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The cosmological constant and dark energy

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Physics welcomes the idea that space contains energy whose gravitational effect approximates that of Einstein's cosmological constant, Λ; today the concept is termed dark energy or quintessence. Physics also suggests that dark energy could be dynamical, allowing for the arguably appealing picture of an evolving dark-energy density approaching its natural value, zero, and small now because the expanding universe is old. This would alleviate the classical problem of the curious energy scale of a millielectron volt associated with a constant Λ. Dark energy may have been detected by recent cosmological tests. These tests make a good scientific case for the context, in the relativistic Friedmann-Lemaître model, in which the gravitational inverse-square law is applied to the scales of cosmology. We have well-checked evidence that the mean mass density is not much more than one-quarter of the critical Einstein–de Sitter value. The case for detection of dark energy is not yet as convincing but still serious; we await more data, which may be derived from work in progress. Planned observations may detect the evolution of the dark-energy density; a positive result would be a considerable stimulus for attempts at understanding the microphysics of dark energy. This review presents the basic physics and astronomy of the subject, reviews the history of ideas, assesses the state of the observational evidence, and comments on recent developments in the search for a fundamental theory.

CONTENTS

I. INTRODUCTION

There is significant observational evidence for the detection of Einstein's cosmological constant, Λ, or a component of the material content of the universe that varies only slowly with time and space and so acts like Λ. We shall use the term *dark energy* for Λ or a component that acts like it. Detection of dark energy would be a new clue to an old puzzle: the gravitational effect of the zero-point energies of particles and fields. The total with other energies, that are close to homogeneous and nearly independent of time, acts as dark energy. What is puzzling is that the value of the dark-energy density has to be tiny compared to what is suggested by dimensional analysis; the startling new evidence is that it may be different from the only other natural value, zero.

The main question to consider now is whether to accept the evidence for detection of dark energy. We outline the nature of the case in this section. After reviewing the basic concepts of the relativistic world model in Sec. II, and in Sec. III reviewing the history of ideas, we present in Sec. IV a more detailed assessment of the cosmological tests and the evidence for detection of Λ or its analog in dark energy.

There is little new to report on the big issue for physics—why the dark-energy density is so small—since Weinberg's (1989) review in this journal.[1] But there have been analyses of a simpler idea: can we imagine that the present dark-energy density is evolving, perhaps approaching zero? Models are introduced in Secs. II.C and III.E, and recent work is summarized in more detail in the Appendix. Feasible advances in cosmological tests could detect evolution of the dark-energy density, and perhaps its gravitational response to large-scale fluctuations in the mass distribution. This would substantially motivate the search for a more fundamental physics model for dark energy.

A. The issues for observational cosmology

We will make two points. First, cosmology has a substantial observational and experimental basis, which supports many aspects of the standard model as almost certainly being good approximations to reality. Second, the empirical basis is not nearly as strong for cosmology as it is for the standard model of particle physics: in cosmology it is not yet a matter of measuring the parameters in a well-established theory.

To explain the second point we direct our attention to those more accustomed to experiments in the laboratory than to astronomy-related observations of astronomers' Tantalus principle: one can look at distant objects but never touch them. For example, the observations of supernovae in distant galaxies offer evidence of dark energy, under the assumption that distant and nearby supernovae are drawn from the same statistical sample (that is, that they are statistically similar enough for the purpose of this test). There is no direct way to check this, and it is easy to imagine differences between distant and nearby supernovae of the same nominal type. More distant supernovae are seen in younger galaxies, because of the travel time of light, and these younger galaxies tend to have more massive rapidly evolving stars with lower heavy-element abundances. How do we know that the properties of the supernovae are not also different? The reader is referred to Leibundgut's (2001, Sec. 4) discussion of astrophysical hazards. Astronomers have checks for this and other issues of interpretation when considering the observations used in cosmological tests. But it takes nothing away from this careful and elegant work to note that the checks are seldom convincing, because the astronomy is complicated and what can be observed is sparse. What is more, we do not know ahead of time that the physics well tested on scales ranging from the laboratory to the Solar System survives the enormous extrapolation to cosmology.

The situation is by no means hopeless. We now have significant cross-checks from the consistency of results based on independent applications of the astronomy and of the physics of the cosmological model. If the physics or astronomy was faulty we would not expect consistency from independent lines of evidence—apart from the occasional accident and the occasional tendency to stop the analysis when it approaches the "right answer." We have to demand abundant evidence of consistency, and that is starting to appear.

The case for detection of Λ or dark energy commences with the Friedmann-Lemaître cosmological model. In this model the expansion history of the universe is determined by a set of dimensionless parameters whose sum is normalized to unity,

$$\Omega_{M0}+\Omega_{R0}+\Omega_{\Lambda 0}+\Omega_{K0}=1. \tag{1}$$

The first, Ω_{M0}, is a measure of the present mean mass density in nonrelativistic matter, mainly baryons and nonbaryonic dark matter. The second, $\Omega_{R0}\sim 1\times 10^{-4}$, is a measure of the present mass in the relativistic 3-K thermal cosmic microwave background radiation, which almost homogeneously fills space, and the accompanying low-mass neutrinos. The third is a measure of Λ or the present value of the dark-energy equivalent. The fourth, Ω_{K0}, is an effect of the curvature of space. We review some details of these parameters in the next section, and of their measurements in Sec. IV.

The most direct evidence for detection of dark energy comes from observations of supernovae of a type whose intrinsic luminosities are close to uniform (after subtle astronomical corrections, a few details of which are discussed in Sec. IV.B.4). The observed brightness as a function of the wavelength shift of the radiation probes the geometry of spacetime, in what has come to be called the redshift-magnitude relation.[2] The measurements agree with the relativistic cosmological model with $\Omega_{K0}=0$, meaning no space curvature, and $\Omega_{\Lambda 0}\sim 0.7$, meaning nonzero Λ. A model with $\Omega_{\Lambda 0}=0$ is two

[1] Sahni and Starobinsky (2000); Carroll (2001); Weinberg (2001); Witten (2001); and Ellwanger (2002) present more recent reviews.

[2] The apparent magnitude is $m=-2.5\log_{10}f$ plus a constant, where f is the detected energy flux density in a chosen wavelength band. The standard measure of the wavelength shift, due to the expansion of the universe, is the redshift z defined in Eq. (7) below.

or three standard deviations off the best fit, depending on the data set and analysis technique. This is an important indication, but 2 to 3 σ is not convincing, even when we can be sure that systematic errors are under reasonable control. And we have to consider that there may be a significant systematic error from differences between distant, high-redshift, and nearby, low-redshift, supernovae.

There is a check, based on the cold-dark-matter (CDM) model[3] for structure formation. The fit of the model to the observations reviewed in Sec. IV.B yields two key constraints. First, the angular power spectrum of fluctuations in the temperature of the 3-K thermal cosmic microwave background radiation across the sky indicates that Ω_{K0} is small. Second, the power spectrum of the spatial distribution of the galaxies requires $\Omega_{M0} \sim 0.25$. Similar estimates of Ω_{M0} follow from independent lines of observational evidence. The rate of gravitational lensing prefers a somewhat larger value (if Ω_{K0} is small), and some dynamical analyses of systems of galaxies prefer lower Ω_{M0}. But the differences could all result from measurement uncertainties. Since Ω_{R0} in Eq. (1) is small, the conclusion is that $\Omega_{\Lambda 0}$ is large, in excellent agreement with the supernovae result.

Caution is in order, however, because this check depends on the CDM model for structure formation. We cannot see the dark matter, so we naturally assign it the simplest properties possible. Maybe it is significant that the model has observational problems with galaxy formation, as discussed in Sec. IV.A.2, or maybe these problems are only apparent, due to the complications of the astronomy. We are going to have to determine which is correct before we can have confidence in the role of the CDM model in cosmological tests. We will get a strong hint from current precision angular distribution measurements of the 3-K thermal cosmic microwave background radiation.[4] If the results match precisely the prediction of the relativistic model for cosmology and the CDM model for structure formation, with parameter choices that agree with the constraints from all the other cosmological tests, there will be strong evidence that we are approaching a good approximation to reality, and the completion of the great program of cosmological tests that commenced in the 1930s. But all that is in the future.

We wish to emphasize that the advances in the empirical basis for cosmology already are very real and substantial. How firm the conclusion is depends on the issue, of course. Every competent cosmologist we know accepts as established beyond reasonable doubt that the universe is expanding and cooling in a near homogeneous and isotropic way from a hotter denser state: how else could space, which is transparent now, have been filled with radiation that has relaxed to a thermal spectrum? The debate is when the expansion commenced or became a meaningful concept. Some whose opinions and research we respect question the extrapolation of the gravitational inverse-square law, in its use in estimates of masses in galaxies and systems of galaxies, and of Ω_{M0}. We agree that this law is one of the hypotheses to be tested. Our conclusion from the cosmological tests of Sec. IV is that the law passes significant, though not yet complete, tests, and that we already have a strong scientific case, resting on the abundance of cross-checks, that the matter density parameter Ω_{M0} is about one-quarter. The case for detection of $\Omega_{\Lambda 0}$ is significant too, but not yet as compelling.

For the most part the results of the cosmological tests agree wonderfully well with accepted theory. But the observational challenges to the tests are substantial: we are drawing profound conclusions from very limited information. We have to be liberal when considering ideas about what the universe is like, and conservative when accepting ideas into the established canon.

B. The opportunity for physics

Unless there is some serious and quite unexpected flaw in our understanding of the principles of physics we can be sure the zero-point energy of the electromagnetic field at laboratory wavelengths is real and measurable, as in the Casimir (1948) effect.[5] Like all energy, this zero-point energy has to contribute to the source term in Einstein's gravitational field equation. If, as seems likely, the zero-point energy of the electromagnetic field is close to homogeneous and independent of the velocity of the observer, it manifests itself as a positive contribution to Einstein's Λ, or dark energy. The zero-point energies of the fermions make a negative contribution. Other contributions, perhaps including the energy densities of fields that interact only with themselves and

[3]The model is named after the nonbaryonic cold dark matter that is assumed to dominate the masses of galaxies in the present universe. There are more assumptions in the CDM model, of course; they are discussed in Secs. III.D and IV.A.2.

[4]At the time of writing the Microwave Anisotropy Probe (MAP) satellite is collecting data; the project is described in http://map.gsfc.nasa.gov/

[5]See Bordag, Mohideen, and Mostepanenko (2001) for a recent review. The attractive Casimir force between two parallel conducting plates results from the boundary condition that suppresses the number of modes of oscillation of the electromagnetic field between the plates, thus suppressing the energy of the system. One can understand the effect at small separation without reference to the quantum behavior of the electromagnetic field, such as in the analysis of the van der Waals interaction in quantum mechanics, by taking account of the term in the particle Hamiltonian for the Coulomb potential energy between the charged particles in the two separate neutral objects. But a more complete treatment, as discussed by Cohen-Tannoudji, Dupont-Roc, and Grynberg (1992), replaces the Coulomb interaction with the coupling of the charged particles to the electromagnetic-field operator. In this picture the van der Waals interaction is mediated by the exchange of virtual photons. With either way of looking at the Casimir effect—the perturbation of the normal modes or the exchange of virtual quanta of the unperturbed modes—the effect is the same, the suppression of the energy of the system.

gravity, might have either sign. The value of the sum suggested by dimensional analysis is much larger than what is allowed by the relativistic cosmological model. The only other natural value is $\Lambda=0$. If Λ really is tiny but not zero, this introduces a most stimulating though enigmatic clue to the physics yet to be discovered.

To illustrate the problem we outline an example of a contribution to Λ. The energy density in the 3-K thermal cosmic microwave background radiation, which amounts to $\Omega_{R0}\sim 5\times 10^{-5}$ in Eq. (1) (ignoring the neutrinos), peaks at wavelength $\lambda\sim 2$ mm. At this Wien peak the photon occupation number is about one-fifteenth. The zero-point energy amounts to half the energy of a photon at the given frequency. This means the zero-point energy in the electromagnetic field at wavelengths $\lambda\sim 2$ mm amounts to a contribution $\delta\Omega_{\Lambda 0}\sim 4\times 10^{-4}$ to the density parameter in Λ or the dark energy. The sum over the modes scales as λ^{-4} [as illustrated in Eq. (37)]. Thus a naive extrapolation to visible wavelengths determines that the contribution amounts to $\delta\Omega_{\Lambda 0}\sim 5\times 10^{10}$, already a ridiculous number.

The situation can be compared to the development of the theory of weak interactions. The Fermi pointlike interaction model is strikingly successful for a considerable range of energies, but it was clear from the start that the model fails at high energy. A fix was discussed—mediate the interaction by an intermediate boson—and eventually incorporated into the even more successful electroweak theory. General relativity and quantum mechanics are extremely successful over a considerable range of length scales, provided we agree not to use the rules of quantum mechanics to count the zero-point energy density in the vacuum, even though we know we have to count the zero-point energies in all other situations. There are thoughts on improving the situation, though they seem to be less focused than was the case for the Fermi model. Perhaps a new energy component spontaneously cancels the vacuum energy density or the new component varies slowly with position and here and there happens to cancel the vacuum energy density well enough to allow observers like us to exist. Whatever the nature of the more perfect theory, it must reproduce the successes of general relativity and quantum mechanics. That includes the method of representing the material content of the observable universe—all forms of mass and energy—by the stress-energy tensor, and the relation between the stress-energy tensor and the curvature of macroscopic spacetime. One part has to be adjusted.

The numerical values of the parameters in Eq. (1) also are enigmatic, and possibly trying to tell us something. The evidence is that the parameters have the approximate values

$$\Omega_{\Lambda 0}\sim 0.7, \quad \Omega_{DM0}\sim 0.2, \quad \Omega_{B0}\sim 0.05. \tag{2}$$

We have written Ω_{M0} in two parts: Ω_{B0} measures the density of the baryons we know exist and Ω_{DM0} measures the hypothetical nonbaryonic cold dark matter we need to fit the cosmological tests. The parameters Ω_{B0} and Ω_{DM0} have similar values but represent different things—baryonic and nonbaryonic matter—and $\Omega_{\Lambda 0}$, which is thought to represent something completely different, is not much larger. Also, if the parameters were measured when the universe was one-tenth its present size the time-independent Λ parameter would contribute $\Omega_{\Lambda}\sim 0.003$. That is, we seem to have come on the scene just as Λ has become an important factor in the expansion rate. These curiosities surely are in part accidental, but maybe in part physically significant. In particular, one might imagine that the dark-energy density represented by Λ is rolling to its natural value, zero, but is very small now because we measure it when the universe is very old. We shall discuss efforts along this line to at least partially rationalize the situation.

C. Some explanations

We have to explain our choice of nomenclature. Basic concepts of physics say that space contains homogeneous zero-point energy, and perhaps also energy that is homogeneous or nearly so in other forms, real or effective (such as from counter terms in gravity physics, which make the net energy density cosmologically acceptable). In the literature this near homogeneous energy has been termed vacuum energy, the sum of vacuum energy and quintessence (Caldwell, Davé, and Steinhardt, 1998), and dark energy (Turner, 1999). We have adopted the last term, and we shall refer to the dark-energy density ρ_{Λ} that manifests itself as an effective version of Einstein's cosmological constant, but one that may vary slowly with time and position.[6]

Our subject involves two quite different traditions, in physics and astronomy. Each has familiar notation, and familiar ideas that may be "in the air" but not in recent literature. Our attempt to take account of these traditions commences with the summary in Sec. II of the basic notation with brief explanations. We expect that readers will find some of these concepts trivial and others of some use, and that the useful parts will be different for different readers.

We offer in Sec. III our reading of the history of ideas on Λ and its generalization to dark energy. This is a fascinating and we think edifying illustration of how science may advance in unexpected directions. It is relevant to an understanding of the present state of research in cosmology, because traditions inform opinions, and people have had mixed feelings about Λ ever since Einstein (1917) introduced it 85 years ago. The concept never entirely disappeared in cosmology because a series of observations hinted at its presence, and because to some cosmologists Λ fits the formalism too well to be ignored. The search for the physics of the vacuum, and its possible relation to Λ, has a long history too. Despite

[6]The dark energy should of course be distinguished from a hypothetical gas of particles with velocity dispersion large enough that the distribution is close to homogeneous.

the common and strong suspicion that Λ must be negligibly small, because any other acceptable value is absurd, all this history has made the community well prepared for the recent observational developments that argue for the detection of Λ.

Our approach in Sec. IV to the discussion of the evidence for detection of Λ, from the cosmological tests, also requires explanation. One occasionally reads that the tests will show us how the world will end. That certainly seems interesting, but it is not the main point: why should we trust an extrapolation into the indefinite future of a theory that we can at best show is a good approximation to reality?[7] As we remarked in Sec. I.A, the purpose of the tests is to check the approximation to reality, by checking the physics and astronomy of the standard relativistic cosmological model, along with any viable alternatives that may be discovered. We take our task to be the identification of the aspects of the standard theory that enter the interpretation of the measurements and thus are or may be empirically checked or measured.

II. BASIC CONCEPTS

A. The Friedmann-Lemaître model

The standard world model is close to homogeneous and isotropic on large scales, and lumpy on small scales—the effect of mass concentrations in galaxies, stars, people, etc. The length scale at the transition from nearly smooth to strongly clumpy is about 10 Mpc. We use here and throughout the standard astronomers' length unit,

$$1\ \mathrm{Mpc}=3.1\times 10^{24}\ \mathrm{cm}=3.3\times 10^{6}\ \text{light years}. \tag{3}$$

To be more definite, imagine that many spheres of radius 10 Mpc are placed at random, and the mass within each is measured. At this radius the rms fluctuation in the set of values of masses is about equal to the mean value. On smaller scales the departures from homogeneity are progressively more nonlinear; on larger scales the density fluctuations are perturbations to the homogeneous model. From now on we mention these perturbations only when relevant for the cosmological tests.

The expansion of the universe means the distance $l(t)$ between two well-separated galaxies varies with world time t as

$$l(t)\propto a(t), \tag{4}$$

where the expansion or scale factor $a(t)$ is independent of the choice of galaxies. It is an interesting exercise, for those who have not already thought to do so, to check that Eq. (4) is required to preserve homogeneity and isotropy.[8]

The rate of change of the distance in Eq. (4) is the speed

$$v=dl/dt=Hl, \quad H=\dot{a}/a, \tag{5}$$

where the overdot means the derivative with respect to world time t and H is the time-dependent Hubble parameter. When v is small compared to the speed of light this is Hubble's law. The present value of H is Hubble's constant, H_0. When needed we will use[9]

$$H_0=100h\ \ \mathrm{km\,s^{-1}\,Mpc^{-1}}=67\pm 7\ \ \mathrm{km\,s^{-1}\,Mpc^{-1}} \\ =(15\pm 2\ \ \mathrm{Gyr})^{-1}, \tag{6}$$

at two standard deviations. The first equation defines the dimensionless parameter h.

Another measure of the expansion follows by considering the stretching of the wavelength of light received from a distant galaxy. The observed wavelength λ_{obs} of a feature in the spectrum that had wavelength λ_{em} at emission satisfies

$$1+z=\lambda_{\mathrm{obs}}/\lambda_{\mathrm{em}}=a(t_{\mathrm{obs}})/a(t_{\mathrm{em}}), \tag{7}$$

[7] Observations may now have detected Λ, at a characteristic energy scale of a millielectron volt [Eq. (47)]. We have no guarantee that an even lower-energy scale does not exist; such a scale could first become apparent through cosmological tests.

[8] We feel we have to comment on a few details about Eq. (4) to avoid contributing to debates that are more intense than seem warranted. Think of the world time t as the proper time kept by each of a dense set of observers, each moving so that all the others are isotropically moving away, and with the times synchronized to a common energy density, $\rho(t)$, in the near homogeneous expanding universe. The distance $l(t)$ is the sum of the proper distances between neighboring observers, all measured at time t, and along the shortest distance between the two observers. The rate of increase of the distance, dl/dt, may exceed the velocity of light. This is no more problematic in relativity theory than is the large speed at which the beam of a flashlight on Earth may swing across the face of the Moon (assuming an adequately tight beam). Space sections at fixed t may be noncompact, and the total mass of a homogeneous universe formally infinite. As far as is known this is not meaningful: we can only assert that the universe is close to homogeneous and isotropic over observable scales, and that what can be observed is a finite number of baryons and photons.

[9] The numerical values in Eq. (6) are determined from an analysis of all available measurements of H_0 prior to mid-1999 (Gott *et al.*, 2001). They are a very reasonable summary of the current situation. For instance, the Hubble Space Telescope Key Project summary measurement value $H_0=72\pm 8\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$ (1σ uncertainty; Freedman *et al.*, 2001) is in very good agreement with Eq. (6), as is the recent Tammann *et al.* (2001) summary value $H_0=60\pm 6\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$ (approximate 1σ systematic uncertainty). This is an example of the striking change in the observational situation over the previous five years: the uncertainty in H_0 has decreased by more than a factor of 3, making it one of the better-measured cosmological parameters.

where the expansion factor a is defined in Eq. (4) and z is the redshift. That is, the wavelength of freely traveling radiation stretches in proportion to the factor by which the universe expands. To understand this, imagine that a large part of the universe is enclosed in a cavity with perfectly reflecting walls. The cavity expands with the general expansion, the widths proportional to $a(t)$. Electromagnetic radiation is a sum of the normal modes that fit the cavity. At interesting wavelengths the mode frequencies are much larger than the rate of expansion of the universe, so adiabaticity says a photon in a mode stays there, and its wavelength thus must vary as $\lambda \propto a(t)$, as stated in Eq. (7). The cavity disturbs the long-wavelength part of the radiation, but the disturbance can be made exceedingly small by choosing a large cavity.

Equation (7) defines the redshift z. The redshift is a convenient label for epochs in the early universe, where z exceeds unity. A good exercise for the student is to check that when z is small Eq. (7) reduces to Hubble's law, where λz is the first-order Doppler shift in the wavelength λ, and Hubble's parameter H is given by Eq. (5). Thus Hubble's law may be written as $cz=Hl$ (where we have put in the speed of light).

These results follow from the symmetry of the cosmological model and conventional local physics; we do not need general relativity theory. When $z \gtrsim 1$ we need relativistic theory to compute the relations among the redshift and other observables. An example is the relation between redshift and apparent magnitude used in the supernova test. Other cosmological tests check consistency among these relations, and this checks the world model.

In general relativity the second time derivative of the expansion factor satisfies

$$\frac{\ddot{a}}{a}=-\frac{4}{3}\pi G(\rho+3p). \tag{8}$$

The gravitational constant is G. Here and throughout we choose units to set the velocity of light to unity. The mean mass density, $\rho(t)$, and the pressure, $p(t)$, counting all contributions including dark energy, satisfy the local energy-conservation law

$$\dot{\rho}=-3\frac{\dot{a}}{a}(\rho+p). \tag{9}$$

The first term on the right-hand side represents the decrease of mass density due to the expansion that more broadly disperses the matter. The pdV work in the second term is a familiar local concept, and is meaningful in general relativity. But one should note that energy does not have a general global meaning in this theory.

The first integral of Eqs. (8) and (9) is the Friedmann equation

$$\dot{a}^2=\frac{8}{3}\pi G\rho a^2+\text{const}. \tag{10}$$

It is conventional to rewrite this as

$$\begin{aligned}\left(\frac{\dot{a}}{a}\right)^2 &= H_0^2E(z)^2\\ &= H_0^2[\Omega_{\mathrm{M0}}(1+z)^3+\Omega_{\mathrm{R0}}(1+z)^4\\ &\quad +\Omega_{\Lambda 0}+\Omega_{\mathrm{K0}}(1+z)^2].\end{aligned} \tag{11}$$

The first equation defines the function $E(z)$ that is introduced for later use. The second equation assumes constant Λ; the time-dependent dark-energy case is reviewed in Secs. II.C and III.E. The first term in the last part of Eq. (11) represents nonrelativistic matter with negligibly small pressure; one sees from Eqs. (7) and (9) that the mass density in this form varies with the expansion of the universe as $\rho_{\mathrm{M}}\propto a^{-3}\propto(1+z)^3$. The second term represents radiation and relativistic matter, with pressure $p_{\mathrm{R}}=\rho_{\mathrm{R}}/3$, whence $\rho_{\mathrm{R}}\propto(1+z)^4$. The third term is the effect of Einstein's cosmological constant, or a constant dark-energy density. The last term, discussed in more detail below, is the constant of integration in Eq. (10). The four density parameters Ω_{i0} are the fractional contributions to the square of Hubble's constant, H_0^2, that is, $\Omega_{i0}(t)=8\pi G\rho_{i0}/(3H_0^2)$. At the present epoch, $z=0$, the present value of $\dot{a}/a$ is H_0, and the Ω_{i0} sum to unity [Eq. (1)].

In this notation, Eq. (8) is

$$\frac{\ddot{a}}{a}=-H_0^2[\Omega_{\mathrm{M0}}(1+z)^3/2+\Omega_{\mathrm{R0}}(1+z)^4-\Omega_{\Lambda 0}]. \tag{12}$$

The constant of integration in Eqs. (10) and (11) is related to the geometry of spatial sections at constant world time. Recall that in general relativity events in spacetime are labeled by the four coordinates x^μ of time and space. Neighboring events at separation dx^μ have invariant separation ds defined by the line element

$$ds^2=g_{\mu\nu}dx^\mu dx^\nu. \tag{13}$$

The repeated indices are summed, and the metric tensor $g_{\mu\nu}$ is a function of position in spacetime. If ds^2 is positive then ds is the proper (physical) time measured by an observer who moves from one event to the other; if negative, $|ds|$ is the proper distance between the events measured by an observer who is moving so the events are seen to be simultaneous.

In the flat spacetime of special relativity one can choose coordinates so the metric tensor has the Minkowskian form

$$\eta_{\mu\nu}=\begin{pmatrix}1&0&0&0\\0&-1&0&0\\0&0&-1&0\\0&0&0&-1\end{pmatrix}. \tag{14}$$

A freely falling, inertial observer can choose locally Minkowskian coordinates, such that along the path of the observer $g_{\mu\nu}=\eta_{\mu\nu}$, and the first derivatives of $g_{\mu\nu}$ vanish.

In the homogeneous world model we can choose coordinates so the metric tensor is of the form that results in the line element

$$ds^2 = dt^2 - a(t)^2\left[\frac{dr^2}{1+Kr^2} + r^2(d\theta^2 + \sin^2\theta d\phi^2)\right]$$
$$= dt^2 - K^{-1}a(t)^2[d\chi^2 + \sinh^2\chi(d\theta^2 + \sin^2\theta d\phi^2)]. \quad (15)$$

In the second expression, which assumes $K>0$, the radial coordinate is $r = K^{-1/2}\sinh\chi$. The expansion factor $a(t)$ appears in Eq. (4). If a were constant and the constant K vanished this would represent the flat spacetime of special relativity in polar coordinates. The key point for now is that Ω_{K0} in Eq. (11), which represents the constant of integration in Eq. (10), is related to the constant K:

$$\Omega_{K0} = K/(H_0 a_0)^2, \quad (16)$$

where a_0 is the present value of the expansion factor $a(t)$. Cosmological tests that are sensitive to the geometry of space constrain the value of the parameter Ω_{K0}, and Ω_{K0} and the other density parameters Ω_{i0} in Eq. (11) determine the expansion history of the universe.

It is useful for what follows to recall that the metric tensor in Eq. (15) satisfies Einstein's field equation, a differential equation we can write as

$$G_{\mu\nu} = 8\pi G T_{\mu\nu}. \quad (17)$$

The left side is a function of $g_{\mu\nu}$ and its first two derivatives; it represents the geometry of spacetime. The stress-energy tensor $T_{\mu\nu}$ represents the material contents of the universe, including particles, radiation, fields, and zero-point energies. An observer in a homogeneous and isotropic universe, moving so the universe is observed to be isotropic, would measure the stress-energy tensor to be

$$T_{\mu\nu} = \begin{pmatrix} \rho & 0 & 0 & 0 \\ 0 & p & 0 & 0 \\ 0 & 0 & p & 0 \\ 0 & 0 & 0 & p \end{pmatrix}. \quad (18)$$

This diagonal form is a consequence of the symmetry; the diagonal components define the pressure and energy density. With Eq. (18), the differential equation (17) yields the expansion-rate equations (11) and (12).

B. The cosmological constant

Special relativity is very successful in laboratory physics. Thus one might guess that any inertial observer would see the same vacuum. A freely moving inertial observer represents spacetime in the neighborhood by locally Minkowskian coordinates, with the metric tensor $\eta_{\mu\nu}$ given in Eq. (14). A Lorentz transformation to an inertial observer with another velocity does not change this Minkowski form. The same must be true of the stress-energy tensor of the vacuum, if all observers see the same vacuum, so it has to be of the form

$$T^{\Lambda}_{\mu\nu} = \rho_\Lambda g_{\mu\nu}, \quad (19)$$

where ρ_Λ is a constant, in a general coordinate labeling. When writing this contribution to the stress-energy tensor separately from the rest, we bring the field equation (17) to

$$G_{\mu\nu} = 8\pi G(T_{\mu\nu} + \rho_\Lambda g_{\mu\nu}). \quad (20)$$

This is Einstein's (1917) revision of the field equation of general relativity, where ρ_Λ is proportional to his cosmological constant Λ; his reason for writing down this equation is discussed in Sec. III.A. In many dark-energy scenarios ρ_Λ is a slowly varying function of time and its stress-energy tensor differs slightly from Eq. (19), so the observed properties of the vacuum depend on the observer's velocity.

One sees from Eqs. (14), (18), and (19) that the new component in the stress-energy tensor resembles an ideal fluid with negative pressure,

$$p_\Lambda = -\rho_\Lambda. \quad (21)$$

This fluid picture is of limited use, but the following properties are worth noting.[10]

The stress-energy tensor of an ideal fluid with four-velocity u^μ generalizes from Eq. (18) to $T^{\mu\nu} = (\rho + p)u^\mu u^\nu - pg^{\mu\nu}$. The equations of fluid dynamics follow from the vanishing of the divergence of $T^{\mu\nu}$. Let us consider the simple case of locally Minkowskian coordinates, meaning free fall, and a fluid that is close to homogeneous. By the latter we mean the fluid velocity $\vec{v}$—the space part of the four-velocity u^μ—and the density fluctuation $\delta\rho$ from homogeneity may be treated in linear perturbation theory. Then the equations of energy and momentum conservation are

$$\delta\dot{\rho} + (\langle\rho\rangle + \langle p\rangle)\nabla\cdot\vec{v} = 0, \quad (\langle\rho\rangle + \langle p\rangle)\dot{\vec{v}} + c_s^2\nabla\delta\rho = 0, \quad (22)$$

where $c_s^2 = dp/d\rho$ and the mean density and pressure are $\langle\rho\rangle$ and $\langle p\rangle$. These combine to

$$\delta\ddot{\rho} = c_s^2\nabla^2\delta\rho. \quad (23)$$

If c_s^2 is positive this is a wave equation, and c_s is the speed of sound.

The first of Eqs. (22) is the local energy-conservation law, as in Eq. (9). If $p = -\rho$, the pdV work cancels the ρdV part: the work done to increase the volume cancels the effect of the increased volume. This has to be so for a Lorentz-invariant stress-energy tensor, of course, where all inertial observers see the same vacuum. Another way to see this is to note that the energy flux density in Eqs. (22) is $(\langle\rho\rangle + \langle p\rangle)\vec{v}$. This vanishes when

[10]These arguments have been familiar, in some circles, for a long time, though in our experience, discussed more often in private than in the literature. Early statements of elements are in Lemaître (1934) and McCrea (1951); see Kragh (1999, pp. 397 and 398) for a brief historical account.

$p = -\rho$: the streaming velocity loses meaning. When c_s^2 is negative Eq. (23) shows that the fluid is unstable, in general. But when $p = -\rho$ the vanishing divergence of $T^{\mu\nu}$ becomes the condition shown in Eq. (22), that $\rho = \langle\rho\rangle + \delta\rho$ is constant.

There are two measures of gravitational interactions with a fluid: the passive gravitational mass density determines how the fluid streaming velocity is affected by an applied gravitational field, and the active gravitational mass density determines the gravitational field produced by the fluid. When the fluid velocity is nonrelativistic the expression for the former in general relativity is $\rho + p$, as one can determine by writing out the covariant divergence of $T^{\mu\nu}$. This vanishes when $p = -\rho$, consistent with the loss of meaning of the streaming velocity. The latter is $\rho + 3p$, as one can see from Eq. (8). Thus a fluid with $p = -\rho/3$, if somehow kept homogeneous and static, would produce no gravitational field.[11] In the model in Eqs. (19) and (21) the active gravitational mass density is negative when ρ_Λ is positive. When this positive ρ_Λ dominates the stress-energy tensor, $\ddot{a}$ is positive: the rate of expansion of the universe increases. In the language of Eq. (20), this cosmic repulsion is a gravitational effect of the negative active gravitational mass density, not a new force law.

The homogeneous active mass represented by Λ changes the equation of relative motion of freely moving test particles in the nonrelativistic limit to

$$\frac{d^2\vec{r}}{dt^2} = \vec{g} + \Omega_{\Lambda 0} H_0^2 \vec{r}, \tag{24}$$

where $\vec{g}$ is the relative gravitational acceleration produced by the distribution of ordinary matter.[12] For an illustration of the size of the last term consider its effect on our motion in a nearly circular orbit around the center of the Milky Way galaxy. The Solar System is moving at speed $v_c = 220\ \mathrm{km\,s^{-1}}$ at radius $r = 8$ kpc. The ratio of the acceleration g_Λ produced by Λ to the total gravitational acceleration $g = v_c^2/r$ is

$$g_\Lambda / g = \Omega_{\Lambda 0} H_0^2 r^2 / v_c^2 \sim 10^{-5}, \tag{25}$$

a small number. Since we are near the edge of the luminous part of our galaxy, a search for the effect of Λ on the internal dynamics of galaxies such as the Milky Way does not look promising. The precision of celestial dynamics in the Solar System is much greater, but the effect of Λ is very much smaller; $g_\Lambda / g \sim 10^{-22}$ for the orbit of the Earth.

One can generalize Eq. (19) to a variable ρ_Λ, by taking p_Λ to be negative but different from $-\rho_\Lambda$. But if the dynamics were that of a fluid, with pressure a function of ρ_Λ, stability would require $c_s^2 = dp_\Lambda / d\rho_\Lambda > 0$, from Eq. (23), which seems quite contrived. A viable working model for a dynamical ρ_Λ is the dark energy of a scalar field with self-interaction potential chosen to make the variation of the field energy acceptably slow, as discussed next.

C. Inflation and dark energy

The negative active gravitational mass density associated with a positive cosmological constant is an early precursor of the inflation picture of the early universe; inflation in turn is one precursor of the idea that Λ might generalize into evolving dark energy.

To begin, we review some aspects of causal relations between events in spacetime. Neglecting space curvature, a light ray moves a proper distance $dl = a(t)dx = dt$ in time interval dt, so the integrated coordinate displacement is

$$x = \int dt/a(t). \tag{26}$$

If $\Omega_{\Lambda 0} = 0$ this integral converges in the past—we see distant galaxies that at the time of observation cannot have seen us since the singular start of expansion at $a = 0$. This "particle horizon problem" is curious: how could distant galaxies in different directions in the sky know to look so similar? The inflation idea is that in the early universe the expansion history approximates that of de Sitter's (1917) solution to Einstein's field equation for $\Lambda > 0$ and $T_{\mu\nu} = 0$ in Eq. (20). We can choose the coordinate labels in this de Sitter spacetime so space curvature vanishes. Then Eqs. (11) and (12) show that the expansion parameter is

$$a \propto e^{H_\Lambda t}, \tag{27}$$

where H_Λ is a constant. As one sees by working the integral in Eq. (26), here everyone can have seen everyone else in the past. The details need not concern us; for the following discussion two concepts are important. First, the early universe acts like an approximation to de Sitter's solution because it is dominated by a large effective cosmological "constant," or dark-energy density. Second, the dark energy is modeled as that of a near homogeneous field, Φ.

In this scalar field model, motivated by grand unified models of very-high-energy particle physics, the action of the real scalar field Φ (in units chosen so that Planck's constant $\hbar$ is unity) is

[11]Lest we contribute to a wrong problem for the student we note that a fluid with $p = -\rho/3$ held in a container would have net positive gravitational mass, from the pressure in the container walls required for support against the negative pressure of the contents. We have finessed the walls by considering a homogeneous situation. We believe Whittaker (1935) gives the first derivation of the relativistic active gravitational mass density. Whittaker also presents an example of the general proposition that the active gravitational mass of an isolated stable object is the integral of the time-time part of the stress-energy tensor in the locally Minkowskian rest frame. Misner and Putman (1959) give the general demonstration.

[12]This assumes that the particles are close enough for application of the ordinary operational definition of proper relative position. The parameters in the last term follow from Eqs. (8) and (21).

$$S=\int d^4x\sqrt{-g}\left[\frac{1}{2}g^{\mu\nu}\partial_\mu\Phi\partial_\nu\Phi-V(\Phi)\right]. \tag{28}$$

The potential-energy density V is a function of the field Φ, and g is the determinant of the metric tensor. When the field is spatially homogeneous [in the line element of Eq. (15)], and space curvature may be neglected, the field equation is

$$\ddot{\Phi}+3\frac{\dot{a}}{a}\dot{\Phi}+\frac{dV}{d\Phi}=0. \tag{29}$$

The stress-energy tensor of this homogeneous field is diagonal (in the rest frame of an observer moving so that the universe is seen to be isotropic), with time and space parts along the diagonal

$$\rho_\Phi=\frac{1}{2}\dot{\Phi}^2+V(\Phi),$$

$$p_\Phi=\frac{1}{2}\dot{\Phi}^2-V(\Phi). \tag{30}$$

If the scalar field varies slowly in time, so that $\dot{\Phi}^2\ll V$, the field energy approximates the effect of Einstein's cosmological constant, with $p_\Phi\simeq-\rho_\Phi$.

The inflation picture assumes that the near exponential expansion of Eq. (27) in the early universe lasts long enough so that every bit of the present observable universe has seen every other bit, and presumably has discovered how to relax to almost exact homogeneity. The field Φ may then start varying rapidly enough to produce the entropy of our universe, and the field or the entropy may produce the baryons, leaving ρ_Φ small or zero. But one can imagine that the late time evolution of ρ_Φ is slow. If it is slower than the evolution in the mass density in matter, there comes a time when ρ_Φ again dominates, and the universe appears to have a cosmological constant.

A model for this late time evolution assumes a potential of the form

$$V=\kappa/\Phi^\alpha, \tag{31}$$

where the constant κ has dimensions of mass raised to the power $\alpha+4$. For simplicity let us suppose the universe after inflation, but at high redshift, is dominated by matter or radiation, with mass density ρ, that drives the power-law expansion, $a\propto t^n$. Then the power-law solution to the field Eq. (29) with the potential in Eq. (31) is

$$\Phi\propto t^{2/(2+\alpha)}, \tag{32}$$

and the ratio of the mass densities in the scalar field and in matter or radiation is

$$\rho_\Phi/\rho\propto t^{4/(2+\alpha)}. \tag{33}$$

In the limit at which the parameter α approaches zero, ρ_Φ is constant, and this model is equivalent to Einstein's Λ.

When $\alpha>0$ the field Φ in this model grows arbitrarily large at large time, so $\rho_\Phi\rightarrow 0$, and the universe approaches the Minkowskian spacetime of special relativity. This is within a simple model, of course. It is easy to imagine that in other models ρ_Φ approaches a constant positive value at large time, and spacetime approaches the de Sitter solution, or ρ_Φ passes through zero and becomes negative, causing spacetime to collapse to a Big Crunch.

The power-law model with $\alpha>0$ has two properties that seem desirable. First, the solution in Eq. (32) is said to be an *attractor* (Ratra and Peebles, 1988) or a *tracker* (Steinhardt, Wang, and Zlatev, 1999), meaning it is the asymptotic solution for a broad range of initial conditions at high redshift. That includes relaxation to a near homogeneous energy distribution even when gravity has collected the other matter into nonrelativistic clumps. Second, the energy density in the attractor solution decreases less rapidly than that of matter and radiation. This allows us to realize the scenario: after inflation but at high redshift the field energy density ρ_Φ is small so it does not disturb the standard model for the origin of the light elements, but eventually ρ_Φ dominates and the universe acts as if it had a cosmological constant, but one that varies slowly with position and time. We comment on details of this model in Sec III.E.

III. HISTORICAL REMARKS

These comments on what people were thinking are gleaned from the literature and supplemented by private discussions and our own recollections. More is required for a complete history of the subject, of course, but we hope we have captured the main themes and the way in which these themes have evolved into the present appreciation of the situation.

A. Einstein's thoughts

Einstein disliked the idea of an island universe in asymptotically flat spacetime, because a particle could leave the island and move arbitrarily far from all the other matter in the universe, yet preserve all its inertial properties, which he considered a violation of Mach's idea of the relativity of inertia. Einstein's (1917) cosmological model accordingly assumes that the universe is homogeneous and isotropic, on average, thus removing the possibility of arbitrarily isolated particles. Einstein had no empirical support for this assumption, yet it agrees with modern precision tests. There is no agreement as to whether this is more than a lucky guess.

Motivated by the observed low velocities of the then known stars, Einstein assumed that the large-scale structure of the universe is static. He introduced the cosmological constant to reconcile this picture with his general relativity theory. In the notation of Eq. (12), one sees that a positive value of $\Omega_{\Lambda0}$ can balance the positive values of Ω_{M0} and Ω_{R0} for consistency with $\ddot{a}=0$. The balance is unstable: a small perturbation to the mean mass density or the mass distribution causes expansion or contraction of the whole or parts of the universe. One sees this in Eq. (24): the mass distribution can be chosen so the two terms on the right-hand side cancel, but the balance can be upset by redistributing the mass.[13]

Einstein did not consider the cosmological constant to be part of the stress-energy term: his form for the field equation [in the streamlined notation of Eq. (17)] is

$$G_{\mu\nu}-8\pi G\rho_\Lambda g_{\mu\nu}=8\pi G T_{\mu\nu}. \quad (34)$$

The left-hand side contains the metric tensor and its derivatives; a new constant of nature, Λ, appears in the addition to Einstein's original field equation. One can equally place Einstein's new term on the right-hand side of the equation, as in Eq. (20), and count $\rho_\Lambda g_{\mu\nu}$ as part of the source term in the stress-energy tensor. The distinction becomes interesting when ρ_Λ takes part in the dynamics, and the field equation is properly written with ρ_Λ, or its generalization, as part of the stress-energy tensor. One would then be able to say that the differential equation of gravity physics has not changed from Einstein's original form; instead there is a new component in the content of the universe.

Having assumed that the universe is static, Einstein did not write down the differential equation for $a(t)$, and so did not see the instability. Friedmann (1922, 1924) found the evolving homogeneous solution, but had the misfortune to do so before the astronomy became suggestive. Slipher's measurements of the spectra of the spiral nebulae—galaxies of stars—showed most are shifted toward the red, and Eddington (1924, pp. 161 and 162) remarked that that might be a manifestation of the second, repulsive term in Eq. (24). Lemaître (1927) introduced the relation between Slipher's redshifts and a homogeneous matter-filled expanding relativistic world model. He may have been influenced by Hubble's work, which led to the publication (Hubble, 1929) of the linear redshift-distance relation [Eq. (5)]: as a graduate student at MIT Lemaître attended a lecture by Hubble.

In Lemaître's (1927) solution, the expanding universe traces asymptotically back to Einstein's static case. Lemaître then turned to what he called the primeval atom, which is now termed the Big Bang model. This solution expands from densities so large that they require some sort of quantum treatment, passes through a quasistatic approximation to Einstein's solution, and then continues expanding to de Sitter's (1917) empty space solution. To modern tastes, this "loitering" model requires incredibly special initial conditions, as will be discussed. Lemaître liked it because the loitering epoch allows the expansion time to be acceptably long for Hubble's (1929) estimate of H_0, which is an order-of-magnitude high.

The record shows Einstein never liked the Λ term. His view of how general relativity might fit Mach's principle was disturbed by de Sitter's (1917) solution to Eq. (34) for empty space ($T_{\mu\nu}=0$) with $\Lambda>0$.[14,15] Pais (1982, p. 288) pointed out that Einstein, in a letter to Weyl in 1923, commented on the effect of Λ in Eq. (24): "According to De Sitter, two material points that are sufficiently far apart, continue to be accelerated and move apart. If there is no quasistatic world, then away with the cosmological term." We do not know whether at this time Einstein was influenced by Slipher's redshifts or Friedmann's expanding world model.

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[13]To help motivate the introduction of Λ, Einstein (1917) mentioned a modification of Newtonian gravity physics that could render the theory well defined when the mass distribution is homogeneous. In Einstein's example, similar to what was considered by Seeliger and Neumann in the mid-1890s, the modified field equation for the gravitational potential φ is $\nabla^2\varphi-\lambda\varphi=4\pi G\rho_{\rm M}$. This allows the nonsingular homogeneous static solution $\varphi=-4\pi G\rho_{\rm M}/\lambda$. In this example the potential for an isolated point mass is the Yukawa form, $\varphi \propto e^{-\sqrt{\lambda}r}/r$. Trautman (1965) pointed out that this is not the nonrelativistic limit of general relativity with the cosmological term. Rather, Eq. (24) follows from $\nabla^2\varphi=4\pi G(\rho_{\rm M}-2\rho_\Lambda)$, where the active gravitational mass density of the Λ term is $\rho_\Lambda+3p_\Lambda=-2\rho_\Lambda$. Norton (1999) reviewed the history of ideas of this Seeliger-Neumann Yukawa-type potential in gravity physics.

[14]North (1965) reviews the confused early history of ideas on the possible astronomical significance of de Sitter's solution for an empty universe with $\Lambda>0$; we add a few comments regarding the physics that contributed to the discovery of the expanding world model. Suppose an observer in de Sitter's spacetime holds a string tied to a source of light, so the source stays at fixed physical distance $r \ll H_\Lambda^{-1}$. The source is much less massive than the observer, the gravitational frequency shift due to the observer's mass may be neglected, and the observer is moving freely. Then the observer receives light from the source shifted to the red by $\delta\lambda/\lambda=-(H_\Lambda r)^2/2$. The observed redshifts of particles moving on geodesics depend on the initial conditions. Stars in the outskirts of our galaxy are held at fixed mean distances from Earth by their motions. The mean shifts of the spectra of light from these stars include this quadratic de Sitter term as well as the much larger Doppler and ordinary gravitational shifts. The prescription for initial conditions that reproduces the linear redshift-distance relation for distant galaxies follows Weyl's (1923) principle: the world particle geodesics trace back to a near common position in the remote past, in the limiting case of the Friedmann-Lemaître model at $\Omega_{\rm M0} \rightarrow 0$. This spatially homogeneous coordinate labeling of de Sitter's spacetime, with space sections with negative curvature, already appears in de Sitter [1917, Eq. (15)], and is repeated in Lanczos (1922). This line element is the second expression in our Eq. (15) with $a \propto \cosh H_\Lambda t$. Lemaître (1925) and Robertson (1928) present the coordinate labeling for the spatially flat case, where the line element is $ds^2=dt^2-e^{2H_\Lambda t}(dx^2+dy^2+dz^2)$ [in the choice of symbols and signature in Eqs. (15) and (27)]. Lemaître (1925) and Robertson (1928) note that particles at rest in this coordinate system present a linear redshift-distance relation, $v=H_\Lambda r$, at small v. Robertson (1928) estimated H_Λ, and Lemaître (1927) its analog for the Friedmann-Lemaître model, from published redshifts and Hubble's galaxy distances. Their estimates are not far off Hubble's (1929) published value.

[15]To the present way of thinking the lengthy debate about the singularity in de Sitter's static solution, chronicled by North (1965), seems surprising, because de Sitter (1917) and Klein (1918) had presented de Sitter's solution as a sphere embedded in 4-plus-1-dimensional flat space, with no physical singularity.

The earliest published comments we have found on Einstein's opinion of Λ within the evolving world model (Einstein, 1931; Einstein and de Sitter, 1932) make the point that, since not all the terms in the expansion-rate equation (11) are logically required, and the matter term surely is present and likely dominates over radiation at low redshift, a reasonable working model drops Ω_{K0} and $\Omega_{\Lambda 0}$ and ignores Ω_{R0}. This simplifies the expansion-rate equation to what has come to be called the Einstein–de Sitter model,

$$\frac{\dot{a}^2}{a^2}=\frac{8}{3}\pi G\rho_{\mathrm{M}}, \tag{35}$$

where ρ_{M} is the mass density in nonrelativistic matter; here $\Omega_{\mathrm{M}}=8\pi G\rho_{\mathrm{M}}/(3H^2)$ is unity. The left side is a measure of the kinetic energy of expansion per unit mass, and the right-hand side a measure of the negative of the gravitational potential energy. In effect, this model universe expands with escape velocity.

Einstein and de Sitter point out that Hubble's estimate of H_0 and de Sitter's estimate of the mean mass density in galaxies are not inconsistent with Eq. (35) (and since both quantities scale with distance in the same way, this result is not affected by the error in the distance scale that affected Hubble's initial measurement of H_0). But the evidence shows now that the mass density is about one-quarter of what is predicted by this equation, as we will discuss.

Einstein and de Sitter (1932) remarked that the curvature term in Eq. (11) is "essentially determinable, and an increase in the precision of the data derived from observations will enable us in the future to fix its sign and determine its value." This is happening, 70 years later. The cosmological constant term is measurable, in principle, too, and may now have been detected. But Einstein and de Sitter said only that the theory of an expanding universe with finite mean mass density "can be reached without the introduction of" Λ.

Further to this point, in the appendix of the second edition of his book, *The Meaning of Relativity*, Einstein (1945, p. 127) states that the "introduction of the 'cosmologic member' "—Einstein's terminology for Λ—"into the equations of gravity, though possible from the point of view of relativity, is to be rejected from the point of view of logical economy," and that if "Hubble's expansion had been discovered at the time of the creation of the general theory of relativity, the cosmologic member would never have been introduced. It seems now so much less justified to introduce such a member into the field equations, since its introduction loses its sole original justification,—that of leading to a natural solution of the cosmologic problem." Einstein knew that without the cosmological constant the expansion time derived from Hubble's estimate of H_0 is uncomfortably short compared to estimates of the ages of the stars, and opined that that might be a problem with the star ages. The big error, the value of H_0, was corrected by 1960 (Sandage, 1958, 1962).

Gamow (1970, p. 44) recalls that "when I was discussing cosmological problems with Einstein, he remarked that the introduction of the cosmological term was the biggest blunder he ever made in his life." This certainly is consistent with all of Einstein's written comments that we have seen on the cosmological constant *per se*; we do not know whether Einstein was also referring to the missed chance to predict the evolution of the universe.

B. The development of ideas

1. Early indications of Λ

In the classic book, *The Classical Theory of Fields*, Landau and Lifshitz (1951, p. 338) second Einstein's opinion of the cosmological constant Λ, stating there is "no basis whatsoever" for adjustment of the theory to include this term. The empirical side of cosmology is not much mentioned in this book, however (though there is a perceptive comment on the limited empirical support for the homogeneity assumption; p. 332). In the Supplementary Notes to the English translation of his book, *Theory of Relativity*, Pauli (1958, p. 220) also endorses Einstein's position.

Discussions elsewhere in the literature on how one might find empirical constraints on the values of the cosmological parameters usually take account of Λ. The continued interest was at least in part driven by indications that Λ might be needed to reconcile theory and observations. Here are three examples.

First, the expansion time is uncomfortably short if $\Lambda=0$. Sandage's recalibration of the distance scale in the 1960s indicates $H_0\simeq 75\ \mathrm{km\,s^{-1}\,Mpc^{-1}}$. If $\Lambda=0$ this shows that the time of expansion from densities too high for stars to have existed is $<H_0^{-1}\simeq 13$ Gyr, maybe less than the ages of the oldest stars, then estimated to be greater than about 15 Gyr. Sandage (1961a) points out that the problem is removed by adding a positive Λ. The present estimates reviewed below (Sec. IV.B.3) are not far from these numbers, but still too uncertain for a significant case for Λ.

Second, counts of quasars as a function of redshift show a peak at $z\sim 2$, as would be produced by the loitering epoch in Lemaître's Λ model (Petrosian, Salpeter, and Szekeres, 1967; Shklovsky, 1967; Kardashev, 1967). The peak is now well established, centered at $z\sim 2.5$ (Croom *et al.*, 2001; Fan *et al.*, 2001). It is usually interpreted as the evolution in the rate of violent activity in the nuclei of galaxies, though in the absence of a loitering epoch the indicated sharp variation in quasar activity with time is curious (but certainly could be a consequence of astrophysics that is not well understood).

The third example is the redshift-magnitude relation. Sandage's (1961a) analysis indicates this is a promising method of distinguishing world models. The Gunn and Oke (1975) measurement of this relation for giant elliptical galaxies, with Tinsley's (1972) correction for evolution of the star population from assumed formation at high redshift, indicates curvature away from the linear relation in the direction that, as Gunn and Tinsley

(1975) discuss, could only be produced by Λ (within general relativity theory). The new application of the redshift-magnitude test, to type-Ia supernovae (Sec. IV.B.4), is not inconsistent with the Gunn-Oke measurement; we do not know whether this agreement of the measurements is significant, because Gunn and Oke were worried about galaxy evolution.[16]

2. The coincidences argument against Λ

An argument against an observationally interesting value of Λ, from our distrust of accidental coincidences, has been in the air for decades, and became very influential in the early 1980s with the introduction of the inflation scenario for the very early universe.

If the Einstein–de Sitter model in Eq. (35) were a good approximation at the present epoch, an observer measuring the mean mass density and Hubble's constant when the age of the universe was one-tenth the present value, or ten times the present age, would reach the same conclusion, that the Einstein–de Sitter model is a good approximation. That is, we would flourish at a time that is not special in the course of evolution of the universe. If, on the other hand, two or more of the terms in the expansion-rate equation (11) made substantial contributions to the present value of the expansion rate, it would mean that we are present at a special epoch, because each term in Eq. (11) varies with the expansion factor in a different way. To put this in more detail, we imagine that the physics of the very early universe, when the relativistic cosmological model became a good approximation, set the values of the cosmological parameters. The initial values of the contributions to the expansion-rate equation had to have been very different from each other, and exceedingly specially fixed, to yield two Ω_{i0}'s with comparable values. This would be a most remarkable and unlikely coincidence. The multiple coincidences required for the near vanishing of $\dot{a}$ and $\ddot{a}$ at a redshift not much larger than unity makes an even stronger case against Lemaître's loitering model, with this line of argument.

The earliest published comment we have found on this point is by Bondi (1960, p. 166), in the second edition of his book *Cosmology*. Bondi notes the "remarkable property" of the Einstein–de Sitter model: the dimensionless parameter we now call Ω_M is independent of the time at which it is computed (since it is unity). The coincidences argument follows and extends Bondi's comment. It is presented in McCrea (1971, p. 151). When Peebles was a postdoctoral research associate, in the early 1960s, in R. H. Dicke's gravity research group, the coincidences argument was discussed, but published much later (Dicke, 1970, p. 62; Dicke and Peebles, 1979). We do not know its provenance in Dicke's group, whether from Bondi, McCrea, Dicke, or someone else. We would not be surprised to learn others had similar thoughts.

The coincidences argument is sensible but not a proof, of course. The discovery of the 3-K thermal cosmic microwave background radiation gave us a term in the expansion-rate equation that is down from the dominant one by four orders of magnitude, not such a large factor by astronomical standards. This might be counted as a first step away from the argument. From the dynamics of galaxies the evidence that Ω_{M0} is less than unity is another step (Peebles, 1984, p. 442; 1986). And yet another is the development of the evidence that the Λ and dark-matter terms differ by only a factor of 3 [Eq. (2)]. This last piece is the most curious, but the community has come to accept it, for the most part. The precedent makes Lemaître's loitering model more socially acceptable.

A socially acceptable value of Λ cannot be such as to make life impossible, of course.[17] But perhaps the most productive interpretation of the coincidences argument is that it demands a search for a more fundamental underlying model. This is discussed further in Sec. III.E and the Appendix.

3. Vacuum energy and Λ

Another tradition to consider is the relation between Λ and the vacuum or dark-energy density. In one approach to the motivation for the Einstein field equation, taken by McVittie (1956) and others, Λ appears as a constant of integration (of the expression for local conservation of energy and momentum). McVittie (1956, p. 35) emphasizes that, as a constant of integration, Λ "cannot be assigned any particular value on *a priori* grounds." Interesting variants of this line of thought are still under discussion (Weinberg, 1989; Unruh, 1989; and references therein).

The notion of Λ as a constant of integration may be related to the issue of the zero point of energy. In laboratory physics one measures and computes energy differences. But the net energy matters for gravity physics, and one can imagine that Λ represents the difference between the true energy density and the sum at which one arrives by laboratory physics. Eddington (1939) and Lemaître (1934, 1949) make this point.

[16]Early measurements of the redshift-magnitude relation were meant in part to test the Steady State cosmology of Bondi and Gold (1948) and Hoyle (1948). Since Steady State cosmology assumes spacetime is independent of time its line element has to have the form of the de Sitter solution with $\Omega_{K0}=0$ and the expansion parameter in Eq. (27). The measured curvature of the redshift-magnitude relation is in the direction predicted by Steady State cosmology. But this cosmology fails other tests discussed in Sec. IV.B.

[17]If Λ were negative and the magnitude too large there would not be enough time for the emergence of life such as ours. If Λ were positive and too large the universe would expand too rapidly to allow galaxy formation. Our existence, which requires something resembling the Milky Way galaxy to contain and recycle heavy elements, thus provides an upper bound on the value of Λ. Such anthropic considerations are discussed by Weinberg (1987, 2001) and Vilenkin (2001), and references therein.

Bronstein[18] (1933) carries the idea further, allowing for transfer of energy between ordinary matter and that represented by Λ. In our notation, Bronstein expresses this picture by generalizing Eq. (9) to

$$\dot{\rho}_\Lambda = -\dot{\rho} - 3\frac{\dot{a}}{a}(\rho + p), \tag{36}$$

where ρ and p are the energy density and pressure of ordinary matter and radiation. Bronstein goes on to propose a violation of local energy conservation, a thought that no longer seems interesting. North (1965, p. 81) finds Eddington's (1939) interpretation of the zero point of energy also somewhat hard to defend. But for our purpose the important point is that the idea of Λ as a form of energy has been present, in at least some circles, for many years.

The zero-point energy of fields contributes to the dark-energy density. To make physical sense the sum over the zero-point mode energies must be cut off at a short distance or a high frequency up to which the model under consideration is valid. The integral of the zero-point energy ($k/2$) of normal modes (of wave number k) of a massless real bosonic scalar field (Φ), up to the wave-number cutoff k_c, gives the vacuum energy density the quantum-mechanical expectation value[19]

$$\rho_\Phi = \int_0^{k_c} \frac{4\pi k^2 dk}{(2\pi)^3}\frac{k}{2} = \frac{k_c^4}{16\pi^2}. \tag{37}$$

Nernst (1916) seems to have been the first to write down this equation, in connection with the idea that the zero-point energy of the electromagnetic field fills the vacuum, as a light aether, that could have physically significant properties.[20] This was before Heisenberg and Schrödinger: Nernst's hypothesis is that each degree of freedom, to which classical statistical mechanics assigns energy $kT/2$, has *Nullpunktsenergie* $h\nu/2$. This would mean that the ground-state energy of a one-dimensional harmonic oscillator is $h\nu$, twice the correct value. Nernst's expression for the energy density in the electromagnetic field thus differs from Eq. (37) by a factor of 2 (after taking account of the two polarizations), which is wonderfully close. For a numerical example, Nernst noted that if the cutoff frequency were $\nu = 10^{20}$ Hz, or ~ 0.4 MeV, the energy density of the *Lichtäther* (light aether) would be $10^{23}\,\mathrm{erg\,cm^{-3}}$, or about $100\,\mathrm{g\,cm^{-3}}$.

By the end of the 1920s Nernst's hypothesis was replaced with the demonstration that in quantum mechanics the zero-point energy of the vacuum is as real as any other. W. Pauli, in unpublished work in the 1920s,[21] repeated Nernst's calculation, with the correct factor of 2, taking k_c to correspond to the classical electron radius. Pauli knew the value of ρ_Λ was quite unacceptable: the radius of the static Einstein universe with this value of ρ_Λ "would not even reach to the moon" (Rugh and Zinkernagel, 2002, p. 5).[22] The modern version of this "physicists' cosmological constant problem" is even

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[18]Kragh (1996, p. 36) describes Bronstein's motivation and history. We discuss this model in more detail in Sec. III.E, and comment on why decay of dark energy into ordinary matter or radiation would be hard to reconcile with the thermal spectrum of the 3-K cosmic microwave background radiation. Decay into the dark sector may be interesting.

[19]Equation (37), which usually figures in discussions of the vacuum energy puzzle, gives a helpful indication of the situation: the zero-point energy of each mode is real and the sum is large. The physics is seriously incomplete, however. The elimination of spatial momenta with magnitudes $k > k_c$ only makes sense if there is a preferred reference frame in which k_c is defined. Magueijo and Smolin (2002) mention a related issue: In which reference frame is the Planck momentum of a virtual particle at the threshold for new phenomena? In both cases one may implicitly choose the rest frame for the large-scale distribution of matter and radiation. It seems strange to think that the microphysics is concerned about large-scale structure, but perhaps this happens in a sea of interacting fields. The cutoff in Eq. (37) might be applied at a fixed comoving wave number $k_c \propto a(t)^{-1}$, or at a fixed physical value of k_c. The first prescription can be described by an action written as a sum of terms $\dot{\Phi}_k^2/2 + k^2\Phi_k^2/[2a(t)^2]$ for the allowed modes. The zero-point energy of each mode scales with the expansion of the universe as $a(t)^{-1}$, and the sum over modes scales as $\rho_\Phi \propto a(t)^{-4}$, consistent with $k_c \propto a(t)^{-1}$. In the limit of exact spatial homogeneity, an equivalent approach uses the spatial average of the standard expression for the field stress-energy tensor. Indeed, DeWitt (1975) and Akhmedov (2002) show that the vacuum expectation value of the stress-energy tensor, expressed as an integral cutoff at $k = k_c$, and computed in the preferred coordinate frame, is diagonal with a space part $p_\Phi = \rho_\Phi/3$, for the massless field we are considering. That is, in this prescription the vacuum zero-point energy acts like a homogeneous sea of radiation. This defines a preferred frame of motion, where the stress-energy tensor is diagonal, which is not unexpected because we need a preferred frame to define k_c. It is unacceptable as a model for the properties of dark energy, of course. For example, if the dark-energy density were normalized to the value now under discussion, and varied as $\rho_\Lambda \propto a(t)^{-4}$, it would quite mess up the standard model for the origin of the light elements. We get a more acceptable model for the behavior of ρ_Λ from the second prescription, with the cutoff at a fixed physical momentum. If we also want to satisfy local energy conservation we must take the pressure to be $p_\Phi = -\rho_\Phi$. This does not contradict the derivation of p_Φ in the first prescription, because the second situation cannot be described by an action: the pressure must be stipulated, not derived. What is worse is that the known fields at laboratory momenta certainly do not allow this stipulation; they are well described by analogs of the action in the first prescription. This quite unsatisfactory situation illustrates how far we are from a theory of vacuum energy.

[20]A helpful discussion of Nernst's ideas on cosmology is that of Kragh (1996, pp. 151–157).

[21]This is discussed by Enz and Thellung (1960); Enz (1974); Rugh and Zinkernagel (2002, pp. 4 and 5), and Straumann (2002).

[22]In an unpublished letter in 1930, G. Gamow considered the gravitational consequences of the Dirac sea (Dolgov, 1989, p. 230). We thank A. Dolgov for helpful correspondence on this point.

more acute, because a natural value for k_c is thought to be much larger than what Nernst or Pauli used.[23]

While there was occasional discussion of this issue in the middle of the 20th century (as in the quote from N. Bohr in Rugh and Zinkernagel, 2002, p. 5), the modern era begins with the paper by Zel'dovich (1967) that convinced the community to consider the possible connection between the vacuum energy density of quantum physics and Einstein's cosmological constant.[24]

If the physics of the vacuum looks the same to any inertial observer its contribution to the stress-energy tensor is the same as Einstein's cosmological constant [Eq. (19)]. Lemaître (1934) notes this: "in order that absolute motion, i.e., motion relative to the vacuum, may not be detected, we must associate a pressure $p=-\rho c^2$ to the energy density ρc^2 of vacuum." Gliner (1965) goes further, presenting the relation between the metric tensor and the stress-energy tensor of a vacuum that is the same to any inertial observer. But it was Zel'dovich (1968) who presented the argument clearly enough and at the right time to catch the attention of the community.

With the development of the concept of broken symmetry in the now standard model for particle physics came the idea that the expansion and cooling of the universe is accompanied by a sequence of first-order phase transitions accompanying the symmetry breaking. Each first-order transition has a latent heat that appears as a contribution to an effective time-dependent $\Lambda(t)$ or dark-energy density.[25] The decrease in value of the dark-energy density at each phase transition is much larger than the acceptable present value (within relativistic cosmology); the natural presumption is that the dark energy is negligible now. This final condition seems bizarre, but the picture led to the very influential concept of inflation. We discussed the basic elements in connection with Eq. (27); we now turn to some implications.

[23]In terms of an energy scale ϵ_Λ defined by $\rho_\Lambda=\epsilon_\Lambda^4$, the Planck energy $G^{-1/2}$ is about 30 orders of magnitude larger than the "observed" value of ϵ_Λ. This is, of course, an extreme case, since many of the theories of interest break down well below the Planck scale. Furthermore, in addition to other contributions, one may add a counterterm to Eq. (37) to predict any value of ρ_Λ. With reference to this point, it is interesting to note that while Pauli did not publish his computation of ρ_Λ, he remarks in his famous 1933 *Handbuch der Physik* review on quantum mechanics that it is more consistent to "exclude a zero-point energy for each degree of freedom as this energy, evidently from experience, does not interact with the gravitational field" (Rugh and Zinkernagel, 2002, p. 5). Pauli was fully aware that one must take account of zero-point energies in the binding energies of molecular structure, for example (and we expect he was aware that what contributes to the energy contributes to the gravitational mass). He chose to drop the section with the above comment from the second (1958) edition of the review (Pauli, 1980, pp. iv and v). In a globally supersymmetric field theory there are equal numbers of bosonic and fermionic degrees of freedom, and the net zero-point vacuum energy density ρ_Λ vanishes (Iliopoulos and Zumino, 1974; Zumino, 1975). However, supersymmetry is not a symmetry of low-energy physics, or even at the electroweak unification scale. It must be broken at low energies, and the proper setting for a discussion of the zero-point ρ_Λ in this case is locally supersymmetric supergravity. Weinberg (1989, p. 6) notes "it is very hard to see how any property of supergravity or superstring theory could make the effective cosmological constant sufficiently small." Witten (2001) and Ellwanger (2002) review more recent developments on this issue in the superstring/M theory/branes scenario.

[24]For subsequent more detailed discussions of this issue, see Zel'dovich (1981), Weinberg (1989), Carroll, Press, and Turner (1992), Sahni and Starobinsky (2000), Carroll (2001), and Rugh and Zinkernagel (2000).

C. Inflation

1. The scenario

The deep issue addressed by inflation is the origin of the large-scale homogeneity of the observable universe. In a relativistic model with positive pressure we can see distant galaxies that have not been in causal contact with each other since the singular start of expansion [Sec. II.C, Eq. (26)]; they are said to be outside each other's particle horizon. Why do apparently causally unconnected parts of space look so similar?[26] Kazanas (1980), Guth (1981), and Sato (1981a, 1981b) make the key point: if the early universe were dominated by the energy density of a relatively flat real scalar field (inflaton) potential $V(\Phi)$ that acts like Λ, the particle horizon could spread beyond the universe we can see. This would allow for the possibility that microphysics during inflation could smooth inhomogeneities sufficiently to provide an explanation of the observed large-scale homogeneity. (We are unaware of a definitive demonstration of this idea, however.)

In the inflation scenario the field Φ rolls down its potential until eventually $V(\Phi)$ steepens enough to terminate inflation. Energy in the scalar field is supposed to decay to matter and radiation, heralding the usual Big Bang expansion of the universe. With the modifications of Guth's (1981) scenario by Linde (1982) and Albrecht and Steinhardt (1982), the community quickly accepted this promising and elegant way to understand the origin of our homogeneous expanding universe.[27]

In Guth's (1981) picture the inflation kinetic-energy density is subdominant during inflaton, $\dot\Phi^2 \ll V(\Phi)$, so from Eqs. (30) the pressure p_Φ is very close to the negative of the mass density ρ_Φ, and the expansion of the universe approximates the de Sitter solution, $a \propto \exp(H_\Lambda t)$ [Eq. (27)].

[25]Early references to this point were made by Linde (1974), Dreitlein (1974), Kirzhnitz and Linde (1974), Veltman (1975), Bludman and Ruderman (1977), Canuto and Lee (1977), and Kolb and Wolfram (1980).

[26]Early discussions of this question are reviewed by Rindler (1956), more recent examples are Misner (1969); Dicke and Peebles (1979); and Zee (1980).

[27]Aspects of the present state of the subject are reviewed by Guth (1997); Brandenberger (2001); and Lazarides (2002).

For our comments on the spectrum of mass density fluctuations produced by inflation and the properties of solutions of the dark-energy models in Sec. III.E we shall find it useful to have another scalar field model. Lucchin and Matarrese (1985a, 1985b) considered the potential

$$V(\Phi)=\frac{A}{G^2}\exp[-\Phi\sqrt{8\pi qG}], \tag{38}$$

where q and A are parameters.[28] They showed that the scale factor and the homogeneous part of the scalar field evolve in time as

$$a(t)=a_0[1+Nt]^{2/q},$$

$$\Phi(t)=\frac{1}{\sqrt{2\pi qG}}\ln[1+Nt], \tag{39}$$

where $N=2q\sqrt{\pi A}/\sqrt{G(6-q)}$. If $q<2$ this model inflates. Halliwell (1987) and Ratra and Peebles (1988) showed that the solution (39) of the homogeneous equation of motion has the attractor property[29] mentioned in connection with Eq. (31). This exponential potential is of historical interest: it provided the first clear illustration of an attractor solution. We will return to this point in Sec. III.E.

A signal achievement of inflation is that it offers a theory for the origin of the departures from homogeneity. Inflation tremendously stretches length scales, so that cosmologically significant lengths now correspond to extremely short lengths during inflation. On these tiny length scales quantum mechanics governs: the wavelengths of zero-point field fluctuations generated during inflation are stretched by the inflationary expansion,[30] and these fluctuations are converted to classical density fluctuations in the late time universe.[31]

The power spectrum of the fluctuations depends on the model for inflation. If the expansion rate during inflation is close to exponential [Eq. (27)], the zero-point fluctuations are frozen into primeval mass density fluctuations with the power spectrum

$$P(k)=\langle|\delta(k,t)|^2\rangle=AkT^2(k). \tag{40}$$

Here $\delta(k,t)$ is the Fourier transform at wave number k of the mass density contrast $\delta(\vec{x},t)=\rho(\vec{x},t)/\langle\rho(t)\rangle-1$, where ρ is the mass density and $\langle\rho\rangle$ the mean value. After inflation, but at very large redshifts, the spectrum in this model is $P(k)\propto k$ on all interesting length scales. This means the curvature fluctuations produced by the mass fluctuations diverge only as $\log k$. The form $P(k)\propto k$ thus need not be cut off anywhere near observationally interesting lengths, and in this sense it is scale invariant.[32] The transfer function $T(k)$ accounts for the effects of radiation pressure and the dynamics of nonrelativistic matter on the evolution of $\delta(k,t)$, computed in linear perturbation theory, at redshifts $z\lesssim 10^4$. The constant A is determined by details of the chosen inflation model we need not describe.

The exponential potential model in Eq. (38) produces the power spectrum[33]

$$P(k)=Ak^nT^2(k), \quad n=(2-3q)/(2-q). \tag{41}$$

When $n\neq 1 (q\neq 0)$ the power spectrum is said to be tilted. This offers a parameter n to be adjusted to fit the observations of the large-scale structure, though as we will discuss, the simple scale-invariant case $n=1$ is close to the best fit to the observations.

The mass fluctuations in these inflation models are said to be adiabatic, because they are what results from adiabatically compression or decompression of parts of an exactly homogeneous universe. This means the initial conditions for the mass distribution are described by one function of position, $\delta(\vec{x},t)$. This function is a realization of a spatially stationary random Gaussian process, because it is frozen out of almost free quantum field fluctuations. Thus the single function of position is statistically prescribed by its power spectrum, as in Eqs. (40) and (41). More complicated models for inflation produce density fluctuations that are not Gaussian, or do not have simple power-law spectra, or have parts that break adiabaticity, such as gravitational waves (Rubakov, Sazhin, and Veryaskin, 1982), magnetic fields (Turner and Widrow, 1988; Ratra, 1992b), or new hypothetical fields. All these extra features may be invoked to fit the observations, if needed. It may be significant that none seem to be needed to fit the main cosmological constraints we have now.

2. Inflation in a low-density universe

We do need an adjustment from the simplest case—an Einstein–de Sitter cosmology—to account for the measurements of the mean mass density. In the two models that lead to Eqs. (40) and (41) the enormous expansion factor during inflation suppresses the curvature of space sections, making Ω_{K0} negligibly small. If $\Lambda=0$, this fits the Einstein–de Sitter model [Eq. (35)], which in the

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[28]Similar exponential potentials appear in some higher-dimensional Kaluza-Klein models. For an early discussion see Shafi and Wetterich (1985).

[29]Ratra (1989, 1992a) shows that spatial inhomogeneities do not destroy this property, that is, for $q<2$ the spatially inhomogeneous scalar field perturbation has no growing mode.

[30]The strong curvature of spacetime during inflation makes the vacuum state quite different from that of Minkowski spacetime (Ratra, 1985). This is somewhat analogous to how the Casimir metal plates modify the usual Minkowski spacetime vacuum state.

[31]For the development of these ideas see Hawking (1982); Starobinsky (1982); Guth and Pi (1982); Bardeen, Steinhardt, and Turner (1983); and Fischler, Ratra, and Susskind (1985).

[32]The virtues of a spectrum that is scale invariant in this sense were noted before inflation by Harrison (1970), Peebles and Yu (1970), and Zel'dovich (1972).

[33]This is discussed by Abbott and Wise (1984), Lucchin and Matarrese (1985a, 1985b), and Ratra (1989, 1992a).

absence of data clearly is the elegant choice. But the high mass density in this model was already seriously challenged by the data available in 1983, on the low streaming flow of the nearby galaxies toward the nearest known large mass concentration, in the Virgo cluster of galaxies, and the small relative velocities of galaxies outside the rich clusters of galaxies.[34] A striking and long familiar example of the latter is that the galaxies immediately outside the Local Group of galaxies, at distances of a few megaparsecs, are moving away from us in a good approximation to Hubble's homogeneous flow, despite the very clumpy distribution of galaxies on this scale.[35] The options (within general relativity) are that the mass density is low, so its clumpy distribution has little gravitational effect, or the mass density is high and the mass is more smoothly distributed than the galaxies. We comment on the first option here, and the second in connection with the cold-dark-matter model for structure formation in Sec. III.D.

Under the first option we have these choices: introduce a cosmological constant, or space curvature, or perhaps even both. In the conventional inflation picture space curvature is unacceptable, but there is another line of thought that leads to a universe with open space sections. Gott's (1982) scenario commences with a large energy density in an inflaton at the top of its potential. This behaves like Einstein's cosmological constant and produces a near de Sitter universe expanding as $a \propto \exp(H_\Lambda t)$, with sufficient inflation to allow for a microphysical explanation of the large-scale homogeneity of the observed universe. As the inflaton gradually rolls down the potential it reaches a point where there is a small bump in the potential. The inflaton tunnels through this bump by nucleating a bubble. Symmetry forces the interior of the bubble to have open spatial sections (Coleman and De Luccia, 1980), and the continuing presence of a nonzero $V(\Phi)$ inside the bubble acts like Λ, resulting in an open inflating universe. The potential is supposed to steepen, bringing the second limited epoch of inflation to an end before space curvature has been completely redshifted away. The region inside the open bubble at the end of inflation is a radiation-dominated Friedmann-Lemaître open model, with $0<\Omega_{K0}<1$ [Eq. (16)]. This can fit the dynamical evidence for low Ω_{M0} with $\Lambda=0$.[36]

The decision on which scenario, spatially flat or open, is elegant, if either, depends ultimately on which Nature has chosen, if either.[37] But it is natural to make judgments in advance of the evidence. Since the early 1980s there have been occasional explorations of the open case, but the community generally has favored the flat case, $\Omega_{K0}=0$, without or, more recently, with a cosmological constant, and indeed the evidence now shows that space sections are close to flat. The earlier preference for the Einstein–de Sitter case with $\Omega_{K0}=0$ and $\Omega_{\Lambda 0}=0$ led to considerable interest in the picture of biased galaxy formation in the cold-dark-matter model, as we now describe.

D. The cold-dark-matter model

Some of the present cosmological tests have been understood since the 1930s; others are based on new ideas about structure formation. A decade ago a half dozen models for structure formation were under discussion,[38] now the known viable models have been winnowed to

[34]This is discussed by Davis and Peebles (1983a, 1983b) and Peebles (1986). Relative velocities of galaxies in rich clusters are large, but the masses in clusters are known to add up to a modest mean mass density. Thus most of the Einstein–de Sitter mass would have to be outside the dense parts of the clusters, where the relative velocities are small.

[35]The situation a half century ago is illustrated by the compilation of galaxy redshifts by Humason, Mayall, and Sandage (1956). In this sample of 806 galaxies, 14 have negative redshifts (after correction for the rotation of the Milky Way galaxy and for the motion of the Milky Way toward the other large galaxy in the Local Group, the Andromeda Nebula), indicating motion toward us. Nine are members of the Local Group, at distances $\lesssim 1$ Mpc. Four are in the direction of the Virgo cluster, at redshift $\sim 1200\ \mathrm{km\,s^{-1}}$ and distance ~ 20 Mpc. Subsequent measurements indicate two of these four really have negative redshifts, and plausibly are members of the Virgo cluster on the tail of the distribution of peculiar velocities of the cluster members. (Astronomers use the term *peculiar velocity* to denote the deviation from the uniform Hubble expansion velocity.) The last of the 14, NGC 3077, is in the M 81 group of galaxies at 3-Mpc distance. It is now known to have a small positive redshift.

[36]Gott's scenario is resurrected by Ratra and Peebles (1994, 1995). See Bucher and Turok (1995), Yamamoto, Sasaki, and Tanaka (1995), and Gott (1997) for further discussions of this model. In this case spatial curvature provides a second cosmologically relevant length scale (in addition to that set by the Hubble radius H^{-1}), so there is no natural preference for a power-law power spectrum (Ratra, 1994; Ratra and Peebles, 1995).

[37]At present, high-energy physics considerations do not provide a compelling specific inflation model, but there are strong indications that inflation occurs in a broad range of models, so it might not be unreasonable to think that future advances in high-energy physics could give us a compelling and observationally successful model of inflation that will determine whether the scenario is flat or open.

[38]A scorecard is given in Peebles and Silk (1990). Structure-formation models that assume all matter is baryonic, and those that augment baryons with hot dark matter such as low-mass neutrinos, were already seriously challenged a decade ago. Vittorio and Silk (1985) showed that the Uson and Wilkinson (1984) bound on the small-scale anisotropy of the 3-K cosmic microwave background temperature rules out a baryon-dominated universe with adiabatic initial conditions. This is because the dissipation of the baryon density fluctuations by radiation drag as the primeval plasma combines to neutral hydrogen (at redshift $z\sim 1000$) unacceptably suppresses structure formation on the scale of galaxies. Cold dark matter avoids this problem by eliminating radiation drag. This is one of the reasons attention turned to the hypothetical nonbaryonic cold dark matter. There has not been a thorough search for more baroque initial conditions that might save the baryonic dark-matter model, however.

one class: cold dark matter (CDM) and variants. We comment on the present state of tests of the CDM model in Sec. IV.A.2, and in connection with the cosmological tests in Sec. IV.B.

The CDM model assumes that the mass of the universe now is dominated by dark matter that is nonbaryonic and acts like a gas of massive, weakly interacting particles with negligibly small primeval velocity dispersion. Structure is supposed to have formed as a result of the gravitational growth of primeval departures from homogeneity that are adiabatic, scale invariant, and Gaussian. The early discussions also assume an Einstein–de Sitter universe. These features all are naturally implemented in simple models for inflation, and the CDM model may have been inspired in part by the developing ideas of inflation. But the motivation for writing down this model was to find a simple way to show that the observed present-day mass fluctuations can agree with growing evidence that the anisotropy of the 3-K thermal cosmic microwave background radiation is very small (Peebles, 1982). The first steps toward turning this picture into a model for structure formation were taken by Blumenthal *et al.* (1984).

In the decade commencing about 1985 the standard cosmology for many active in research in this subject was the Einstein–de Sitter model, and for good reason: it eliminates the coincidences problem, it avoids the curiosity of nonzero dark energy, and it fits the condition from conventional inflation that space sections have zero curvature. But unease about the astronomical problems with the high-mass density of the Einstein–de Sitter model led to occasional discussions of a low-density universe with or without a cosmological constant, and the CDM model played an important role in these considerations, as we now discuss.

When the CDM model was introduced it was known that the observations disfavor the high-mass density of the Einstein–de Sitter model, unless the mass is more smoothly distributed than the visible matter (Sec. III.C). Kaiser (1984) and Davis *et al.* (1985) showed that this wanted biased distribution of visible galaxies, relative to the distribution of all of the mass, can follow in a natural way in the CDM theory. In short, where the mass density is high enough to lead to the gravitational assembly of a large galaxy the mass density tends to be high nearby, favoring the formation of neighboring large galaxies.

The biasing concept is important and certainly had to be explored. But in 1985 there was little empirical evidence for the effect and significant arguments against it, mainly involving the empty state of the voids between the concentrations of large galaxies.[39] In the biasing picture the voids contain most of the mass of an Einstein–de Sitter universe, but few of the galaxies, since galaxy formation there is supposed to have been suppressed. But it is hard to see how galaxy formation could be entirely extinguished: the CDM model would be expected to predict a void population of irregular galaxies that show signs of a difficult youth. Many irregular galaxies are observed, but they avoid the voids. The straightforward reading of the observations thus is that the voids are empty, and that the dynamics of the motions of the visible galaxies therefore show that Ω_{M0} is well below unity, and that the mass is not more smoothly distributed than that of the visible galaxies.

In a low-density open universe, with $\Omega_{\Lambda 0}=0$ and positive Ω_{K0}, the growth of mass clustering is suppressed at $z \lesssim \Omega_{M0}^{-1}-1$. Thus to agree with the observed low-redshift mass distribution, density fluctuations at high redshift must be larger in the open model than in the Einstein–de Sitter case. This makes it harder to understand the small 3-K cosmic microwave background anisotropy. In a low-density spatially flat universe with $\Omega_{K0}=0$ and a cosmological constant, the transition from matter-dominated expansion to Λ-dominated expansion is more recent than the transition from matter-dominated expansion to space-curvature-dominated expansion in an open universe with $\Lambda=0$, as one sees from Eq. (11). This makes density fluctuations grow almost as much as those in the Einstein–de Sitter model, thus allowing smaller peculiar velocities in the flat-Λ case, a big help in understanding the observations.[40]

An argument for low Ω_{M0}, with or without Λ, developed out of the characteristic length scale for structure in the CDM model. In the Friedmann-Lemaître cosmology the mass distribution is gravitationally unstable. This simple statement has a profound implication: the early universe has to have been very close to homogeneous, and the growing departures from homogeneity at high redshift are well described by linear perturbation theory. The linear density fluctuations may be decomposed into Fourier components (or generalizations for open- or closed-space sections). At high enough redshift the wavelength of a mode is much longer than the time-dependent Hubble length H^{-1}, and gravitational instability makes the mode amplitude grow. Adiabatic fluctuations remain adiabatic, because different regions behave as if they were parts of different homogeneous universes. When the Hubble length becomes comparable to the mode proper wavelength, the baryons and radiation, strongly coupled by Thomson scattering at high redshift, oscillate as an acoustic wave, and the mode amplitude for the cold dark matter stops growing.[41] The mass densities in dark matter and radia-

[39]The issue is presented by Peebles (1986); the data and history of ideas are reviewed in Peebles (2001).

[40]The demonstration that the suppression of peculiar velocities is a lot stronger than the suppression of the growth of structure is in Peebles (1984) and Lahav *et al.* (1991).

[41]At high redshift the dark-matter mass density is less than that of the radiation. The radiation thus fixes the expansion rate, which is too rapid for the self-gravity of the dark matter to have any effect on its distribution. Early discussions of this effect are those of Guyot and Zel'dovich (1970) and Mészáros (1974).

tion are equal at redshift $z_{\mathrm{eq}}=2.4\times10^4\Omega_{\mathrm{M0}}h^2$; thereafter the dark-matter mass density dominates and the fluctuations in its distribution start to grow again. The suppressed growth of density fluctuations within the Hubble length at $z>z_{\mathrm{eq}}$ produces a bend in the power spectrum of the dark mass distribution, from $P(k)\propto k$ at long wavelengths, if scale invariant [Eq. (40)], to $P(k)\propto k^{-3}$ at short wavelengths.[42] This means that at small scales (large k) the contribution to the variance of the mass density per logarithmic interval of wavelength is constant, and at small k the contribution to the variance of the Newtonian gravitational potential per logarithmic interval of k is constant.

The wavelength at the break in the spectrum is set by the Hubble length at equal radiation and matter mass densities, $t_{\mathrm{eq}}\propto z_{\mathrm{eq}}^{-2}$. This characteristic break scale grows by the factor z_{eq} to $\lambda_{\mathrm{break}}\sim z_{\mathrm{eq}}t_{\mathrm{eq}}$ at the present epoch. The numerical value is [Peebles, 1980a, Eq. (92.47)]

$$\lambda_{\mathrm{break}}\sim50\Omega_{\mathrm{M0}}^{-1}h^{-2}\ \mathrm{Mpc}. \tag{42}$$

If Λ is close to constant, or $\Lambda=0$ and space sections are curved, it does not appreciably affect the expansion rate at redshift z_{eq}, so this characteristic length is the same in flat and open cosmological models. In an Einstein–de Sitter model, with $\Omega_{\mathrm{M0}}=1$, the predicted length scale at the break in the power spectrum of CDM mass fluctuations is uncomfortably small relative to structures such as are observed in clusters of clusters of galaxies (superclusters), and relative to the measured galaxy two-point correlation function. That is, more power is observed on scales ~100 Mpc than is predicted in the CDM model with $\Omega_{\mathrm{M0}}=1$. Since λ_{break} scales as $\Omega_{\mathrm{M0}}^{-1}$, a remedy is to go to a universe with small Ω_{M0}, either with $\Lambda=0$ and open-space sections or $\Omega_{\mathrm{K0}}=0$ and a nonzero cosmological constant. The latter case is now known as ΛCDM.[43]

E. Dark energy

The idea that the universe contains close to homogeneous dark energy that approximates a time-variable cosmological "constant" arose in particle physics, through the discussion of phase transitions in the early universe and through the search for a dynamical cancellation of the vacuum energy density; in cosmology, through the discussions of how to reconcile a cosmologically flat universe with the small mass density indicated by galaxy peculiar velocities; and on both sides by the thought that Λ might be very small now because it has been rolling toward zero for a very long time.[44]

The idea that dark energy decays by emission of matter or radiation is now strongly constrained by the condition that the decay energy must not significantly disturb the spectrum of the 3-K cosmic microwave background radiation. But the history of the idea is interesting, and decay to dark matter still a possibility, so we comment here on both. The picture of dark energy in the form of defects in cosmic fields has not received much attention in recent years, in part because the computations are difficult, but might yet prove to be productive. Much discussed nowadays is dark energy in a slowly varying scalar field. The idea is reviewed at some length here and in even more detail in the Appendix. We begin with another much discussed approach: prescribe the dark energy by parameters in numbers that seem fit for the quality of the measurements.

1. The XCDM parametrization

In the XCDM parametrization the dark energy interacts only with itself and gravity, the dark-energy density $\rho_{\mathrm{X}}(t)>0$ is approximated as a function of world time alone, and the pressure is written as

$$p_{\mathrm{X}}=w_{\mathrm{X}}\rho_{\mathrm{X}}, \tag{43}$$

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[42]That is, the transfer function in Eq. (40) goes from a constant at small k to $T^2(k)\propto k^{-4}$ at large k.

[43]The scaling of λ_{break} with Ω_{M0} was frequently noted in the 1980s. The earliest discussions of the significance for large-scale structure are in Silk and Vittorio (1987) and Efstathiou, Sutherland, and Maddox (1990), who consider a spatially flat universe; Blumenthal, Dekel, and Primack (1988), who consider the open case; and Holtzman (1989), who considers both. For the development of tests of the open case see Lyth and Stewart (1990); Ratra and Peebles (1994); Kamionkowski *et al.* (1994a); and Górski *et al.* (1995). Pioneering steps in the analysis of the anisotropy of the 3-K cosmic microwave background temperature in the ΛCDM model include the work of Kofman and Starobinsky (1985) and Górski, Silk, and Vittorio (1992). We review developments after the Cosmic Background Explorer (COBE) detection of the anisotropy (Smoot *et al.*, 1992) in Secs. IV.B.11 and IV.B.12.

[44]This last idea is similar in spirit to Dirac's (1937, 1938) attempt to explain the large dimensionless numbers of physics. He noted that the gravitational force between two protons is much smaller than the electromagnetic force, and that that might be because the gravitational constant G decreases in inverse proportion to world time. This is the earliest discussion we know of what has come to be called the *hierarchy problem*, that is, the search for a mechanism that might be responsible for the large ratio between a possibly more fundamental high-energy scale, for example, that of grand unification or the Planck scale (where quantum-gravitational effects become significant) and a lower possibly less fundamental energy scale, for example, that of electroweak unification (see, for example, Georgi, Quinn, and Weinberg, 1974). The hierarchy problem in particle physics may be rephrased as a search for a mechanism to prevent the light electroweak symmetry-breaking Higgs scalar field mass from becoming large because of a quadratically divergent quantum-mechanical correction (see, for example, Susskind, 1979). In this sense it is similar in spirit to the physicists' cosmological constant problem of Sec. III.B.

an expression that has come to be known as the cosmic equation of state.[45] Then the local energy-conservation equation (9) is

$$\frac{d\rho_X}{dt} = -3\,\frac{\dot{a}}{a}\,\rho_X(1+w_X). \tag{44}$$

If w_X is constant the dark-energy density scales with the expansion factor as

$$\rho_X \propto a^{-3(1+w_X)}. \tag{45}$$

If $w_X < -1/3$ the dark energy makes a positive contribution to $\ddot{a}/a$ [Eq. (8)]. If $w_X = -1/3$ the dark energy has no effect on $\ddot{a}$, and the energy density varies as $\rho_X \propto 1/a^2$, the same as the space-curvature term in $\dot{a}^2/a^2$ [Eq. (11)]. That is, the expansion time histories are the same in an open model with no dark energy and in a spatially flat model with $w_X = -1/3$, although the space-time geometries differ.[46] If $w_X < -1$ the dark-energy density increases.[47]

Equation (45) with constant w_X has the great advantage of simplicity. An appropriate generalization for the more precise measurements to come might be guided by the idea that the dark-energy density is close to homogeneous, spatial variations rearranging themselves at close to the speed of light, as in the scalar field models discussed below. Then for most of the cosmological tests we have an adequate general description of the dark energy if we let w_X be a free function of time.[48] In scalar field pictures w_X is derived from the field model; it can be a complicated function of time even when the potential is a simple function of the scalar field.

The analysis of the large-scale anisotropy of the 3-K cosmic microwave background radiation requires a prescription for how the spatial distribution of the dark energy is gravitationally related to the inhomogeneous distribution of other matter and radiation (Caldwell *et al.*, 1998). In XCDM this requires at least one more parameter, an effective speed of sound, with $c_{sX}^2 > 0$ (for stability, as discussed in Sec. II.B), in addition to w_X.

2. Decay by emission of matter or radiation

Bronstein (1933) introduced the idea that the dark-energy density ρ_Λ decays as a result of the emission of matter or radiation. The continuing discussions of this and the associated idea of decaying dark matter (Sciama, 2001, and references therein) are testimony to the appeal. Considerations in the decay of dark energy include the effect on the formation of light elements at $z \sim 10^{10}$, the contribution to the γ-ray or optical extragalactic background radiation, and the perturbation to the spectrum of the 3-K cosmic microwave background radiation.[49]

The effect on the 3-K cosmic microwave background was of particular interest a decade ago, as a possible explanation of indications of a significant departure from a Planck spectrum. Precision measurements now show the spectrum is very close to thermal. The measurements and their interpretation are discussed by Fixsen *et al.* (1996). They showed that the allowed addition to the 3-K cosmic microwave background energy density ρ_R is limited to just $\delta\rho_R/\rho_R \lesssim 10^{-4}$ since redshift $z \sim 10^5$, when the interaction between matter and radiation was last strong enough for thermal relaxation. The bound on $\delta\rho_R/\rho_R$ is not inconsistent with what the galaxies are thought to produce, but it is well below an observationally interesting dark-energy density.

Dark energy could decay by emission of dark matter, cold or hot, without disturbing the spectrum of the 3-K cosmic microwave background radiation. For example, let us suppose the dark-energy equation of state is $w_X = -1$, and hypothetical microphysics causes the dark-energy density to decay as $\rho_\Lambda \propto a^{-n}$ by the production of

[45] Other parametrizations of dark energy are discussed by Hu (1998) and Bucher and Spergel (1999). The name, XCDM, for the case $w_X < 0$ in Eq. (43) was introduced by Turner and White (1997). There is a long history in cosmology of applications of such an equation of state, and the related evolution of ρ_Λ; examples are Canuto *et al.* (1977); Lau (1985), Huang (1985), Fry (1985), Hiscock (1986), Özer and Taha (1986), and Olson and Jordan (1987). See Ratra and Peebles (1988) for references to other early work on a time-variable Λ and Overduin and Cooperstock (1998) and Sahni and Starobinsky (2000) for reviews. More recent discussions of this and related models may be found in John and Joseph (2000), Zimdahl *et al.* (2001), Dalal *et al.* (2001), Gudmundsson and Björnsson (2002), Bean and Melchiorri (2002), Mak, Belinchón, and Harko (2002), and Kujat *et al.* (2002), through which other recent work may be traced.

[46] As discussed in Sec. IV, it appears difficult to reconcile the case $w_X = -1/3$ with type-Ia supernova apparent magnitude data (Garnavich *et al.*, 1998; Perlmutter *et al.*, 1999a).

[47] This is quite a step from the thought that the dark-energy density is small because it has been rolling to zero for a long time, but the case has found a context (Caldwell, 2002; Maor *et al.*, 2002). Such models were first discussed in the context of inflation (e.g., Lucchin and Matarrese, 1985b), where it was shown that the $w_X < -1$ component could be modeled as a scalar field with a negative kinetic-energy density (Peebles, 1989a).

[48] The availability of a free function greatly complicates the search for tests as opposed to curve fitting. This is clearly illustrated by Maor *et al.* (2002). For more examples see Perlmutter, Turner, and White (1999b) and Efstathiou (1999).

[49] These considerations generally are phenomenological: the evolution of the dark-energy density, and its related coupling to matter or radiation, is assigned rather than derived from an action principle. Recent discussions include those of Pollock (1980); Kazanas (1980); Freese *et al.* (1987); Gasperini (1987); Sato, Terasawa, and Yokoyama (1989); Bartlett and Silk (1990); Overduin, Wesson, and Bowyer (1993); Matyjasek (1995); and Birkel and Sarkar (1997).

nonrelativistic dark matter. Then Bronstein's Eq. (36) shows that dark-matter density varies with time as

$$\rho_M(t) = \frac{A}{a(t)^3} + \frac{n}{3-n}\rho_\Lambda(t), \tag{46}$$

where A is a constant and $0<n<3$. In the late time limit the dark-matter density is a fixed fraction of the dark energy. But for the standard interpretation of the measured anisotropy of the 3-K background we would have to suppose the first term on the right-hand side of Eq. (46) is not much smaller than the second, so the coincidences issue discussed in Sec. III.B.2 is not much relieved. It does not help relieve the problem with the small present value of ρ_Λ [to be discussed in connection with Eq. (47)].

We are not aware of any work on this decaying dark-energy picture. Attention instead has turned to the idea that dark-energy density evolves without emission, as illustrated in Eq. (45) and by the two classes of physical models to be discussed next.

3. Cosmic field defects

The physics and cosmology of topological defects produced at phase transitions in the early universe are reviewed by Vilenkin and Shellard (1994). An example of dark energy is a tangled web of cosmic string, with a fixed mass per unit length, which self-intersects without reconnection. In Vilenkin's (1984) analysis[50] the mean mass density in strings scales as $\rho_{\rm string} \propto [ta(t)]^{-1}$. When ordinary matter is the dominant contribution to $\dot{a}^2/a^2$, the ratio of mass densities is $\rho_{\rm string}/\rho \propto t^{1/3}$. Thus at late times the string mass dominates. In this limit, $\rho_{\rm string} \propto a^{-2}$, $w_X = -1/3$ for the XCDM parametrization of Eq. (45), and the universe expands as $a \propto t$. Davis (1987) and Kamionkowski and Toumbas (1996) proposed the same behavior for a texture model. One can also imagine that domain walls fill space densely enough not to be dangerous. If the domain walls are fixed in comoving coordinates the domain-wall energy density scales as $\rho_X \propto a^{-1}$ (Zel'dovich, Kobzarev, and Okun, 1974; Battye, Bucher, and Spergel, 1999). The corresponding equation of state parameter is $w_X = -2/3$, which is thought to be easier to reconcile with the supernova measurements than $w_X = -1/3$ (Garnavich *et al.*, 1998; Perlmutter *et al.*, 1999a). The cosmological tests of defect models for the dark energy have not been very thoroughly explored, at least in part because an accurate treatment of the behavior of the dark energy is difficult (as seen, for example, in Spergel and Pen, 1997; Friedland, Muruyama, and Perelstein, 2003), but this class of models is worth bearing in mind.

4. Dark-energy scalar field

At present the popular picture for dark energy is a classical scalar field with a self-interaction potential $V(\Phi)$ that is shallow enough that the field energy density decreases with the expansion of the universe at a slower rate than the energy density in matter. This idea grew in part out of the inflation scenario, in part from ideas from particle physics. Early examples are Weiss (1987) and Wetterich (1988).[51] The former considers a quadratic potential with an ultralight effective mass, an idea that reappears in Frieman *et al.* (1995). The latter considers the time variation of the dark-energy density in the case of the Lucchin and Matarrese (1985a) exponential self-interaction potential [Eq. (38)].[52]

In the exponential potential model the scalar field energy density varies with time in constant proportion to the dominant energy density. The evidence is that radiation dominates at redshifts in the range $10^3 \lesssim z \lesssim 10^{10}$, from the success of the standard model for light element formation, and matter dominates at $1 \lesssim z \lesssim 10^3$, from the success of the standard model for the gravitational growth of structure. This would leave the dark energy subdominant today, contrary to what is wanted. This led to the proposal of the inverse-power-law potential in Eq. (31) for a single real scalar field.[53]

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[50]The string flops at speeds comparable to light, making the coherence length comparable to the expansion time t. Suppose a string randomly walks across a region of physical size $a(t)R$ in N steps, where $aR \sim N^{1/2}t$. The total length of this string within the region R is $l \sim Nt$. Thus the mean mass density of the string scales with time as $\rho_{\rm string} \propto l/a^3 \propto [ta(t)]^{-1}$. One randomly walking string does not fill space, but we can imagine many randomly placed strings produce a nearly smooth mass distribution. Spergel and Pen (1997) compute the 3-K cosmic microwave background radiation anisotropy in a related model, where the string network is fixed in comoving coordinates so the mean mass density scales as $\rho_{\rm string} \propto a^{-2}$.

[51]Other early examples include those cited in Ratra and Peebles (1988), as well as Endō and Fukui (1977), Fujii (1982), Dolgov (1983), Nilles (1985), Zee (1985), Wilczek (1985), Bertolami (1986), Ford (1987), Singh and Padmanabhan (1988), and Barr and Hochberg (1988).

[52]For recent discussions of this model see Ferreira and Joyce (1998), Ott (2001), Hwang and Noh (2001), and references therein.

[53]In what follows we focus on this model, which was proposed by Peebles and Ratra (1988). The model assumes a conventionally normalized scalar field kinetic energy and spatial gradient term in the action, and it assumes that the scalar field is coupled only to itself and gravity. The model is then completely characterized by the form of the potential (in addition to all the other usual cosmological parameters, including initial conditions). Models based on other forms for $V(\Phi)$, with a more general kinetic energy and spatial gradient term, or with more general couplings to gravity and other fields, are discussed in the Appendix.

We do not want the hypothetical field Φ to couple too strongly to baryonic matter and fields, because that would produce a "fifth force" that is not observed.[54,55] Within quantum field theory the inverse-power-law scalar field potential renders the model nonrenormalizable and thus pathological. But the model is meant to describe what might emerge out of a more fundamental quantum theory, which perhaps also resolves the physicists' cosmological constant problem (Sec. III.B), as the effective classical description of the dark energy.[56] The potential of this classical effective field is chosen *ad hoc*, to fit the scenario. But one can adduce analogs within supergravity, superstring/M, and brane theory, as reviewed in the Appendix.

The solution for the mass fraction in dark energy in the inverse-power-law potential model [in Eq. (33) when $\rho_\Phi \ll \rho$, and the numerical solution at lower redshifts] is not unique, but it behaves as what has come to be termed an attractor or tracker: it is the asymptotic solution for a broad range of initial conditions.[57] The solution also has the property that ρ_Φ is decreasing, but less rapidly than the mass densities in matter and radiation. This may help alleviate two troubling aspects of the cosmological constant. The coincidences issue is discussed in Sec. III.B. The other is the characteristic energy scale set by the value of Λ,

$$\epsilon_\Lambda(t_0) = \rho_\Lambda(t_0)^{1/4} = 0.003(1-\Omega_{M0})^{1/4} h^{1/2} \text{ eV}, \tag{47}$$

when Ω_{R0} and Ω_{K0} may be neglected. In the limit of constant dark-energy density, cosmology seems to indicate new physics at an energy scale more typical of chemistry. If ρ_Λ is rolling toward zero the energy scale might look more reasonable, as follows (Peebles and Ratra, 1988; Steinhardt *et al.*, 1999; Brax *et al.*, 2000).

Suppose that as conventional inflation ends the scalar field potential switches over to the inverse-power-law form in Eq. (31). Let the energy scale at the end of inflation be $\epsilon(t_I) = \rho(t_I)^{1/4}$, where $\rho(t_I)$ is the energy density in matter and radiation at the end of inflation, and let $\epsilon_\Lambda(t_I)$ be the energy scale of the dark energy at the end of inflation. Since the present value $\epsilon_\Lambda(t_0)$ of the dark-energy scale [Eq. (47)] is comparable to the present energy scale belonging to the matter, we have from Eq. (33)

$$\epsilon_\Lambda(t_I) \simeq \epsilon(t_I) \left[\frac{\epsilon_\Lambda(t_0)}{\epsilon(t_I)} \right]^{2/(\alpha+2)}. \tag{48}$$

For parameters of common inflation models, $\epsilon(t_I) \sim 10^{13}$ GeV and $\epsilon_\Lambda(t_0)/\epsilon(t_I) \sim 10^{-25}$. If, say, $\alpha = 6$, then

$$\epsilon_\Lambda(t_I) \sim 10^{-6} \epsilon(t_I) \sim 10^{7} \text{ GeV}. \tag{49}$$

As this example illustrates, one can arrange the scalar field model so it has a characteristic energy scale that exceeds the energy $\sim 10^3$ GeV below which physics is thought to be well understood: in this model cosmology does not force upon us the idea that there is as yet undiscovered physics at the very small energy in Eq. (47). Of course, where the factor $\sim 10^{-6}$ in Eq. (49) comes from is still an open question, but, as discussed in the Appendix, perhaps it is easier to resolve than the origin of the factor $\sim 10^{-25}$ in the constant Λ case.

When we can describe the dynamics of the departure from a spatially homogeneous field in linear perturba-

[54]The current value of the mass associated with spatial inhomogeneities in the field is $m_\phi(t_0) \sim H_0 \sim 10^{-33}$ eV, as one would expect from the dimensions. More explicitly, one arrives at this mass by writing the field as $\Phi(t,\vec{x}) = \langle \Phi \rangle(t) + \phi(t,\vec{x})$ and Taylor expanding the scalar field potential-energy density $V(\Phi)$ about the homogeneous mean background $\langle \Phi \rangle$ to quadratic order in the spatially inhomogeneous part ϕ, to get $m_\phi^2 = V''(\langle \Phi \rangle)$. Within the context of the inverse-power-law model, the tiny value of the mass follows from the requirements that V vary slowly with the field value and that the current value of V be observationally acceptable. The difference between the roles of m_ϕ and the constant m_q in the quadratic potential model $V = m_q^2 \Phi^2/2$ is worth noting. The mass m_q has an assigned and arguably fine-tuned value. The effective mass $m_\phi \sim H$ belonging to $V \propto \Phi^{-\alpha}$ is a derived quantity that evolves as the universe expands. The small value of $m_\phi(t_0)$ explains why the scalar field energy cannot be concentrated with the nonrelativistic mass in galaxies and clusters of galaxies. Because of the tiny mass a scalar field would mediate a new long-range fifth force if it were not weakly coupled to ordinary matter. Weak coupling also ensures that the contributions to coupling constants (such as the gravitational constant) from the exchange of dark-energy bosons are small, so such coupling constants are not significantly time variable in this model. See, for example, Carroll (1998); Chiba (1999); Horvat (1999); Amendola (2000); Bartolo and Pietroni (2000); and Fujii (2000), for recent discussions of this and related issues.

[55]Coupling between dark energy and dark matter is not constrained by conventional fifth-force measurements. An example is discussed by Amendola and Tocchini-Valentini (2002). Perhaps the first consideration is that the fifth-force interaction between neighboring dark-matter halos must not be so strong as to shift regular galaxies of stars away from the centers of their dark-matter halos.

[56]Of course, the zero-point energy of the quantum-mechanical fluctuations around the mean-field value contributes to the physicists' cosmological constant problem, and renormalization of the potential could destroy the attractor solution (however, see Doran and Jäckel, 2002) and generate couplings between the scalar field and other fields leading to an observationally inconsistent "fifth force." The problems within quantum field theory related to the idea that the energy of a classical scalar field is the dark energy, or drives inflation, are further discussed in the Appendix. The best we can hope for is that the effective classical model is a useful approximation to what is actually happening, which might lead us to a more satisfactory theory.

[57]A recent discussion is that of Brax and Martin (2000). Brax, Martin, and Riazuelo (2000) present a thorough analysis of the evolution of spatial inhomogeneities in the inverse-power-law scalar field potential model and confirm that these inhomogeneities do not destroy the homogeneous attractor solution. For other recent discussions of attractor solutions in a variety of contexts see Liddle and Scherrer (1999); Uzan (1999); de Ritis *et al.* (2000); Holden and Wands (2000); Baccigalupi, Matarrese, and Perrotta (2000); and Huey and Tavakol (2002).

tion theory, a scalar field model is generally characterized by the time-dependent values of w_X [Eq. (43)] and the speed of sound c_{sX} (e.g., Ratra, 1991; Caldwell *et al.*, 1998). In the inverse-power-law potential model the relation between the power-law index α and the equation of state parameter in the matter-dominated epoch is independent of time (Ratra and Quillen, 1992),

$$w_X = -\frac{2}{\alpha+2}. \tag{50}$$

When the dark-energy density starts to make an appreciable contribution to the expansion rate the parameter w_X starts to evolve. The use of a constant value of w_X to characterize the inverse-power-law potential model thus can be misleading. For example, Podariu and Ratra (2000, Fig. 2) show that, when applied to the type-Ia supernova measurements, the XCDM parametrization in Eq. (50) leads to a significantly tighter apparent upper limit on w_X, at fixed Ω_{M0}, than is warranted by the results of a computation of the evolution of the dark-energy density in this scalar field model. Caldwell *et al.* (1998) deal with the relation between scalar field models and the XCDM parametrization by fixing w_X, as a constant or some function of redshift, deducing the scalar field potential $V(\Phi)$ that produces this w_X, and then computing the gravitational response of Φ to the large-scale mass distribution.

IV. THE COSMOLOGICAL TESTS

Our intention is to supplement recent discussions of parameter determinations within standard relativistic cosmology[58] with a broader consideration of the issues summarized in two questions: what is the purpose of the cosmological tests, and how well is the purpose addressed by recent advances and work in progress?

The short answer to the first question used to be that we seek to check the underlying physical theory, general relativity, applied on the time and length scales of cosmology; the model for the stress-energy tensor in Einstein's field equation, suitably averaged over the rich small-scale structure we cannot describe in any detail; and the boundary condition, that the universe we can observe is close to homogeneous and isotropic on the scale of the Hubble length. Recent advances make use of the CDM prescription for the stress-energy tensor and the boundary condition, so we must add the elements of the CDM model to the physics to be checked.

The short answer to the second question is that we now have checks of the standard cosmology, which the model passes. But we believe it takes nothing away from the remarkable advances of the tests, and the exemplary care in the measurements, to note that there is a lot of room for systematic errors. As we discussed in Sec. I.A, the empirical basis for the standard model for cosmology is not nearly as substantial as is the empirical basis for the standard model for particle physics: in cosmology it is not yet a matter of measuring parameters in a well-established physical theory.

We comment on the two main components of physics, general relativity and the CDM model, in Secs. IV.A.1 and IV.A.2. In Sec. IV.B we discuss the state of 13 cosmological tests, proceeding roughly in order of increasing model dependence. We conclude that there is a well-established scientific case for the physical significance of the matter density parameter, and for the result of the measurements, $0.15 \lesssim \Omega_{M0} \lesssim 0.4$ (in the sense of a two standard-deviation range). Our reasoning is summarized in Sec. IV.C, along with an explanation of why we are not so sure about the detection of Λ or dark energy.

A. The theories

1. General relativity

Some early discussions of cosmological tests, as in Robertson (1955) and Bondi (1960), make the point that observationally important elements of spatially homogeneous cosmology follow by symmetry, independent of general relativity. This means some empirical successes of cosmology are not tests of relativity. The point was important in the 1950s, because Steady State theory was a viable alternative to Friedmann-Lemaître cosmology, and because the experimental tests of relativity were quite limited.

The tests of general relativity are much better now, but cosmology still is a considerable extrapolation. The length scale characteristic of the precision tests of general relativity in the Solar System and binary pulsar is $\lesssim 10^{13}$ cm. An important scale for cosmology is the Hubble length, $H_0^{-1} \sim 5000$ Mpc $\sim 10^{28}$ cm, 15 orders of magnitude larger. An extrapolation of 15 orders of magnitude in energy from that achieved at the largest accelerators, $\sim 10^{12}$ eV, brings us to the very different world of Planck energy. Why is the community not concerned about an extrapolation of similar size in the opposite direction? One reason is that the known open issues of physics have to do with small length scales; there is no credible reason to think general relativity may fail on large scales. This is comforting, to be sure, but, as indicated in footnote 7, not the same as a demonstration that we really know the physics of cosmology. Another reason is that if the physics of cosmology were very different from general relativity it surely would have already been manifest in serious problems with cosmological tests. This is also encouraging, but we have to consider details, as follows.

One sobering detail is that in standard cosmology the two dominant contributions to the stress-energy tensor—dark energy and dark matter—are hypothetical, introduced to make the theories fit the observations [Eq. (2)]. This need not mean there is anything wrong with general relativity—we have no reason to expect Nature to have made all matter readily observable other than by its gravity—but it is a cautionary example of the chal-

[58]See Bahcall *et al.* (1999), Schindler (2001), Sarkar (2002), Freedman (2002), Plionis (2002), and references therein.

lenges. Milgrom's (1983) modified Newtonian dynamics (MOND) replaces the dark-matter hypothesis with a hypothetical modification of the gravitational force law. MOND gives remarkably successful fits to observed motions within galaxies, without dark matter (de Blok *et al.*, 2001). So why should we believe that there really is cosmologically significant mass in nonbaryonic dark matter? Unless we are lucky enough to get a laboratory detection, the demonstration must be through the tests of relativistic cosmology (and any other viable cosmological models that may come along, perhaps including an extension of MOND). This indirect chain of evidence for dark matter is becoming tight. A new example—the prospect for a test of the inverse-square law for gravity on the length scales of cosmology—is striking enough for special mention here.[59]

Consider the equation of motion[60] of a freely moving test particle with nonrelativistic peculiar velocity $\vec{v}$ in a universe with expansion factor $a(t)$,

$$\frac{\partial \vec{v}}{\partial t}+\frac{\dot{a}}{a}\vec{v}=\vec{g}=-\frac{1}{a}\vec{\nabla}\varphi. \tag{51}$$

The particle is always moving toward receding observers, which produces the second term in the left-most expression. The peculiar gravitational acceleration $\vec{g}$ relative to the homogeneous background model is computed from the Poisson equation for the gravitational potential φ,

$$\nabla^2\varphi=4\pi G a^2[\rho(\vec{x},t)-\langle\rho\rangle]. \tag{52}$$

The mean mass density $\langle\rho\rangle$ is subtracted because $\vec{g}$ is computed relative to the homogeneous model. The equation of mass conservation expressed in terms of the density contrast $\delta=\rho/\langle\rho\rangle-1$ of the mass distribution modeled as a continuous pressureless fluid is

$$\frac{\partial \delta}{\partial t}+\frac{1}{a}\vec{\nabla}\cdot(1+\delta)\vec{v}=0. \tag{53}$$

In linear perturbation theory in $\vec{v}$ and δ these equations give

$$\frac{\partial^2\delta}{\partial t^2}+2\frac{\dot{a}}{a}\frac{\partial\delta}{\partial t}=4\pi G\langle\rho\rangle\delta, \quad \delta(\vec{x},t)=f(\vec{x})D(t). \tag{54}$$

Here $D(t)$ is the growing solution to the first equation.[61] The velocity field belonging to the solution $D(t)$ is the inhomogeneous solution to Eq. (53) in linear perturbation theory,

$$\vec{v}=\frac{fH_0 a}{4\pi}\int\frac{\vec{y}-\vec{x}}{|\vec{y}-\vec{x}|^3}\delta(\vec{y})\,d^3y, \quad f\simeq\Omega_{M0}^{0.6}. \tag{55}$$

The factor $f=d\ln D/d\ln a$ depends on the cosmological model; the second equation is a good approximation if $\Lambda=0$ or space curvature vanishes.[62] One sees from Eq. (55) that the peculiar velocity is proportional to the gravitational acceleration, as one would expect in linear theory.

The key point of Eq. (54) for the present purpose is that the evolution of the density contrast δ at a given position is not affected by the value of δ anywhere else. This is a consequence of the inverse-square law. The mass fluctuation in a chosen volume element produces a peculiar gravitational acceleration $\delta\vec{g}$ that produces a peculiar velocity field $\delta\vec{v}\propto\vec{g}$. This field has zero divergence and so the mass inside the volume element does not affect the exterior.

For a "toy" model of the effect of an inverse-square law failure, suppose we adjust the expression for the peculiar gravitational acceleration produced by a given mass distribution to

$$\vec{g}=a^3R\int d^3y\;\delta(\vec{y})\frac{\vec{y}-\vec{x}}{|\vec{y}-\vec{x}|}Q(a|\vec{y}-\vec{x}|), \tag{56}$$

where R is some function of world time only. In standard gravity physics $Q(w)=w^{-2}$. We have no basis in fundamental physics for any other function of w. Although Milgrom's (1983) MOND provides a motivation, Eq. (56) is not meant to be an extension of MOND to large-scale flows. It is an *ad hoc* model that illustrates an important property of the inverse-square law.

We noted that in linear theory $\vec{v}\propto\vec{g}$. Thus we find that the divergence of Eq. (56), with the mass conservation Eq. (53) in linear perturbation theory, gives

$$\frac{\partial^2\delta_{\vec{k}}}{\partial t^2}+2\frac{\dot{a}}{a}\frac{\partial\delta_{\vec{k}}}{\partial t}=S(k,t)\delta_{\vec{k}},$$

$$S(k,t)=(4\pi Rk/a)\int_0^\infty w^2dw\,Q(w)j_1(kw/a), \tag{57}$$

where $\delta_{\vec{k}}(t)$ is the Fourier transform of the mass density contrast $\delta(\vec{x},t)$ and j_1 is a spherical Bessel function. The inverse-square law, $Q=w^{-2}$, makes the factor S independent of the wave number k. This means all Fourier amplitudes grow by the same factor in linear perturbation theory (when the growing mode dominates), so the functional form of $\delta(\vec{x},t)$ is preserved and the amplitude grows, as Eq. (54) shows. When Q is some other function, the phases of the $\delta_{\vec{k}}$ are preserved but the functional form of the power spectrum $|\delta_{\vec{k}}|^2$ evolves. For example, if $Q\propto w^{n-2}$ with $-2<n<1$ [so the integral in Eq. (57) does not diverge] Eq. (57) becomes

[59]Binétruy and Silk (2001) and Uzan and Bernardeau (2001) pioneered this probe of the inverse-square law. Related probes, based on the relativistic dynamics of gravitational lensing and the anisotropy of the 3-K thermal background, are discussed by these authors and White and Kochanek (2001).

[60]These relations are discussed in many books on cosmology, including Peebles (1980a).

[61]The general solution is a sum of the growing and decaying solutions, but because the universe has expanded by a large factor since nongravitational forces were last important on large scales we can ignore the decaying part.

[62]This is illustrated in Fig. 13.14 in Peebles (1993). An analytic expression for spherical symmetry is derived by Lightman and Schechter (1990).

$$\frac{\partial^2 \delta_{\vec{k}}}{\partial t^2}+2\frac{\dot{a}}{a}\frac{\partial \delta_{\vec{k}}}{\partial t}=U(t)\left(\frac{a}{k}\right)^n \delta_{\vec{k}}, \tag{58}$$

where U is some function of world time.

If $n>0$, density fluctuations grow more rapidly on larger scales. If $Q(w)$ follows Newtonian gravity on the scale of galaxies and bends to $n>0$ on larger scales it reduces the mean mass density needed to account for the measured large-scale galaxy flows, and perhaps reduces the need for dark matter. But there are testable consequences: the apparent value of Ω_{M0} would vary with the length scale of the measurement, and the form of the power spectrum of the present mass distribution would not agree with the form at redshift $z\sim1000$ when it produced the observed angular power spectrum of the 3-K cosmic microwave background. Thus we are very interested in evidence of consistency of these tests (as discussed in Sec. IV.B.13).

2. The cold-dark-matter model for structure formation

Important cosmological tests assume the CDM model for structure formation (Sec. III.C), so we must consider tests of the model. The model has proven to be a useful basis for analyses of the physics of formation of galaxies and clusters of galaxies (e.g., Colberg *et al.*, 2000; Kay *et al.*, 2002, and references therein). There are issues to consider, however; Sellwood and Kosowsky (2001) provide a useful survey of the situation. We remark on recent developments and what seem to us to be critical issues.

Numerical simulations of the dark mass distribution in the CDM model predict that massive halos have many low-mass satellites, perhaps significantly more than the number observed around the Milky Way galaxy (Klypin *et al.*, 1999; Moore *et al.*, 1999a). The issue is of great interest but not yet a critical test, because it is difficult to predict the nature of star formation in a low-mass dark halo: what does a dark halo look like when star formation or the neutral gas content makes it visible? For recent discussions see Stoehr *et al.* (2002) and Tully *et al.* (2002).

The nature of the dark mass distribution within galaxies is a critical issue, because we know where to look for a distinctive CDM feature: a cusplike central mass distribution, the density varying with radius r as $\rho\propto r^{-\alpha}$ with $\alpha\gtrsim1$. The power law is not unexpected, because there is nothing in the CDM model to set an astronomically interesting value for a core radius.[63] A measure of the mass distribution in disk galaxies is the rotation curve: the circular velocity as a function of radius for matter supported by rotation. In some galaxies with low surface brightness the observed rotation curves are close to a solid body, $v_c\propto r$, near the center, consistent with a near homogeneous core, and inconsistent with the cusplike CDM mass distribution.[64]

The circular velocity produced by the mass distribution $\rho\propto r^{-1}$ is not very different from a solid body, or from the observations, and the difference might be erased by gravitational rearrangement of the dark mass by the fluctuations in the distribution of baryonic mass driven by star formation, winds, or supernovae. This is too complicated to assess by current numerical simulations. But we do have a phenomenological hint: central solid-body rotation is most clearly seen in the disklike galaxies with the lowest surface brightnesses, the objects in which the baryon mass seems least likely to have had a significant gravitational effect on the dark mass. This challenge to the CDM model is pressing.

The challenge may be resolved in a warm-dark-matter model, where the particles are assigned a primeval velocity dispersion that suppresses the initial power spectrum of density fluctuations on small scales (Moore *et al.*, 1999b; Sommer-Larsen and Dolgov, 2001; Bode, Ostriker, and Turok, 2001). But it seems to be difficult to reconcile the wanted suppression of small-scale power with the observation of small-scale clustering in the Lyman-α forest—the neutral hydrogen observed at $z\sim3$ in the Lyman-α resonance absorption lines in quasar spectra (Narayanan *et al.*, 2000; Knebe *et al.*, 2002). Spergel and Steinhardt (2000) point out that the scattering cross section of self-interacting cold-dark-matter particles can be adjusted to suppress the cusplike core.[65] Davé *et al.* (2001) demonstrate the effect in numerical simulations. But Miralda-Escudé (2002) points out that the collisions would tend to make the velocity distribution isotropic, contrary to the evidence for ellipsoidal distributions of dark matter in clusters of galaxies. For recent surveys of the very active debate on these issues see Primack (2002) and Tasitsiomi (2002); for references to still other possible solutions see Davé *et al.* (2001).

Another critical issue traces back to the biasing picture discussed in Sec. III.D. If Ω_{M0} is well below unity there need not be significant mass in the voids defined by the large galaxies. But the biasing process still operates, and might be expected to cause dwarf or irregular galaxies to trespass into the voids outlined by the large regular galaxies. This seems to happen in CDM model simulations to a greater extent than is observed. Mathis and White (2002) discuss voids in ΛCDM simulations,

[63] Pioneering work on the theory of the central mass distribution in a dark mass halo can be found in Dubinski and Carlberg (1991). Moore (1994) and Flores and Primack (1994) were among the first to point out the apparent disagreement between theory and observation.

[64] The situation is reviewed by de Blok *et al.* (2001), and de Blok and Bosma (2002). The galaxy NGC 3109 is a helpful example because it is particularly close—just outside the Local Group—and so particularly well resolved. An optical image is in plate 39 of the *Hubble Atlas of Galaxies* (Sandage, 1961b). The radial velocity measurements across the face of the galaxy, in Figs. 1 and 2 in Blais-Ouellette, Amram, and Carignan (2001), are consistent with circular motion with $v_c\propto r$ at $r\lesssim2$ kpc.

[65] In a power-law halo with $\rho\propto r^{-\gamma}$, the velocity dispersion varies with radius as $\langle v^2\rangle\sim GM(<r)/r\propto r^{2-\gamma}$. The particle-scattering cross section must be adjusted to erase the effective temperature gradient, thus lowering the mass density at small radii, without promoting unacceptable core collapse.

but do not address the trespassing issue. The reader is invited to compare the relative distributions of big and little galaxies in the simulation in Fig. 1 of Mathis and White (2002) with the examples of observed distributions in Figs. 1 and 2 in Peebles (1989b) and in Figs. 1–3 in Peebles (2001).

The community belief is that the trespassing issue need not be a problem for the CDM model: the low-mass density in voids disfavors formation of galaxies from the debris left in these regions. But we have not seen an explanation of why local upward mass fluctuations, of the kind that produce normal galaxies in populated regions, and appear also in the predicted debris in CDM voids, fail to produce dwarf or irregular void galaxies. An easy explanation is that the voids contain no matter, having been gravitationally emptied by the growth of primeval non-Gaussian mass density fluctuations. The evidence in tests 10 and 11 in Sec. IV.B is that the initial conditions are close to Gaussian. But non-Gaussian initial conditions that reproduce the character of the galaxy distribution, including suppression of the trespassing effect, would satisfy test 10 by construction.

We mention finally the related issues of when the large elliptical galaxies formed and when they acquired the central compact massive objects that are thought to be remnant quasar engines (Lynden-Bell, 1969).

In the CDM model large elliptical galaxies form in substantial numbers at a redshift $z<1$. Many astronomers do not see this as a problem, because ellipticals tend to contain relatively young star populations, and some elliptical galaxies have grown by recent mergers, as predicted by the CDM model.[66] But prominent merger events are rare, and the young stars seen in ellipticals generally seem to be a "frosting" (Trager *et al.*, 2000) of recent star formation on a dominant old star population. The straightforward reading of the evidence assembled in Peebles (2002) is that most of the large ellipticals are present as assembled galaxies of stars at $z=2$.[67] The ΛCDM model prediction is uncertain because it depends on the complex processes of star formation that are so difficult to model. The reading of the situation by Thomas and Kauffmann (1999) showed that the predicted abundance of giant ellipticals at $z=2$ is less than about one-third of what it is now. Deciding whether the gap between theory and observation can be closed is not yet straightforward.

A related issue is the significance of the observations of quasars at a redshift $z\sim6$. By conventional estimates[68] these quasars are powered by black holes with masses at the upper end of the range of masses of the central compact objects—let us call them black-hole quasar remnants—in the largest present-day elliptical galaxies. Here are some options to consider. First, the high-redshift quasars may be in the few large galaxies that have already formed at $z\sim6$. Wyithe and Loeb (2002), following Efstathiou and Rees (1988), showed that this fits the ΛCDM model if the quasars at $z\sim6$ have a black-hole mass $\sim10^9 M_\odot$ in dark halos with a mass $\sim10^{12} M_\odot$. In the ΛCDM picture these early galaxies would be considerably denser than normal galaxies; to be checked is whether they would be rare enough to be observationally acceptable. Second, the quasars at $z\sim6$ may have been in more modest star clusters that later grew by merging into giant ellipticals. To be established is whether this growth would preserve the remarkably tight correlation between the central black-hole mass and the velocity dispersion of the stars,[69] and whether growth by merging would produce an acceptable upper bound on black-hole masses at the present epoch. Third, large ellipticals might have grown by accretion around preexisting black holes, without much merging. This is explored by Danese *et al.* (2002).

There does not seem to be a coherent pattern to the present list of challenges for the CDM model. The rota-

[66]The classic merger example is also the nearest large elliptical galaxy, Centaurus A (NGC 5128). The elliptical image is crossed by a band of gas and dust that is likely the result of a merger with one of the spiral galaxies in the group around this elliptical. For a thorough review of what is known about this galaxy see Israel (1998).

[67]Papovich, Dickinson, and Ferguson (2002) found evidence to show that the comoving number density of all galaxies with a star mass greater than $1\times10^{10} M_\odot$, where $M_\odot$ is the mass of the Sun, is significantly less at a redshift $z>1$ than it is now. This is at least roughly in line with the distribution of star ages in the Milky Way spiral galaxy: the bulge stars are old, while the stars in the thin disk have a broad range of ages. Thus if this galaxy evolved from $z=2$ without significant growth by mergers its star mass at $z=2$ would be significantly less than the present value, which is about $5\times10^{10} M_\odot$. Cimatti *et al.* (2002) showed that the redshift distribution of faint galaxies selected at wavelength $\lambda\sim2\ \mu$m is not inconsistent with the fact that galaxy evolution at $z<2$ is dominated by ongoing star formation rather than by merging.

[68]The quasars discovered in the Sloan Digital Sky Survey are discussed by Fan *et al.* (2001). If the quasar radiation is not strongly beamed toward us, its luminosity translates to an Eddington mass (the mass at which the gravitational pull on unshielded plasma balances the radiation pressure) $M_{\rm bh}\sim10^{9.3} M_\odot$. In a present-day elliptical galaxy with this mass in the central compact object, the line-of-sight velocity dispersion is $\sigma\simeq350\ {\rm km\,s^{-1}}$. This is close to the highest velocity dispersion observed in low-redshift elliptical galaxies. For example, in the Faber *et al.* (1989) catalog of 500 ellipticals, 15 have $300<\sigma<400\ {\rm km\,s^{-1}}$, and none a larger σ. From the present-day relation between σ and luminosity, an elliptical galaxy with $\sigma=350\ {\rm km\,s^{-1}}$ has mass $\sim10^{12.3} M_\odot$ in stars. The dark matter associated with this many baryons is $M_{\rm DM}\sim10^{13} M_\odot$. This is a large mass to assemble at $z\sim6$, but it helps that such objects are rare. The present number density of giant elliptical galaxies with $\sigma>300\ {\rm km\,s^{-1}}$ is about $10^{-5}\ {\rm Mpc^{-3}}$, four orders of magnitude more than the comoving number density of quasars detected at $z\sim6$.

[69]Ferrarese and Merritt (2000) and Gebhardt *et al.* (2000) showed that the black-hole mass correlates with the velocity dispersion of the stars in an elliptical galaxy and the velocity dispersion of the bulge stars in a spiral galaxy. This is not a direct gravitational effect: the black-hole mass is less than 1% of the star mass in the bulge or the elliptical galaxy.

tion curves of galaxies with low surface brightness suggest we want to suppress the primeval density fluctuations on small scales, but the observations of what seem to be mature elliptical galaxies at high redshifts suggest we want to increase small-scale fluctuations, or maybe postulate non-Gaussian fluctuations that grow into the central engines for quasars at $z \sim 6$. We do not want these central engines to appear in galaxies with low surface brightness, of course.

It would not be at all surprising if the confusion of challenges proved to be at least in part due to the difficulty in comparing necessarily schematic analytic and numerical model analyses to the limited and indirect empirical constraints. But it is also easy to imagine that the CDM model has to be refined because the physics of the dark sector of matter and energy is more complicated than ΛCDM, and maybe even more complicated than any of the alternatives now under discussion. Perhaps some of the structure-formation ideas considered a decade ago, which invoke good physics, also will prove to be significant factors in relieving the problems with structure formation. And the important point for our purpose is that we do not know how the relief might affect the cosmological tests.

B. The tests

The literature on the cosmological tests is enormous compared to just a decade ago, and it is growing. Our references to this literature are much sparser than in Sec. III, on the principle that no matter how complete the list it will be out of date by the time this review is published. For the same reason, we do not attempt to present the best values of the cosmological parameters based on their joint fit to the full suite of present measurements. The situation will continue to evolve as the measurements improve, and the state of the art is best followed on astro-ph. We do take it to be our assignment to consider what the tests are testing, and to assess the directions in which the results seem to be leading. The latter causes us to return many times to two results that seem secure because they are so well checked by independent lines of evidence, as follows.

First, at present, the tests indicate that optically selected galaxies are useful mass tracers. By that we mean that the assumption that visible galaxies trace mass does not seriously degrade the accuracy of analyses of the observations. This will change as the measurements improve, of course, but the case is strong enough now that we suspect the evidence will continue to point to the fact that optically selected galaxies are good indicators of the location of most of the mass at the present epoch. Second, the mass density in matter is significantly less than the critical Einstein–de Sitter value. The case is compelling because it is supported by so many different lines of evidence (as summarized in Sec. IV.C). Each could be compromised by systematic error, to be sure, but it seems quite unlikely that the evidence would be so consistent yet misleading. A judgement of the range of likely values of the mass density is more difficult. Our estimate, based on the measurements we most trust, is

$$0.15 \lesssim \Omega_{M0} \lesssim 0.4, \tag{59}$$

and we would put the central value at $\Omega_{M0} \simeq 0.25$. The spread is meant in the sense of two standard deviations: we would be surprised to find Ω_{M0} is outside this range.

Several other policy decisions should be noted. The first is that we do not comment on tests that have been considered but not yet applied in a substantial campaign of measurements. A widely discussed example is the Alcock and Paczyński (1979) comparison of the apparent depth and width of a system from its angular size and depth in redshift.

In analyses of the tests of models for evolving dark-energy density, simplicity recommends the XCDM parametrization with a single constant parameter w_X, as is demonstrated by the large number of recent papers on this approach. But the more complete physics recommends the scalar field model with an inverse-power-law potential. This includes the response of the spatial distribution of the dark energy to the peculiar gravitational field. Thus our comments on variable dark-energy density are more heavily weighted to the scalar field model than is the case with the recent literature.

The gravitational deflection of light appears not only as a tool in cosmological tests, such as gravitational lensing, but also as a source of systematic error. The gravitational deflections caused by mass concentrations magnify the image of a galaxy along a line of sight at which the mass density is larger than the average, and reduce the solid angle of the image when the mass density along the line of sight is low. The observed energy flux density is proportional to the solid angle (because the surface brightness, $\mathrm{erg\,cm^{-2}\,s^{-1}\,ster^{-1}\,Hz^{-1}}$, is conserved at a fixed redshift). Selection can be biased either way, by the magnification effect or obscuration by the dust that tends to accompany mass.[70] When the tests are more precise we will have to correct them for these biases, through models for the mass distribution (as in Premadi *et al.*, 2001), and the measurements of the associated gravitational shear of the shapes of the galaxy images. But the biases seem to be small and will not be discussed here.

And finally, as the cosmological tests improve a satisfactory application will require a joint fit of all of the parameters to all of the relevant measurements and constraints. Until recently it made sense to impose prior conditions, most famously the hope that if the universe is not well described by the Einstein–de Sitter model then surely either Λ is negligibly small or space curva-

[70]This was recognized by Zel'dovich (1964), R. Feynman, in 1964, and S. Refsdal, in 1965. Feynman's comments in a colloquium are noted by Gunn (1967). Peebles attended Refsdal's lecture at the International Conference on General Relativity and Gravitation, London, July 1965; Refsdal (1970) mentions the lecture.

ture may be neglected. We suspect that the majority of the community still expect this to be true, on the basis of the coincidences argument in Sec. II.B.2, but it will be important to see the results of joint fits of both Ω_{M0} and $\Omega_{\Lambda 0}$, as well as all the other parameters, as is becoming current practice. We believe our test-by-test discussion is useful for sorting out the physics and astronomy; however, it is not the prototype for the coming generations of precision applications of the tests.

Our remarks are ordered by our estimates of the model dependence.

1. Thermal cosmic microwave background radiation

We are in a sea of radiation with a spectrum very close to Planck at $T=2.73$ K, and isotropic to one part in 10^5 (after correction for a dipole term that usually is interpreted as the result of our motion relative to the rest frame defined by the radiation).[71] The thermal spectrum indicates thermal relaxation, for which the optical depth has to be large on the scale of the Hubble length H_0^{-1}. We know that space now is close to being transparent at the wavelengths of this radiation, because radio galaxies are observed at a high redshift. Thus the universe has to have expanded from a state quite different from now, when it was hotter, denser, and optically thick. This is strong evidence that our universe is evolving.

This interpretation depends on, and checks, conventional local physics with a single metric description of spacetime. Under these assumptions the expansion of the universe preserves the thermal spectrum and cools the temperature as[72]

$$T \propto (1+z). \tag{60}$$

Bahcall and Wolf (1968) point out that one can test this temperature-redshift relation by measurements of the excitation temperatures of fine-structure absorption line systems in gas clouds along quasar lines of sight. The corrections for excitations by collisions and the local radiation field are subtle, however, and perhaps not yet fully sorted out (as discussed by Molaro *et al.*, 2002, and references therein).

[71]The history of the discovery and measurement of this radiation, and its relation to the light-element abundances in test 2, are presented in Peebles (1971, pp. 121–129 and 240 and 241); Wilkinson and Peebles (1990); Alpher and Herman (2001). The precision spectrum measurements are summarized in Halpern, Gush, and Wishnow (1991); and Fixsen *et al.* (1996).

[72]To see this, recall the normal modes argument used for Eq. (7). The occupation number in a normal mode with wavelength λ at temperature T is the Planck form $\mathcal{N}=[e^{\hbar c/kT\lambda}-1]^{-1}$. Adiabaticity implies that $\mathcal{N}$ is constant. Since the mode wavelength varies as $\lambda \propto a(t)$, where a is the expansion factor in Eq. (4), and $\mathcal{N}$ is close to constant, the mode temperature varies as $T\propto 1/a(t)$. Since the same temperature scaling applies to each mode, an initially thermal sea of radiation remains thermal in the absence of interactions. We do not know the provenance of this argument; it was familiar in Dicke's group when the 3-K cosmic microwave background radiation was discovered.

The 3-K thermal cosmic background radiation is a centerpiece of modern cosmology, but its existence does not test general relativity.

2. Light-element abundances

The best evidence that the expansion and cooling of the universe traces back to a high redshift is the success of the standard model for the origin of deuterium and isotopes of helium and lithium, by reactions among radiation, leptons, and atomic nuclei as the universe expands and cools through temperature $T\sim 1$ MeV at a redshift $z\sim 10^{10}$. The free parameter in the standard model is the present baryon number density. The model assumes the baryons are uniformly distributed at a high redshift, so this parameter with the known present radiation temperature fixes the baryon number density as a function of temperature and the temperature as a function of time. The latter follows from the expansion rate Eq. (11), which at the epoch of light-element formation may be written as

$$\left(\frac{\dot a}{a}\right)^2 = \frac{8}{3}\pi G \rho_r, \tag{61}$$

where the mass density ρ_r counts radiation, which is now at $T=2.73$ K, the associated neutrinos, and $e^{\pm}$ pairs. The curvature and Λ terms are unimportant, unless the dark-energy mass density varies quite rapidly.

Independent analyses of the fit to the measured element abundances, corrected for synthesis and destruction in stars, by Burles, Nollett, and Turner (2001), and Cyburt, Fields, and Olive (2001), indicate

$$0.018 \leqslant \Omega_{B0}h^2 \leqslant 0.022, \quad 0.006 \leqslant \Omega_{B0}h^2 \leqslant 0.017, \tag{62}$$

both at 95% confidence limits. Other analyses by Coc *et al.* (2002) and Thuan and Izotov (2002) result in ranges that lie between those of Eq. (62). The difference in values may be a useful indication of remaining uncertainties; it is mostly a consequence of the choice of isotopes used to constrain $\Omega_{B0}h^2$. Burles *et al.* (2001) use the deuterium abundance, Cyburt *et al.* (2001) favor the helium and lithium measurements, and the other two groups use other combinations of abundances. Equation (62) is consistent with the summary range, $0.0095 \leqslant \Omega_{B0}h^2 \leqslant 0.023$ at 95% confidence, of Fields and Sarkar (2002).

The baryons observed at a low redshift, in stars and gas, amount to (Fukugita, Hogan, and Peebles, 1998)

$$\Omega_{B0} \sim 0.01. \tag{63}$$

It is plausible that the difference between Eqs. (62) and (63) is in cool plasma, with temperature $T\sim 100$ eV, in groups of galaxies. It is difficult to observationally constrain the idea that there is a good deal more cool plasma in the large voids between the concentrations of galaxies. A more indirect but eventually more precise constraint on Ω_{B0}, from the anisotropy of the 3-K thermal cosmic microwave background radiation, is discussed in test 11.

It is easy to imagine complications, such as inhomogeneous entropy per baryon, or in the physics of neutrinos; examples may be traced back through Abazajian, Fuller, and Patel (2001) and Giovannini, Keihänen, and Kurki-Suonio (2002). It seems difficult to imagine that a more complicated theory would reproduce the successful predictions of the simple model, but Nature fools us on occasion. Thus before concluding that the theory of the prestellar light-element abundances is known, apart from the addition of decimal places to the cross sections, it is best to wait and see what advances in the physics of baryogenesis and of neutrinos teach us.

How is general relativity probed? The only part of the computation that depends specifically on this theory is the pressure term in the active gravitational mass density, in the expansion-rate equation (8). If we did not have general relativity, a simple Newtonian picture might have led us to write down $\ddot{a}/a = -4\pi G\rho_r/3$ instead of Eq. (8). With $\rho_r \sim 1/a^4$, which is appropriate since most of the mass is fully relativistic at the redshifts of light-element production, this would predict that the expansion time $a/\dot{a}$ is $2^{1/2}$ times the standard expression [that from Eq. (61)]. The larger expansion time would hold the neutron-to-proton number density ratio close to that at thermal equilibrium, $n/p = e^{-Q/kT}$, where Q is the difference between the neutron and proton masses, to lower temperature. It would also allow more time for free decay of the neutrons after thermal equilibrium is broken. Both effects decrease the final ^{4}He abundance. The factor of the $2^{1/2}$ increase in expansion time would reduce the helium abundance by mass to $Y \sim 0.20$. This is significantly less than what is observed in objects with the lowest heavy-element abundances, and so seems to be ruled out (Steigman, 2002).[73] That is, we have positive evidence for the relativistic expression for the active gravitational mass density at the redshift $z \sim 10^{10}$, a striking result.

3. Expansion times

The predicted time of expansion from the very early universe to redshift z is

$$t(z) = \int \frac{da}{\dot{a}} = H_0^{-1} \int_z^\infty \frac{dz}{(1+z)E(z)}, \tag{64}$$

where $E(z)$ is defined in Eq. (11). If $\Lambda = 0$ the present age is $t_0 < H_0^{-1}$. In the Einstein–de Sitter model the present age is $t_0 = 2/(3H_0)$. If the dark-energy density is significant and evolving, we may write $\rho_\Lambda = \rho_{\Lambda 0} f(z)$, where the function of the redshift is normalized to $f(0) = 1$. Then $E(z)$ generalizes to

$$E(z) = [\Omega_{M0}(1+z)^3 + \Omega_{R0}(1+z)^4 + \Omega_{K0}(1+z)^2 + \Omega_{\Lambda 0} f(z)]^{1/2}. \tag{65}$$

In the XCDM parametrization with constant w_X [Eq. (45)], $f(z) = (1+z)^{3(1+w_X)}$. Olson and Jordan (1987) presented the earliest discussion we have found of $H_0 t_0$ in this picture (before it acquired the name). In scalar field models, $f(z)$ generally must be evaluated numerically; examples are in Peebles and Ratra (1988).

The relativistic correction to the active gravitational mass density [Eq. (8)] is not important at the redshifts at which galaxies can be observed and the ages of their star populations estimated. At a moderately high redshift, where the nonrelativistic matter term dominates, Eq. (64) is approximately

$$t(z) \simeq \frac{2}{3H_0 \Omega_{M0}^{1/2}} (1+z)^{-3/2}. \tag{66}$$

That is, the ages of star populations at high redshift are an interesting probe of Ω_{M0} but they are not very sensitive to space curvature or to a near constant dark-energy density.[74]

Recent analyses of the ages of old stars[75] indicate that the expansion time is in the range

$$11 \text{ Gyr} \lesssim t_0 \lesssim 17 \text{ Gyr}, \tag{67}$$

at 95% confidence, with a central value $t_0 \simeq 13$ Gyr. Following Krauss and Chaboyer (2001) these numbers add 0.8 Gyr to the star ages, under the assumption that star formation commenced no earlier than $z = 6$ [Eq. (66)]. A naive addition in quadrature to the uncertainty in H_0 [Eq. (6)] indicates that the dimensionless age parameter is in the range

$$0.72 \lesssim H_0 t_0 \lesssim 1.17, \tag{68}$$

at 95% confidence, with a central value $H_0 t_0 \simeq 0.89$. The uncertainty here is dominated by that in t_0. In the spatially flat ΛCDM model ($\Omega_{K0} = 0$), Eq. (68) translates to $0.15 \lesssim \Omega_{M0} \lesssim 0.8$, with the central value $\Omega_{M0} \simeq 0.4$. In the open model with $\Omega_{\Lambda 0} = 0$, the constraint is $\Omega_{M0} \lesssim 0.6$

[73]There is a long history of discussions of this probe of the expansion rate at the redshifts of light-element production. The reduction of the helium abundance to $Y \sim 0.2$ if the expansion time is increased by the factor $2^{1/2}$ is seen in Figs. 1 and 2 in Peebles (1966). Dicke (1968) introduced the constraint on evolution of the strength of the gravitational interaction; see Uzan (2003) for a recent review. The effect of the number of neutrino families on the expansion rate and hence the helium abundance is noted by Hoyle and Tayler (1964) and Shvartsman (1969). Steigman, Schramm, and Gunn (1977) discuss the importance of this effect as a test of cosmology and of the particle physics measures of the number of neutrino families.

[74]The predicted maximum age of star populations in galaxies at redshifts $z \gtrsim 1$ still depends on $\Omega_{\Lambda 0}$ and Ω_{K0}, and there is the advantage that the predicted maximum age is much shorter than it is today. This variant of the expansion time test is discussed by Nolan *et al.* (2001), Lima and Alcaniz (2001), and references therein.

[75]See Carretta *et al.* (2000), Krauss and Chaboyer (2001), Chaboyer and Krauss (2002), and references therein.

with the central value $\Omega_{M0} \simeq 0.1$. In the inverse-power-law scalar field dark-energy case (Sec. II.C) with power-law index $\alpha = 4$, the constraint is $0.05 \lesssim \Omega_{M0} \lesssim 0.8$.

We should pause to admire the unification of the theory and measurements of stellar evolution in our galaxy, which yield the estimate of t_0, and the measurements of the extragalactic distance scale, which yield H_0, in the product of Eq. (68) that agrees with the relativistic cosmology with dimensionless parameters in the range now under discussion. As we indicated in Sec. III, there is a long history of discussion of the expansion time as a constraint on cosmological models. The measurements now are tantalizingly close to a check of consistency with the values of Ω_{M0} and $\Omega_{\Lambda 0}$ indicated by other cosmological tests.

4. The redshift-angular-size and redshift-magnitude relations

An object at the redshift z with physical length l perpendicular to the line of sight subtends angle θ such that

$$l = a(t) r(z) \theta = a_0 r(z) \theta/(1+z), \tag{69}$$

where $a_0 = a(t_0)$. The angular size distance $r(z)$ is the coordinate position of the object in the first line element in Eq. (15), with the observer placed at the origin. The condition that light moves from source to observer on a radial null geodesic is

$$\int_0^{r(z)} \frac{dr}{\sqrt{1+Kr^2}} = \int \frac{dt}{a(t)}, \tag{70}$$

which gives

$$H_0 a_0 r(z) = \frac{1}{\sqrt{\Omega_{K0}}} \sinh\left[\sqrt{\Omega_{K0}} \int_0^z \frac{dz}{E(z)} \right], \tag{71}$$

where $E(z)$ is defined in Eqs. (11) and (65).

In the Einstein–de Sitter model, the angular-size-redshift relation is

$$\theta = \frac{H_0 l}{2} \frac{(1+z)^{3/2}}{\sqrt{1+z}-1}. \tag{72}$$

At $z \ll 1$, $\theta = H_0 l/z$, consistent with the Hubble redshift-distance relation. At $z \gg 1$ the image is magnified,[76] $\theta \propto 1+z$.

The relation between the luminosity of a galaxy and the energy flux density received by an observer follows from Liouville's theorem: the observed energy flux i_{ν_0} per unit time, area, solid angle, and frequency satisfies

$$i_{\nu_0} \delta\nu_0 = i_{\nu_e} \delta\nu_e/(1+z)^4, \tag{73}$$

with i_{ν_e} the emitted energy flux (surface brightness) at the source and $\delta\nu_e = \delta\nu_0(1+z)$ the bandwidth at the source at redshift z. The redshift factor $(1+z)^4$ appears for the same reason it does in the 3-K cosmic microwave background radiation energy density. With Eq. (69) to fix the solid angle, Eq. (73) shows that the observed energy flux per unit area, time, and frequency from a galaxy at redshift z that has luminosity L_{ν_e} per frequency interval measured at the source is

$$f_{\nu_0} = \frac{L_{\nu_e}}{4\pi a_0^2 r(z)^2 (1+z)}. \tag{74}$$

In conventional local physics with a single metric theory the redshift-angular-size [Eq. (69)] and redshift-magnitude [Eq. (74)] relations are physically equivalent.[77]

The best present measurement of the redshift-magnitude relation uses type Ia supernovae.[78] The results are inconsistent with the Einstein–de Sitter model, at enough standard deviations to make it clear that unless there is something quite substantially and unexpectedly wrong with the measurements, the Einstein–de Sitter model is ruled out. The data require $\Lambda > 0$ at two to three standard deviations, depending on the choice of data and method of analysis (Leibundgut, 2001; Gott *et al.*, 2001). The spatially flat case with Ω_{M0} in the range

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[76] To our best knowledge, Hoyle (1959) published the earliest discussion on the magnification effect. In the coordinate system in Eq. (15), with the observer at the origin, light rays from the object move to the observer along straight radial lines. An image at a high redshift is magnified because the light detected by the observer is emitted when the proper distance to the object measured at a fixed world time is small. Because the proper distance between the object and source increases faster than the speed of light, emitted light directed at the observer initially moves away from the observer.

[77] For a review of measurements of the redshift-magnitude relation (and other cosmological tests) we recommend Sandage (1988). A recent application to the most luminous galaxies in clusters can be found in Aragón-Salamanca, Baugh, and Kauffmann (1998). The redshift-angular-size relation is measured by Daly and Guerra (2001) for radio galaxies, Buchalter *et al.* (1998) for quasars, and Gurvits, Kellermann, and Frey (1999) for compact radio sources. Constraints on the cosmological parameters from the Gurvits *et al.* data are discussed by Vishwakarma (2001), Lima and Alcaniz (2002), Chen and Ratra (2003a), and references therein, and constraints based on the radio galaxy data are discussed by Daly and Guerra (2001), Podariu *et al.* (2003), and references therein.

[78] These supernovae are characterized by the absence of hydrogen lines in the spectra; they are thought to be the result of explosive nuclear burning of white dwarf stars. Pskovskii (1977) and Phillips (1993) pioneered the reduction of the supernovae luminosities to a near universal standard candle. For recent discussions of their use as a cosmological test see Goobar and Perlmutter (1995), Riess *et al.* (1998), Perlmutter *et al.* (1999a), Gott *et al.* (2001), and Leibundgut (2001). We recommend Leibundgut's (2001) cautionary discussion of astrophysical uncertainties, which includes the unknown nature of the trigger for the nuclear burning, the possibility that the Phillips correction to a fiducial luminosity actually depends on the redshift or environment within a galaxy, and possible obscuration by intergalactic dust. There are also issues of physics that may affect this test (and others): strengths of the gravitational or electromagnetic interactions may vary with time, and photon-axion conversion may reduce the number of photons that reach us. All of the above is under active study.

of Eq. (59) is a good fit for constant Λ. The current data do not provide interesting constraints on the models for evolving dark-energy density.[79] Perlmutter *et al.* (http://snap.lbl.gov/) show that a tighter constraint, from supernovae observations to a redshift $z \sim 2$, by the proposed Supernova Acceleration Probe (SNAP) satellite, is feasible and capable of giving a significant detection of Λ and maybe its evolution.[80]

5. Galaxy counts

Counts of galaxies—or of other objects whose number density as a function of the redshift may be modeled—probe the volume element $(dV/dz)dz\,\delta\Omega$ defined by a solid angle $\delta\Omega$ in the sky and a redshift interval dz. The volume is fixed by the angular size distance [Eq. (69)], which determines the area subtended by the solid angle, in combination with the redshift-time relation [Eq. (64)], which fixes the radial distance belonging to the redshift interval.

Sandage (1961a) and Brown and Tinsley (1974) showed that with the technology then available galaxy counts were not a very sensitive probe of the cosmological parameters. Loh and Spillar (1986) opened the modern exploration of the galaxy-count–redshift relation at redshifts near unity, where the predicted counts are quite different in models with and without a cosmological constant (as illustrated in Fig. 13.8 in Peebles, 1993).

The interpretation of galaxy counts requires an understanding of the evolution of galaxy luminosities and the gain and loss of galaxies by merging. Here is an example of the former in a spatially flat cosmological model with $\Omega_{\rm M0}=0.25$. The expansion time from a high redshift is $t_3=2.4$ Gyr at a redshift $z=3$ and $t_0=15$ Gyr now. Consider a galaxy observed at $z=3$. Suppose the bulk of the stars in this galaxy formed at time t_f, and the population then aged and faded without significant later star formation. Then if $t_f \ll t_3$, the ratio of the observed luminosity at $z=3$ to its present luminosity would be (Tinsley, 1972; Worthey, 1994)

$$L_3/L_0 \simeq (t_0/t_3)^{0.8} \simeq 4. \qquad (75)$$

If t_f were larger, but still less than t_3, this ratio would be larger. If t_f were greater than t_3 the galaxy would not be seen, absent earlier generations of stars. In a more realistic picture significant star formation may be distributed over a considerable range of redshifts, and the effect on the typical galaxy luminosity at a given redshift accordingly more complicated. Since there are many more galaxies with low luminosities than with high luminosities, one has to know the luminosity evolution quite well for a meaningful comparison of galaxy counts at high and low redshifts. The present situation is illustrated by the rather different indications from studies by Phillipps *et al.* (2000) and Totani *et al.* (2001).

The understanding of galaxy evolution and the interpretation of galaxy counts will be improved by large samples of galaxy counts as a function of color, apparent magnitude, and redshift. Newman and Davis (2000) point to a promising alternative: count galaxies as a function of the internal velocity dispersion that in spirals correlates with the dispersion in the dark-matter halo. That could eliminate the need to understand the evolution of star populations. There is still the issue of evolution of the dark halos by merging and accretion, but that might be reliably modeled by numerical simulations within the CDM picture. Either way, with further work galaxy counts may provide an important test for dark energy and its evolution (Newman and Davis, 2000; Huterer and Turner, 2001; Podariu and Ratra, 2001).

6. The gravitational lensing rate

The probability of production of multiple images of a quasar or a radio source by gravitational lensing by a foreground galaxy, or of strongly lensed images of a galaxy by a foreground cluster of galaxies, adds the relativistic expression for deflection of light to the physics of the homogeneous cosmological model. Fukugita, Futamase, and Kasai (1990) and Turner (1990) point out the value of this test: at small $\Omega_{\rm M0}$ the predicted lensing rate is considerably larger in a flat model with Λ than in an open model with $\Lambda=0$ (as illustrated in Fig. 13.12 in Peebles, 1993).

The measurement problem for the analysis of quasar lensing is that quasars that are not lensed are not magnified by lensing, making them harder to find and the correction for completeness of detection harder to establish. Present estimates (Falco, Kochanek, and Muñoz, 1998; Helbig *et al.*, 1999) do not seriously constrain $\Omega_{\rm M0}$ in an open model, and in a flat model ($\Omega_{\rm K0}=0$) suggest $\Omega_{\rm M0}>0.36$ at 2σ. This is close to the upper bound in Eq. (59). Earlier indications that the lensing rate in a flat model with constant Λ requires a larger value of $\Omega_{\rm M0}$ than is suggested by galaxy dynamics led Ratra and Quillen (1992) and Waga and Frieman (2000) to investigate the inverse-power-law potential dark-energy scalar field case. They showed this can significantly lower the predicted lensing rate at $\Omega_{\rm K0}=0$ and small $\Omega_{\rm M0}$. The lensing rate is still too uncertain for us to draw conclusions on this point, but advances in the measurement certainly will be followed with interest.

The main problem in the interpretation of the rate of strong lensing of galaxies by foreground clusters as a cosmological test is the sensitivity of the lensing cross section to the mass distribution within the cluster (Wu and Hammer, 1993); for the present, still somewhat uncertain, state of the art see Cooray (1999), and references therein.

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[79]Podariu and Ratra (2000) and Waga and Frieman (2000) discuss the redshift-magnitude relation in the inverse-power-law scalar field model, and Waga and Frieman (2000) and Ng and Wiltshire (2001) discuss this relation in the massive scalar field model.

[80]Podariu, Nugent, and Ratra (2001); Wang and Lovelace (2001); Weller and Albrecht (2002); Gerke and Efstathiou (2002); Eriksson and Amanullah (2002), and references therein, discuss constraints on cosmological parameters from the proposed SNAP mission.

7. Dynamics and the mean mass density

Estimates of the mean mass density from the relation between the mass distribution and the resulting peculiar velocities,[81] and from the gravitational deflection of light, probe gravity physics and constrain Ω_{M0}. The former is not sensitive to Ω_{K0}, $\Omega_{\Lambda 0}$, or the dynamics of the dark energy, the latter only through the angular size distances.

We begin with the redshift space of observed galaxy angular positions and redshift distances z/H_0 in the radial direction. The redshift z has a contribution from the radial peculiar velocity, which is a probe of the gravitational acceleration produced by the inhomogeneous mass distribution. The two-point correlation function, ξ_v, in redshift space is defined by the probability that a randomly chosen galaxy has a neighbor at distance $r_\parallel$ along the line of sight in redshift space and perpendicular distance $r_\perp$,

$$dP = n[1+\xi_v(r_\parallel, r_\perp)]dr_\parallel d^2 r_\perp, \tag{76}$$

where n is the galaxy number density. This is the usual definition of a reduced correlation function. Peculiar velocities make the function anisotropic. On small scales the random relative peculiar velocities of the galaxies broaden ξ_v along the line of sight. On large scales the streaming peculiar velocity of convergence to gravitationally growing mass concentrations flattens ξ_v along the line of sight.[82]

At 10 kpc$\lesssim hr_\perp \lesssim 1$ Mpc the measured line-of-sight broadening is prominent, and indicates the one-dimensional relative velocity dispersion is close to independent of $r_\perp$ at $\sigma \sim 300$ km s^{-1}.[83] This is about what would be expected if the mass two- and three-point correlation functions were well approximated by the galaxy correlation functions, the mass clustering on these scales were close to statistical equilibrium, and the density parameter were in the range of Eq. (59).

We have a check from the motions of the galaxies in and around the Local Group of galaxies, where the absolute errors in the measurements of galaxy distances are least. The two largest group members are the Andromeda Nebula (M 31) and our Milky Way galaxy. If they contain most of the mass their relative motion is the classical two-body problem in Newtonian mechanics (with minor corrections for Λ, mass accretion at low redshifts, and the tidal torques from neighboring galaxies). The two galaxies are separated by 800 kpc and approach at 110 km s^{-1}. In the minimum mass solution the galaxies have completed just over half an orbit in the cosmological expansion time $t_0 \sim 10^{10}$ yr. By this argument Kahn and Woltjer (1959) found that the sum of masses of the two galaxies had to be an order of magnitude larger than what is seen in the luminous parts. An extension to the analysis of the motions and distances of the galaxies within 4-Mpc distance from Earth, taking account of the gravitational effects of the galaxies out to 20-Mpc distance, yields masses that are quite similar to what Kahn and Woltjer found, and consistent with Ω_{M0} in the range of Eq. (59) (Peebles *et al.*, 2001).

We have another check from weak lensing: the shear distortion of images of distant galaxies by the gravitational deflection caused by the inhomogeneous mass distribution.[84] If galaxies trace mass, these measurements show that the matter density parameter measured on scales from about 1 Mpc to 10 Mpc is in the range of Eq. (59). It will be interesting to see whether these measurements can check the factor-of-2 difference between the relativistic gravitational deflection of light and the naive Newtonian deflection angle.

The redshift space correlation function ξ_v [Eq. (76)] is measured well enough at $hr_\perp \sim 10$ Mpc to demonstrate the flattening effect, again consistent with Ω_{M0} in the range of Eq. (59), if galaxies trace mass. Similar numbers follow from galaxies selected as far-infrared Infrared Astronomical Satellite (IRAS) sources (Tadros *et al.*, 1999) and from optically selected galaxies (Padilla *et al.*, 2001; Peacock *et al.*, 2001). The same physics, applied to estimates of the mean relative peculiar velocity of galaxies at separations ~ 10 Mpc, yet again indicates a similar density parameter (Juszkiewicz *et al.*, 2000).

Other methods of analysis of the distributions of astronomical objects and peculiar velocities smoothed over scales $\gtrsim 10$ Mpc give a variety of results for the mass density, some above the range of Eq. (59),[85] others towards the bottom end of the range (Branchini *et al.*, 2001). The measurement of Ω_{M0} from large-scale streaming velocities thus remains open. But we are im-

[81]Early estimates of the mean mass density, by Hubble (1936, p. 189) and Oort (1958), combine the galaxy number density from galaxy counts with estimates of galaxy masses from the internal motions of gas and stars. Hubble (1936, p. 180) was quite aware that this misses mass between the galaxies, and that the motions of galaxies within clusters suggests that there is a lot more intergalactic mass (Zwicky, 1933; Smith, 1936). For a recent review of this subject see Bahcall *et al.* (2000).

[82]This approach grew out of the statistical method introduced by Geller and Peebles (1973); it is derived in its present form in Peebles (1980b) and first applied to a serious redshift sample in Davis and Peebles (1983b). These references give the theory for the second moment σ^2 of ξ_v in the radial direction—the mean-square relative peculiar velocity—in the small-scale stable clustering limit. The analysis of the anisotropy of ξ_v in the linear perturbation theory of large-scale flows [Eq. (55)] is presented in Kaiser (1987).

[83]This measurement requires close attention to clusters that contribute little to the mean mass density but are responsible for a broad and difficult-to-measure tail in the distribution of relative velocities. Details may be traced back through Padilla *et al.* (2001), Peacock *et al.* (2001), and Landy (2002).

[84]Recent studies include Wilson, Kaiser, and Luppino (2001); Van Waerbeke *et al.* (2002); Refregier, Rhodes, and Groth (2002), Bacon *et al.* (2002); and Hoekstra, Yee, and Gladders (2002). See Munshi and Wang (2003), and references therein, for discussions of how weak lensing might probe dark energy.

[85]The methods and results may be traced through Fisher, Scharf, and Lahav (1994), Sigad *et al.* (1998), and Branchini *et al.* (2000).

pressed by an apparently simple local situation, the peculiar motion of the Local Group toward the Virgo cluster of galaxies. This is the nearest known large mass concentration, at distance ~ 20 Mpc. Burstein (2000) found that our Virgo-centric velocity is $v_v = 220\ \mathrm{km\,s^{-1}}$, indicating $\Omega_{M0} \simeq 0.2$ (Davis and Peebles, 1983a, Fig. 1). This leads us to conclude that the weight of the evidence from dynamics on scales ~ 10 Mpc favors low Ω_{M0}, in the range of Eq. (59).

None of these measurements is precise. But many have been under discussion for a long time and seem to us to be reliably understood. Weak lensing is new, but the measurements have been checked by several independent groups. The result, in our opinion, is a well-checked and believable network of evidence that over two decades of well-sampled length scales, 100 kpc to 10 Mpc, the apparent value of Ω_{M0} is constant to a factor of 3 or so, in the range $0.15 \lesssim \Omega_{M0} \lesssim 0.4$. The key point for the purpose of this review is that this result is contrary to what might have been expected from biasing, or from a failure of the inverse-square law (as will be discussed in test 13).

8. The baryon mass fraction in clusters of galaxies

Abell (1958) made the first useful catalog of the rich clusters considered here and in the next test. A typical value of the Abell cluster mass within the Abell radius $r_a = 1.5h^{-1}$ Mpc is $3\times 10^{14} h^{-1} M_\odot$. The cluster masses are reliably measured (within Newtonian gravity) from consistent results from the velocities of the galaxies, the pressure of the intracluster plasma, and the gravitational deflection of light from background galaxies.

White (1992) and White *et al.* (1993) pointed out that rich clusters likely are large enough to contain a close to fair sample of baryons and dark matter, meaning the ratio of baryonic to total mass in a cluster is a good measure of Ω_{B0}/Ω_{M0}. With Ω_{B0} from the model for light elements [Eqs. (62)], this gives a measure of the mean mass density. The baryon mass fraction in clusters is still under discussion.[86] We adopt as the most direct and so perhaps most reliable approach the measurement of the baryonic gas mass fraction of clusters, f_{gas}, through the Sunyaev-Zel'dovich microwave decrement caused by Thomson-Compton scattering of cosmic microwave background radiation by the intracluster plasma. The Carlstrom *et al.* (2001) value for f_{gas} gives $\Omega_{M0} \sim 0.25$,[87] in the range of Eq. (59). This test does not directly constrain Ω_{K0}, $\Omega_{\Lambda 0}$, or the dynamics of the dark energy.

9. The cluster mass function

In the CDM model rich clusters of galaxies grow out of the rare peak upward fluctuations in the primeval Gaussian mass distribution. Within this model one can adjust the amplitude of the mass fluctuations to match the abundance of clusters at one epoch. In the Einstein–de Sitter model it is difficult to see how this one free adjustment can account for the abundance of rich clusters now and at redshifts near unity.[88]

Most authors now agree that the low-density flat ΛCDM model can give a reasonable fit to the cluster abundances as a function of the redshift. The constraint on Ω_{M0} from the present cluster abundance still is under discussion, but generally is found to be close to $\Omega_{M0} \sim 0.3$ if galaxies trace mass.[89] The constraint from the evolution of the cluster number density also is under discussion.[90] The predicted evolution is slower in a lower-density universe, and at given Ω_{M0} the evolution is slower in an open model with $\Lambda = 0$ than in a spatially flat model with Λ (for the reasons discussed in Sec. III.D). Bahcall and Fan (1998) emphasized that we have good evidence for the presence of some massive clusters at $z \sim 1$, and that this is exceedingly difficult to understand in the CDM model in the Einstein–de Sitter cosmology (when biasing is adjusted to yield a reasonable present number density). Low-density models with or without Λ can account for the existence of some massive clusters at a high redshift. Distinguishing between the predictions of the spatially curved and flat low-density cases awaits better measurements.

10. Biasing and the development of nonlinear mass density fluctuations

Elements of the physics of cluster formation in test 9 appear in this test of the early stages in the nonlinear growth of departures from homogeneity. An initially Gaussian mass distribution becomes skew as low-density fluctuations start to bottom out and high-density fluctuations start to develop into prominent mass peaks. The early signature of this nonlinear evolution is the disconnected three-point mass autocorrelation function, $\langle \delta(\vec{x},t)\,\delta(\vec{y},t)\,\delta(\vec{z},t)\rangle$, where $\delta(\vec{x},t) = \delta\rho/\rho$ is the dimensionless mass contrast. If galaxies are useful mass tracers the galaxy three-point function is a good measure of this mass function.

The form for the mass three-point function, for Gaussian initial conditions at a high redshift, in lowest nonzero order in perturbation theory, was worked out in Fry (1984), and Fry (1994) made the point that measurements of the galaxy three-point function test how well

[86]See Hradecky *et al.* (2000); Roussel, Sadat, and Blanchard (2000); Allen, Schmidt, and Fabian (2002), and references therein.

[87]This assumes $\Omega_{B0} h^2 = 0.014$ from Eqs. (62). For the full range of values in Eqs. (6) and (62), $0.1 \lesssim \Omega_{M0} \lesssim 0.4$ at two standard deviations.

[88]Early discussions of this problem include Evrard (1989); Peebles, Daly, and Juszkiewicz (1989); Oukbir and Blanchard (1992).

[89]For recent discussions see Pierpaoli, Scott, and White (2001); Seljak (2002); Viana, Nichol, and Liddle (2002); Ikebe *et al.* (2002); Bahcall *et al.* (2003), and references therein. Wang and Steinhardt (1998) considered this test in the context of the XCDM parametrization; to our knowledge it has not been studied in the scalar field dark-energy case.

[90]Examples include Blanchard *et al.* (2000); and Borgani *et al.* (2001).

galaxies trace mass.[91] There are now two sets of measurements of the galaxy three-point function on scales ~ 10 to 20 Mpc, where the density fluctuations are not far from Gaussian. One uses infrared-selected IRAS galaxies,[92] the other optically selected galaxies (Verde *et al.*, 2002). The latter is consistent with the perturbative computation of the mass three-point function for Gaussian initial conditions. The former shows that infrared-selected galaxies are adequate mass tracers apart from the densest regions, which IRAS galaxies avoid. That has a simple interpretation in astrophysics: galaxies in dense regions tend to be swept clear of the gas and dust that make galaxies luminous in the infrared.

This test provides evidence of consistency for the following three ideas: galaxies are useful mass tracers on scales ~ 10 Mpc, the initial conditions are close to Gaussian, and conventional gravity physics gives an adequate description of this aspect of the growth of structure. It is, in principle, sensitive to $\Omega_{\Lambda 0}$, through the suppression of the growth of small departures from homogeneity at a low redshift, but the effect is small.

11. The anisotropy of the cosmic microwave background radiation

The wonderfully successful CDM prediction of the power spectrum of the angular distribution of the temperature of the 3-K cosmic microwave background radiation has converted many of the remaining skeptics in the cosmology community to the belief that the CDM model likely captures important elements of reality.

Efstathiou (2002) provided a useful measure of the information in the present measurements:[93] the fit to the CDM model significantly constrains three linear combinations of the free parameters. We shall present three sets of considerations that roughly follow Efstathiou's constraints. We begin with reviews of the standard measure of the temperature anisotropy and of the conditions at a redshift $z\sim 1000$ that are thought to produce the observed anisotropy.

The 3-K cosmic microwave background temperature $T(\theta,\phi)$, as a function of position in the sky, is usually expressed as an expansion in spherical harmonics,

$$\delta T(\theta,\phi)=T(\theta,\phi)-\langle T\rangle=\sum_{l,m} a_l^n Y_l^m(\theta,\phi). \tag{77}$$

The square of δT averaged over the sky is

$$\langle \delta T^2\rangle=\frac{1}{4\pi}\sum_l \,(2l+1)\langle |a_l^m|\rangle^2, \tag{78}$$

where $|a_l^m|^2$ is statistically independent of m. This may be rewritten as

$$\langle \delta T^2\rangle=\sum_l \frac{1}{l}\,\delta T_l^2, \quad \delta T_l^2=\frac{l(2l+1)}{4\pi}\langle |a_l^m|^2\rangle. \tag{79}$$

Since Σl^{-1} is close to $\int d\ln l$, δT_l^2 is the variance of the temperature per logarithmic interval of l. A measure of the angular scale belonging to the multipole index l is that the minimum distance between zeros of the spherical harmonic Y_l^m, in longitude or latitude, is $\theta=\pi/l$, except at locations close to the poles, where Y_l^m approaches zero.[94]

Now let us consider the main elements of the physics that determines the 3-K cosmic microwave background anisotropy.[95] At a redshift $z_{\rm dec}\sim 1000$ the temperature reaches the critical value at which the primeval plasma combines to atomic hydrogen (and slightly earlier, to neutral helium). This removes the coupling between baryons and radiation by Thomson scattering, leaving the radiation to propagate nearly freely (apart from residual gravitational perturbations). Ratios of mass densities near the epoch $z_{\rm dec}$, when matter and radiation decouple, are worth noting. At a redshift $z_{\rm eq}=2.4\times 10^4\Omega_{\rm M0}h^2$ the mass density in matter—including the baryonic and nonbaryonic components—is equal to the relativistic mass density in radiation and neutrinos assumed to have low masses. At decoupling the ratio of mass densities is

$$\frac{\rho_{\rm M}(z_{\rm dec})}{\rho_{\rm R}(z_{\rm dec})}=\frac{z_{\rm eq}}{z_{\rm dec}}\sim 20\Omega_{\rm M0}h^2\sim 2, \tag{80}$$

at the central values of the parameters in Eqs. (6) and (59). The ratio of mass densities in baryons and in thermal cosmic microwave background radiation—not counting neutrinos—is

$$\frac{\rho_{\rm B}(z_{\rm dec})}{\rho_{\rm CBR}(z_{\rm dec})}=\frac{4\times 10^4\Omega_{\rm B0}h^2}{z_{\rm dec}}\sim 0.5. \tag{81}$$

That is, the baryons and radiation decouple just as the expansion rate has become dominated by nonrelativistic matter and the baryons are starting to lower the velocity

[91]Other notable contributions to the development of this point include Bernardeau and Schaeffer (1992), Fry and Gaztañaga (1993), and Hivon *et al.* (1995).

[92]Two subsamples of IRAS galaxies are analyzed by Scoccimarro *et al.* (2001) and Feldman *et al.* (2001).

[93]Recent measurements are presented in Lee *et al.* (2001), Netterfield *et al.* (2002), Halverson *et al.* (2002), Miller *et al.* (2002a), Coble *et al.* (2003), Scott *et al.* (2002) and Mason *et al.* (2002).

[94]A more careful analysis distinguishes averages across the sky from ensemble averages. By historical accident the conventional normalization replaces $2l+1$ with $2(l+1)$ in Eq. (79). Kosowsky (2002) reviews the physics of the polarization of the radiation.

[95]The physics is worked out in Peebles and Yu (1970) and Peebles (1982). Important analytic considerations are in Sunyaev and Zel'dovich (1970). The relation of the cosmic microwave background anisotropy to the cosmological parameters is explored in many papers; examples of the development of ideas include Bond (1988); Bond *et al.* (1994); Hu and Sugiyama (1996); Ratra *et al.* (1997, 1999); Zaldarriaga, Spergel, and Seljak (1997), and references therein.

of sound in the coupled baryon-radiation fluid (presenting us with still more cosmic coincidences).

The acoustic peaks in the spectrum of angular fluctuations of the 3-K cosmic microwave background radiation are the result of Fourier modes of the coupled baryon-radiation fluid that have reached maximum or minimum amplitude at decoupling. Since all Fourier components start at zero amplitude at a high redshift—in the growing density perturbation mode—this condition is

$$\int_0^{t_{\rm dec}} k c_s dt/a \simeq n\pi/2, \tag{82}$$

where c_s is the velocity of sound in the baryon-radiation fluid. Before decoupling, the mass density in radiation is greater than that of the baryons, so the velocity of sound is close to $c/\sqrt{3}$. The proper wavelength at the first acoustic peak is thus

$$\lambda_{\rm peak} \sim t_{\rm dec} \propto h^{-1}\Omega_{\rm M0}^{-1/2}. \tag{83}$$

The parameter dependence is derived from Eq. (66). The observed angle subtended by $\lambda_{\rm peak}$ is set by the angular size distance r computed from $z_{\rm eq}$ to the present [Eq. (71)]. If $\Omega_{\rm K0}=0$ or $\Omega_{\Lambda 0}=0$ the angular size distance is

$$H_0 a_0 r \simeq \begin{cases} 2\Omega_{\rm M0}^{-1} & \text{if } \Omega_{\Lambda 0}=0 \\ 2\Omega_{\rm M0}^{-0.4} & \text{if } \Omega_{\rm K0}=0. \end{cases} \tag{84}$$

If $\Lambda=0$ this expression is analytic at large $z_{\rm eq}$. The expression for $\Omega_{\rm K0}=0$ is a reasonable approximation to the numerical solution. So the angular scale of the peak varies with the matter density parameter as

$$\theta_{\rm peak} \sim z_{\rm dec}\lambda_{\rm peak}/a_0 r \propto \begin{cases} \Omega_{\rm M0}^{1/2} & \text{if } \Omega_{\Lambda 0}=0 & (85) \\ \Omega_{\rm M0}^{-0.1} & \text{if } \Omega_{\rm K0}=0. & (86) \end{cases}$$

The key point from these considerations is that the angle defined by the first peak in the fluctuation power spectrum is sensitive to $\Omega_{\rm M0}$ if $\Lambda=0$ [Eq. (85)], but not if $\Omega_{\rm K0}=0$ [Eq. (86)].[96] We have ignored the sensitivity of $z_{\rm dec}$ and $t_{\rm dec}$ to $\Omega_{\rm M0}$, but the effect is weak. More detailed computations, which are needed for a precise comparison with the data, show that the CDM model predicts that the first and largest peak of δT_l appears at multipole index $l_{\rm peak} \simeq 220\Omega_{\rm M0}^{-1/2}$ if $\Lambda=0$, and at $l_{\rm peak} \simeq 220$ if $\Omega_{\rm K0}=0$ and $0.1 \lesssim \Omega_{\rm M0} \lesssim 1$.[97]

The measured spectrum[98] peaks at $\delta T_l \sim 80\ \mu$K at $l \sim 200$, thus requiring small space curvature in the CDM model. This is the first of Efstathiou's constraints. Because of the geometric degeneracy this measurement does not yet seriously constrain $\Omega_{\rm M0}$ if $\Omega_{\rm K0}=0$.

The second constraint comes from the spectrum of temperature fluctuations on large scales, $l \lesssim 30$, where pressure gradient forces were never very important. Under the scale-invariant initial conditions discussed in Sec. III.C the Einstein–de Sitter model predicts that δT_l is nearly independent of l on large scales. A spatially flat model, with $\Omega_{\rm M0} \sim 0.3$, predicts that δT_l decreases slowly with increasing l at small l.[99] The measured spectrum is close to flat at $\delta T_l \sim 30\ \mu$K, but not well-enough constrained for a useful measure of the parameters $\Omega_{\rm M0}$ and $\Omega_{\Lambda 0}$.[100] Because of the simplicity of the physics on large angular scales, this provides the most direct and

[96]This "geometrical degeneracy" was discussed by Efstathiou and Bond (1999). Marriage (2002) presented a closer analysis of the effect. Sugiyama and Gouda (1992); Kamionkowski, Spergel, and Sugiyama (1994b); and Kamionkowski *et al.* (1994a) presented early discussions of the cosmic microwave background anisotropy in an open model.

[97]Brax *et al.* (2000) and Baccigalupi *et al.* (2000) computed the angular spectrum of the cosmic microwave background anisotropy in the dark-energy scalar field model. Doran *et al.* (2001) discussed the angular scale of the peaks in this case, and Corasaniti and Copeland (2002), Baccigalupi *et al.* (2002), and references therein, compared model predictions and observations; it is too early to draw profound conclusions about model viability, and new data are eagerly anticipated. Wasserman (2002) noted that the cosmic microwave background anisotropy data could help discriminate between different dark-energy scalar field models whose predictions do not differ significantly at a low redshift.

[98]For analyses see Knox and Page (2000); Podariu *et al.* (2001); Wang, Tegmark, and Zaldarriaga (2002); Miller *et al.* (2002b); Durrer, Novosyadlyj, and Apunevych (2003), and references therein.

[99]The physics was first demonstrated by Sachs and Wolfe (1967) and applied in the modern context by Peebles (1982). The intermediate Sachs-Wolfe effect that applies if the universe is not Einstein–de Sitter is shown in Eq. (93.26) in Peebles (1980a). This part of the Sachs-Wolfe effect receives a contribution from the low-redshift matter distribution, so cross correlating the observed large-scale cosmic microwave background anisotropy with the low-redshift matter distribution could provide another test of the world model (Boughn and Crittenden, 2001, and references therein).

[100]See, e.g., Górski *et al.* (1998). This ignores the "low" value of the cosmological quadrupole ($l=2$) moment, whose value depends on the model used to remove foreground galactic emission (see, e.g., Kogut *et al.*, 1996). Contamination due to noncosmic microwave background emission is an issue for some of the anisotropy data sets (see, e.g., de Oliveira-Costa *et al.*, 1998; Hamilton and Ganga, 2001; Mukherjee *et al.*, 2002, and references therein). Other issues that need care in such analyses include accounting for the uncertainty in the calibration of the experiment (see, e.g., Ganga *et al.*, 1997; Bridle *et al.*, 2002) and accounting for the shape of the antenna pattern (see, e.g., Wu *et al.*, 2001a; Souradeep and Ratra, 2001; Fosalba, Dore, and Bouchet, 2002).

so perhaps most reliable normalization of the CDM model power spectrum [that is, the parameter A in Eqs. (40) and (41)].

The third constraint is the baryon mass density. It affects the speed of sound c_s [Eq. (82)] in the baryon-radiation fluid prior to decoupling, and the mean free path for the radiation at $z\sim z_{\rm dec}$. These in turn affect the predicted sequence of acoustic peaks (see, e.g., Hu and Sugiyama, 1996). The detected peaks are consistent with a value for the baryon density parameter $\Omega_{\rm B0}$ in a range that includes that derived from the light element abundances [Eqs. (62)].[101] This impressive check may be much improved by the measurements of δT_l in progress.

The measurements of δT_l are consistent with a near scale-invariant power spectrum [Eq. (41) with $n\simeq 1$] with negligible contributions from gravity wave or isocurvature fluctuations (Sec. III.C.1). The 3-K cosmic microwave background temperature fluctuations show no departure from a Gaussian random process.[102] This agrees with the picture in test 10 for the nonlinear growth of structure out of Gaussian initial mass density fluctuations.

The interpretation of the cosmic microwave background temperature anisotropy measurements assumes and tests general relativity and the CDM model. One can write down other models for structure formation that place the peak of δT_l at about the observed angular scale—an example is Hu and Peebles (2000)—but we have seen none so far that seem likely to fit the present measurements of δT_l. Delayed recombination of the primeval plasma in a low-density $\Lambda=0$ CDM model can shift the peak of δT_l to the observed scale.[103] The physics is valid, but the scenario is speculative and arguably quite improbable. On the other hand, we cannot be sure that a fix of the challenges to CDM reviewed in Sec. IV.A.2 will not affect our assessments of such issues, and hence this cosmological test.

12. The mass autocorrelation function and nonbaryonic matter

If the bulk of the nonrelativistic matter, with density parameter $\Omega_{\rm M0}\sim 0.25$, were baryonic, then under adiabatic initial conditions the most immediate problem would be the strong dissipation of primeval mass density fluctuations on the scale of galaxies by diffusion of radiation through the baryons at redshifts near decoupling.[104] Galaxies could form by fragmentation of the first generation of protocluster "pancakes," as Zel'dovich (1978) proposed, but this picture is seriously challenged by the evidence that shows that the galaxies formed before clusters of galaxies.[105] In a baryonic dark-matter model we could accommodate the observations of galaxies already present at $z\sim 3$ by tilting the primeval mass fluctuation spectrum to favor large fluctuations on small scales, but that would disturb the cosmic microwave background anisotropy. The search for isocurvature initial conditions that might fit both conditions in a baryonic dark-matter model has borne no fruit so far (Peebles, 1987).

The most important point of this test is the great difficulty of reconciling the power spectra of matter and radiation without the postulate of nonbaryonic dark matter. The CDM model allows hierarchical growth of structure, from galaxies up, which is what seems to be observed, because the nonbaryonic dark matter interacts with baryons and radiation only by gravity; the dark-matter distribution is not smoothed by the dissipation of density fluctuations in the baryon-radiation fluid at redshifts $z\gtrsim z_{\rm eq}$.

As discussed in Sec. III.D, in the CDM model the small-scale part of the dark-matter power spectrum bends from the primeval scale-invariant form $P(k)\propto k$ to $P(k)\propto k^{-3}$, and the characteristic length at the break scales inversely with $\Omega_{\rm M0}$ [Eq. (42)]. Evidence of such a break in the galaxy power spectrum $P_g(k)$ has been known for more than a decade;[106] it is consistent with a value of $\Omega_{\rm M0}$ in the range of Eq. (59).

13. The gravitational inverse-square law

The inverse-square law for gravity determines the relation between the mass distribution and the gravitationally driven peculiar velocities that enter estimates of the matter density parameter $\Omega_{\rm M0}$. The peculiar velocities also figure in the evolution of the mass distribution, and

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[101]The $\Omega_{\rm B0}h^2$ values estimated from the cosmic microwave background anisotropy measured by Stompor *et al.* (2001), Netterfield *et al.* (2002), and Pryke *et al.* (2002) are more consistent with the higher, deuterium-based, Burles *et al.* (2001) range in Eqs. (62).

[102]Colley, Gott, and Park (1996) presented an early discussion of large angular scale data; more recent discussions can be found in Mukherjee, Hobson, and Lasenby (2000), Phillips and Kogut (2001), and Komatsu *et al.* (2002). Degree and subdegree angular scale anisotropy data were studied in Park *et al.* (2001), Wu *et al.* (2001b), Shandarin *et al.* (2002), and Polenta *et al.* (2002).

[103]The model in Peebles, Seager, and Hu (2000) assumes stellar ionizing radiation at $z\sim 1000$ produces recombination Lyman-α photons. These resonance photons promote photoionization from the $n=2$ level of atomic hydrogen. That allows delayed recombination with a rapid transition to neutral atomic hydrogen, as is required to get the shape of δT_l about right.

[104]Early analyses of this effect are in Peebles (1965); and Silk (1967, 1968).

[105]For example, our Milky Way galaxy is in the Local Group, which seems to be just forming, because the time for a group member to cross the Local Group is comparable to the Hubble time. The Local Group is on the outskirts of the concentration of galaxies around the Virgo cluster. We and neighboring galaxies are moving away from the cluster, but at about 80% of the mean Hubble flow, as if the local mass concentration were slowing the local expansion. That is, our galaxy, which is old, is starting to cluster with other galaxies, in a "bottom up" hierarchical growth of structure, as opposed to the "top down" evolution of the pancake picture.

[106]The first good evidence was discussed in Efstathiou *et al.* (1990); for recent examples see Sutherland *et al.* (1999), Percival *et al.* (2001), and Dodelson *et al.* (2002).

hence the relation between the present mass fluctuation spectrum and the spectrum of cosmic microwave background temperature fluctuations imprinted at a redshift $z \sim 1000$. We are starting to see demanding tests of both aspects of the inverse-square law.

We have a reasonably well-checked set of measurements of the apparent value of Ω_{M0} on scales ranging from 100 kpc to 10 Mpc (as reviewed under test 7). Most agree with a constant value of the apparent Ω_{M0}, within a factor of approximately 3. This is not the precision one would like, but the subject has been under discussion for a long time, and, we believe, is now fairly reliably understood, within that same factor of approximately 3. If galaxies were biased tracers of mass one might have expected to have seen that Ω_{M0} increases with an increasing length scale, since the increasing scale includes the outer parts of extended massive halos. Maybe this is masked by a gravitational force law that decreases more rapidly than the inverse-square law at a large distance. But the much more straightforward reading is that the slow variation of Ω_{M0} sampled over two orders of magnitude in the length scale agrees with the evidence from tests 7 to 10 that galaxies are useful mass tracers, and that the inverse-square law therefore is a useful approximation on these scales.

The toy model in Eq. (57) illustrates how a failure of the inverse-square law would affect the evolution of the shape of the mass fluctuation power spectrum $P(k,t)$ as a function of the comoving wave number k, in linear perturbation theory. This is tested by the measurements of the mass and cosmic microwave background temperature fluctuation power spectra. The galaxy power spectrum $P_g(k)$ varies with wave number at $k \sim 0.1h$ Mpc^{-1} about as expected under the assumptions that the mass distribution grew by gravity out of adiabatic scale-invariant initial conditions, the mass is dominated by dark matter that does not suffer radiation drag at a high redshift, the galaxies are useful tracers of the present mass distribution, the matter density parameter is $\Omega_{M0} \sim 0.3$, and, of course, the evolution is adequately described by conventional physics (Hamilton and Tegmark, 2002, and references therein). If the inverse-square law were significantly wrong at $k \sim 0.1h$ Mpc^{-1}, the near scale-invariant form would have to be an accidental effect of some failure in this rather long list of assumptions. This seems unlikely, but a check certainly is desirable, and one is available from the cosmic microwave background anisotropy measurements. These are consistent with near scale-invariant initial conditions applied at a redshift $z \sim 1000$. This preliminary check on the effect of the gravitational inverse-square law applied on cosmological length scales and back to a redshift $z \sim 1000$ will be improved by a better understanding of the effect on δT_l of primeval tensor perturbations to spacetime, and of the dynamical response of the dark-energy distribution to the large-scale mass distribution.

Another aspect of this check is the comparison of values of the large-scale rms fluctuations in the present distributions of mass and the cosmic microwave background radiation. The latter is largely set at decoupling, after which the former grows by a factor of about 10^3 to the present epoch, in the standard relativistic cosmological model. If space curvature is negligible the growth factor agrees with the observations to about 30%, assuming galaxies trace mass. In a low-density universe with $\Lambda = 0$ the standard model requires either that mass is more smoothly distributed than are galaxies, $\delta N/N \sim 3\,\delta M/M$, or that the gravitational growth factor since decoupling is a factor of 3 off the predicted factor ~ 1000; this factor of 3 is about as large a deviation from unity as is viable. We are not proposing this as an interpretation of the data, rather we are impressed by the modest size of the allowed adjustment to the inverse-square law.

C. The state of the cosmological tests

Precision cosmology is not very interesting if it is based on faulty physics or astronomy. That is why we have emphasized the tests of the standard gravity physics and structure-formation model, in addition to checks of consistency of measures based on different aspects of astronomy.

There are now five main lines of evidence that significantly constrain the value of Ω_{M0} to the range of Eq. (59): the redshift-magnitude relation (test 4), gravitational dynamics and weak lensing (test 7), the baryon mass fraction in clusters of galaxies (test 8), the abundance of clusters as a function of mass and redshift (test 9), and the large-scale galaxy distribution (test 12). There are indications for larger values of Ω_{M0}, from analyses of the rate of strong lensing of quasars by foreground galaxies (test 6) and some analyses of large-scale flows (test 7), though we know of no well-developed line of evidence that points to the Einstein–de Sitter value $\Omega_{M0} = 1$. Each of these measures of Ω_{M0} may suffer from systematic errors: we must bear in mind the Tantalus principle mentioned in Sec. I.A, and we have to remember that the interpretations could be corrupted by a failure of standard physics. But the general pattern of results from a considerable variety of independent approaches seems so close to consistent as to be persuasive. Thus we conclude that there is a well-checked scientific case for the proposition that the measures of the mean mass density of matter in forms capable of clustering are physically meaningful, and that the mass density parameter almost certainly is in the range $0.15 \lesssim \Omega_{M0} \lesssim 0.4$.

In standard cosmology the masses of the galaxies are dominated by dark matter, with a mass density parameter $\Omega_{DM0} \sim 0.2$ that is not baryonic (or acts as such). We do not have direct evidence of a laboratory detection; this is based on two indirect lines of argument. First, the successful model for the origin of the light elements (test 2) requires a baryon density $\Omega_{B0} \sim 0.05$. It is difficult to see how to reconcile a mass density this small with the mass estimates from dynamics and lensing; the hypothesis that Ω_{M0} is dominated by nonbaryonic matter allows us to account for the difference. Second, the nonbaryonic matter allows us to reconcile the theory of the anisotropy of the cosmic microwave background radiation with the distributions of galaxies, and groups and clusters of galaxies, and the presence of galaxies at $z \sim 3$ (tests 11 and 12). This interpretation requires a value for Ω_{B0} that is in line with test 2. The consistency is impressive. But the case is not yet as convincing as is the larger network of evidence that shows that Ω_{M0} is well below unity.

The subject of this review is Einstein's cosmological constant Λ, or its equivalent in dark energy. The evidence for detection of Λ by the redshift-magnitude relation for type-Ia supernovae is checked by the angular distribution of the 3-K cosmic microwave background temperature together with the constraints on Ω_{M0}. This certainly makes a serious case for dark energy. But we keep accounts by the number of significant independent checks, and by this reckoning the case is not yet as strong as the one for nonbaryonic dark matter.

V. CONCLUDING REMARKS

We cannot demonstrate that there does not exist some other physics, applied to some other cosmology, that equally well agrees with the cosmological tests. The same applies to the whole enterprise of physical science, of course. Parts of physics are so densely checked that they are quite convincing approximations to reality. The web of tests is much less dense in cosmology, but, as we have tried to demonstrate, by no means negligible, and growing tighter.

A decade ago there was not much discussion on how to test general relativity theory on the scales of cosmology. That was in part because the theory seemed so logically compelling, and certainly in part also because there was not much evidence to work with. The empirical situation is much better now. We mentioned two tests, namely, of the relativistic active gravitational mass density and the gravitational inverse-square law. The consistency of constraints on Ω_{M0} from dynamics and the redshift-magnitude relation adds a test of the effect of space curvature on the expansion rate. These tests are developing; they will be greatly improved by work in progress. There certainly may be surprises in the gravity physics of cosmology at redshifts $z \lesssim 10^{10}$, but it is already clear that if so the surprises will be subtly hidden.

A decade ago the direction the theory of large-scale structure formation would take was not at all clear. Now the simple CDM model has had enough success that there is good reason to expect the standard model ten years from now will resemble CDM. We have listed challenges to this structure-formation model. Some may merely be a result of the complications of interpreting the theory and observations. Others may prove to be real and drive adjustments of the model. Fixes certainly will include one element from the ideas of structure formation that were under debate a decade ago: explosions that rearrange matter in ways that are difficult to compute. Fixes may also include primeval isocurvature departures from homogeneity, such as in spacetime curvature fluctuations frozen in during inflation, and perhaps in new cosmic fields. It would not be surprising if cosmic field defects, that have such a good pedigree from particle physics, also find a role in structure formation. And a central point to bear in mind is that fixes, which do not seem unlikely, could affect the cosmological tests.

A decade ago we had significant results from the cosmological tests. For example, estimates of the product $H_0 t_0$ suggested we might need positive Λ, though the precision was not quite adequate for a convincing case. That is still so; the community will be watching for further advances. We had fairly good constraints on Ω_{B0} from the theory of the origin of the light elements. The abundance measurements are improving; an important recent development is the detection of deuterium in gas clouds at redshifts $z \sim 3$. Ten years ago we had useful estimates of masses from peculiar motions on relatively small scales, but also more mixed messages from larger scales. The story seems more coherent now, but there is still room for improved consistency. We had a case for nonbaryonic dark matter, from the constraint on Ω_{B0} and from the CDM model for structure formation. The case is tighter now, most notably due to the successful fit of the CDM model prediction to the measurements of the power spectrum of the 3-K cosmic microwave background radiation anisotropy. In 1990 there were believable observations of galaxies identified as radio sources at $z \sim 3$. Now the distributions of galaxies and the intergalactic medium at $z \sim 3$ are mapped out in impressive detail, and we are seeing the development of a semiempirical picture of galaxy formation and evolution. Perhaps that will lead us back to the old dream of using galaxies as markers for cosmological tests.

Until recently it made sense to consider the constraints on one or two of the parameters of the cosmology while holding all the rest at "reasonable" values. That helps us understand what the measurements are probing; it is the path we have followed in Sec. IV.B. But the modern and very sensible trend is to consider joint fits of large numbers of parameters to the full suite of observations.[107] This includes a measure of the biasing or antibiasing of the distribution of galaxies relative to mass, rather than our qualitative argument that one usefully approximates the other. In a fully satisfactory cosmological test the parameter set will also include parametrized departures from standard physics extrapolated to the scales of cosmology.

Community responses to advances in empirical evidence are not always close to linear. The popularity of the Einstein–de Sitter model continued longer into the 1990s than seems logical to us, and the switch to the now standard ΛCDM cosmological model—with flat space sections, nonbaryonic cold dark matter, and dark energy—arguably is more abrupt than warranted by the

[107]Examples are Cole *et al.* (1997), Jenkins *et al.* (1998), Lineweaver (2001), and Percival *et al.* (2002).

advances in the evidence. Our review leads us to conclude that there is now a good scientific case arguing that the matter density parameter is $\Omega_{M0} \sim 0.25$, and a fairly good case that about three-quarters of that is not baryonic. The cases for dark energy and for the ΛCDM model are significant, too, though obscured by observational issues of whether we have an adequate picture of structure formation. But we expect that rapid advances in the observations of structure formation will soon dissipate these clouds, and, considering the record, likely reveal new clouds over the standard model for cosmology a decade from now.

A decade ago the high-energy-physics community had a well-defined challenge to show why the dark-energy density vanishes. Now there seems to be both a new challenge and clue: determine why the dark-energy density is exceedingly small but not zero. The present state of ideas can be compared with the state of research on structure formation a decade ago: in both situations there are many lines of thought but not a clear picture of which is the best direction to take. The big difference is that a decade ago we could be reasonably sure that observations in progress would guide us to a better understanding of how structure formed. Untangling the physics of dark matter and dark energy and their role in gravity physics is a much more subtle challenge, but, we hope, will be guided by advances in the exploration of the phenomenology. Perhaps in another ten years this will include detecting the evolution of the dark energy, and detecting the gravitational response of the dark-energy distribution to the large-scale mass distribution. There may be three unrelated phenomena to deal with: dark energy, dark matter, and a vanishing sum of zero-point energies and whatever goes with them. Or the phenomena may be related. Because our only evidence of dark matter and dark energy is from their gravity, it is a natural and efficient first step to suppose that their properties are as simple as allowed by the phenomenology. However, it makes sense to watch for hints of more complex physics within the dark sector.

The past eight decades have seen steady advances in the technology used for the cosmological tests, from telescopes to computers; advances in the theoretical concepts underlying the tests; and progress through the learning curves on applying the concepts and technology. We see the results: the basis for cosmology is much firmer than it was a decade ago. And the basis surely will be a lot more solid a decade from now.

Einstein's cosmological constant, and the modern variant, dark energy, have figured in a broad range of topics under discussion in physics and astronomy, in at least some circles, for much of the past eight decades. Many of these issues undoubtedly have been discovered more than once. But in our experience such ideas tend to persist for a long time at low visibility and sometimes low fidelity. Thus the community has been very well prepared for the present evidence for detection of dark energy. And for the same reason we believe that dark energy, whether constant, or rolling toward zero, or maybe even increasing, still will be an active topic of research, in at least some circles, a decade from now, whatever the outcome of the present work on the cosmological tests. Though this much is clear, we see no basis for a prediction of whether the standard cosmology a decade from now will be a straightforward elaboration of ΛCDM, or whether there will be more substantial changes of direction.

Note added in proof

The measurements of the angular distribution of the 3-K cosmic microwave background radiation by the Wilkinson Microwave Anisotropy Probe (WMAP) were released (in Bennett *et al.*, 2003; Spergel *et al.*, 2003, and references therein) after this review was completed. It is appropriate to comment on how the results of this superb experiment have changed our assessment of the cosmological tests.

The discussion in Sec. IV.C led us to conclude that the case for the detection of dark energy is not as compelling as the case for dark matter, because there are fewer cross checks. WMAP changes that. The ΛCDM model gives an excellent fit to the WMAP measurements. The parameters required by this fit, including density parameters $0.19 \leqslant \Omega_{M0} \leqslant 0.35$, $0.65 \leqslant \Omega_{\Lambda 0} \leqslant 0.81$, and $-0.02 \leqslant \Omega_{K0} \leqslant 0.06$ (all at two standard deviations), are in good agreement with other constraints (as summarized in Sec. IV). In particular, a pre-WMAP survey of the constraints on Ω_{M0} from a combination of the dynamical, baryon fraction, power spectrum, weak lensing, and cluster abundance measurements indicates $0.2 \lesssim \Omega_{M0} \lesssim 0.35$ (at two standard deviations, Chen and Ratra, 2003b), in striking accord with the WMAP estimate. The fit to the WMAP measurements and the overall consistency of parameter constraints is strong evidence that the ΛCDM model is a good approximation to reality. This evidence increases the weight of parameter estimates that depend on the ΛCDM model. And the model fit to the WMAP measurements requires the presence of dark energy, provided the Hubble constant is within acceptable distance of the astronomical measurements. This is distinct from the line of argument for $\Omega_{\Lambda 0} \sim 0.7$ described in Sec. IV. It provides the wanted cross check that makes a convincing case for the detection of dark energy.

Issues remain. The ΛCDM fit to WMAP indicates the density parameter in baryons is $0.021 \leqslant \Omega_{B0} h^2 \leqslant 0.024$ (at two standard deviations), consistent with what is indicated by the standard Big Bang nucleosynthesis model and measurements of the primeval deuterium abundance [Eq. (62) and footnote 101]. To be resolved are the somewhat different estimates of $\Omega_{B0} h^2$ from the helium and lithium abundances. The temperature anisotropy autocorrelation function is consistent with zero at angular separations greater than about 60° (consistent with the earlier but less emphatic COBE result mentioned in footnote 100), and seems not likely to be consistent with the ΛCDM prediction. Maybe this is an unlikely statistical fluctuation. Or maybe it is telling us about the physics of the dark energy.

We take this opportunity to refer to Padmanabhan's (2002) recently completed review of the cosmological constant, with particular emphasis on possible resolutions of the physicists' cosmological constant problem and the physics of the resulting dark energy models.

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APPENDIX: RECENT DARK-ENERGY SCALAR FIELD RESEARCH

Many dark-energy models are characterized by attractor or tracker behavior. Another goal is also to design the model so that the field energy density is subdominant at a high redshift, where it is not wanted, and dominant at a low redshift, where that is what seems to be observed.

In the simplest such scalar field models the action has a conventionally normalized scalar field kinetic and spatial gradient term, and the real scalar field is coupled only to itself and gravity. Then the scalar field part of the model is fully characterized by the scalar field potential (along with some broad constraints on the initial conditions for the field, if the attractor behavior is realized). The inverse-power-law potential model is described in Secs. II.C and III.E, and used in some of the cosmological tests in Sec. IV. Here we list other scalar field potentials now under discussion for these minimal dark-energy models, modifications of the kinetic part of the action, possible guidance from high-energy particle physics ideas, and constraints on these ideas from cosmological observations in the context of dark-energy models.

To those of us not active in this field the models may seem baroque in their complexity, but that may be the way of Nature. And as we accumulate more and better data it will be possible to test and constrain these models.

This is an active area of research, with frequent introduction of new models, so our discussion will be somewhat sketchy.[108] We limit discussion to a cosmological model with space sections that are flat, thanks to the presence of the dark-energy density.

As mentioned in Sec III.E, the simplest exponential potential scalar field model is unacceptable because it cannot produce the wanted transition from subdominant to dominant energy density. Ratra and Peebles (1988) considered more complex potentials, such as powers of linear combinations of field Φ exponentials. Related models are under active investigation. These include potentials that are powers of $\cosh(\Phi)$ and/or $\sinh(\Phi)$.[109] Sahni and Wang (2000) presented a detailed discussion of a specific example, $V(\Phi)\propto(\cosh\lambda\Phi-1)^p$, where λ and p are constants, emphasizing that this potential interpolates from $V\propto e^{-p\lambda\Phi}$ to $V\propto(\lambda\Phi)^{2p}$ as $|\lambda\Phi|$ increases, thus preserving some of the desirable properties of the simplest exponential potential case. de la Macorra and Piccinelli (2000) considered potentials that are exponentials of more complicated functions of Φ, such as Φ^2 and e^Φ. Skordis and Albrecht (2002) discussed a model with $V(\Phi)\propto[1+(\Phi-A)^2]\exp[-\Phi\sqrt{q/2}]$, while Dodelson, Kaplinghat, and Stewart (2000) studied a potential $V(\Phi)\propto[1+A\sin(B\Phi)]\exp[-\Phi\sqrt{q/2}]$, and Ng, Nunes, and Rosati (2001) considered a model with $V(\Phi)\propto\Phi^A\exp[B\Phi^C]$, a simple example of a class of supergravity-inspired potentials studied by Copeland, Nunes, and Rosati (2000). Here A, B, and C are parameters. Steinhardt *et al.* (1999) considered more complicated functions of inverse powers of Φ, such as $V(\Phi)\propto\exp(1/\Phi)$ and linear combinations of inverse powers of Φ. Brax and Martin (2000) considered a supergravity-inspired generalization with $V(\Phi)=\kappa\Phi^{-\alpha}e^{\Phi^2/2}$.[110]

It will be quite a challenge to select from this wide range of functional forms for the potential ones that have particular theoretical merit and some chance of being observationally acceptable.

Following Dolgov (1983), there have been discussions of nonminimally coupled dark-energy scalar field models.[111] These Jordan-Brans-Dicke-type models have an explicit coupling between the Ricci scalar and a function of the scalar field. This causes the effective gravita-

[108]In particular, we do not discuss variable mass term models, complex or multiple scalar field models, or repulsive matter. We also omit nonscalar field aspects of brane models, Kaluza-Klein models, bimetric theories of gravitation, quantum-mechanical running of the cosmological constant, the Chaplygin gas, and the superstring tachyon.

[109]For examples, see Chimento and Jakubi (1996), Starobinsky (1998), Kruger and Norbury (2000), Di Pietro and Demaret (2001), Ureña-López and Matos (2000), González-Díaz (2000), and Johri (2002).

[110]For still more examples, see Green and Lidsey (2000); Barreiro, Copeland, and Nunes (2000); Rubano and Scudellaro (2002); Sen and Sethi (2002), and references therein.

[111]See Uzan (1999); Chiba (1999); Amendola (1999, 2000); Perrotta, Baccigalupi, and Matarrese (2000); Bartolo and Pietroni (2000); Bertolami and Martins (2000); Fujii (2000); Faraoni (2000); Baccigalupi *et al.* (2000); Chen, Scherrer, and Steigman (2001), and references therein.

tional "constant" G (in units where masses are constant) to vary with time, which may be an observationally interesting effect and useful constraint.[112]

In yet another approach, people have considered modifying the form of the dark energy scalar field kinetic and spatial gradient term in the action.[113]

In Sec. III.E we discussed the idea that at the end of inflation the dark-energy scalar field potential might patch on to the part of the scalar field potential responsible for inflation. In an inverse-power-law model for the dark-energy potential function this requires an abrupt drop in $V(\Phi)$ at the end of inflation. More sophisticated models, now dubbed quintessential inflation, attempt to smooth out this drop by constructing scalar field potentials that interpolate smoothly between the part responsible for inflation and the low-redshift, dark-energy, part. These models assume either minimally or nonminimally (Jordan-Brans-Dicke) coupled scalar fields.[114]

The dark-energy scalar field models we have reviewed here are meant to be classical, effective descriptions of what might result from a more fundamental quantum-mechanical theory. The effective dark-energy scalar field is coupled to itself and gravity, and is supposed to be coupled to the other fields in the universe only by gravity. This might be what Nature chooses, but we lack an understanding of why the coupling of dark and ordinary fields that are allowed by the symmetries are not present, or have coupling strengths that are well below what might be expected by naive dimensional analysis (e.g., Kolda and Lyth, 1999). A satisfactory solution remains elusive. Perhaps this is not unexpected, because it likely requires a proper understanding of how to reconcile general relativity and quantum mechanics.

We turn now to scalar field dark-energy models that arguably are inspired by particle physics. Inverse-power-law scalar field potentials are generated nonperturbatively in models of dynamical supersymmetry breaking. In supersymmetric non-Abelian gauge theories, the resulting scalar field potential may be viewed as being generated by instantons, the potential being proportional to a power of $e^{-1/g^2(\Phi)}$, where $g(\Phi)$ is the gauge coupling constant which evolves logarithmically with the scalar field through the renormalization-group equation. Depending on the parameters of the model, an inverse-power-law scalar field potential can result. This mechanism may be embedded in supergravity and superstring/M theory models.[115] This has not yet led to a model that might be compared to the observations.[116]

In the model considered by Weiss (1987), Frieman *et al.* (1995), and others the dark-energy field potential is of the form $V(\Phi)=M^4[\cos(\Phi/f)+1]$, where M and f are mass scales and the mass of the inhomogeneous scalar field fluctuation $\sim M^2/f$ is on the order of the present value of the Hubble parameter. For discussions of how this model might be more firmly placed on a particle physics foundation see Choi (1999); Kim (2000); Nomura, Watari, and Yanagida (2000); and Barr and Seckel (2001).

There has been much recent interest in the idea of inflation in the brane scenario. Dvali and Tye (1999) note that the potential of the scalar field which describes the relative separation between branes can be of a form that leads to inflation, and will include some inverse-power-law scalar field terms.[117] It will be interesting to learn whether these considerations can lead to a viable dark-energy scalar field model. Brane models allow for a number of other possibilities for dark-energy scalar fields,[118] but it is too early to decide whether any of these options give rise to observationally acceptable dark-energy scalar field models.

Building on earlier work,[119] and Hellerman, Kaloper, and Susskind (2001) and Fischler *et al.* (2001) noted that dark-energy scalar field cosmological models have future event horizons characteristic of the de Sitter model. This means some events have causal futures that do not share any common events. In these dark-energy scalar

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[112]Nucleosynthesis constraints on these and related models are discussed by Arai, Hashimoto, and Fukui (1987); Etoh *et al.* (1997); Perrotta *et al.* (2000); Chen *et al.* (2001); and Yahiro *et al.* (2002).

[113]See Fujii and Nishioka (1990); Chiba, Okabe, and Yamaguchi (2000); Armendariz-Picon, Mukhanov, and Steinhardt (2001); Hebecker and Wetterich (2001), and references therein.

[114]Early work includes Frewin and Lidsey (1993), Spokoiny (1993), Joyce and Prokopec (1998), and Peebles and Vilenkin (1999); more recent discussions can be found in Dimopoulos and Valle (2001); Kaganovich (2001); Huey and Lidsey (2001); Majumdar (2001); Sahni, Sami, and Souradeep (2002).

[115]In the superstring/M theory case, since the coupling constant is an exponential function of the dilaton scalar field, the resulting potential is usually not of the inverse-power-law form. However, it is perhaps not unreasonable to think that after the dilaton has been stabilized, it or one of the other scalar fields in superstring/M theory might be able to play the role of dark energy. Townsend (2001) and Gasperini, Piazza, and Veneziano (2002) considered other ways of using the dilaton as a dark-energy scalar field candidate.

[116]See, Davis, Dine, and Seiberg (1983), and Rossi and Veneziano (1984) for early discussions of supersymmetry breaking, and Quevedo (1996), Dine (1996), Peskin (1997), and Giudice and Rattazzi (1999) for reviews. Applications of dynamical nonperturbative supersymmetry breaking directly relevant to the dark-energy scalar field model are discussed in Binétruy (1999); Masiero, Pietroni, and Rosati (2000); Copeland *et al.* (2000); Brax, Martin, and Riazuelo (2001); de la Macorra and Stephan-Otto (2001), and references therein.

[117]Recent discussions of this setup include Halyo (2001b); Shiu and Tye (2001); Burgess *et al.* (2002); Kyae and Shafi (2002); García-Bellido, Rabadán, and Zamora (2002); Blumenhagen *et al.* (2002); Dasgupta *et al.* (2002).

[118]See Uzawa and Soda (2001); Huey and Lidsey (2001); Majumdar (2001); Chen and Lin (2002); Mizuno and Maeda (2001); Myung (2001); Steinhardt and Turok (2002), and references therein.

[119]See Bousso (2000), Banks and Fischler (2001), and Maldacena and Nuñez (2001), and references therein.

field models, some correlations are therefore unmeasurable, which destroys the observational meaning of the S matrix. This indicates that it is not straightforward to bring superstring/M theory into consistency with dark-energy models in which the expansion of the universe is accelerating.[120]

At the time of writing, while there has been much work, it appears that the dark-energy scalar field scenario still lacks a firm, theoretical foundation based on high-energy physics. While this is a significant drawback, the recent flurry of activity prompted by developments in superstring/M and brane theories appears to hold significant promise for shedding light on dark energy. Whether this happens before the observations rule out or "confirm" dark energy is an intriguing question.

[120]Other early discussions of this issue may be found in He (2001), Moffat (2001), Halyo (2001a), Kolda and Lahneman (2001), and Deffayet, Dvali, and Gabadadze (2002). The more recent literature may be accessed from Larsen, van der Schaar, and Leigh (2002), and Medved (2002).

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