach explains why there is a *surviving* cosmological constant which satisfies $\rho_{\text{DE}} = \sqrt{\rho_{\text{IR}}\rho_{\text{UV}}}$ which — in our opinion — is *the* problem.

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