

Astrophysics and cosmology notes

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March 8, 2020

Introduction and relevant material

These are the (yet to be) revised notes for the course “Fundamentals of Astrophysics and Cosmology” held by professor Sabino Matarrese in fall 2019 at the university of Padua.

They are based on the notes I took during lectures, complemented with notes from the previous years.

They will be revised by the professor in the future, as of yet they have not.

The exam is a traditional oral exam, there are fixed dates but they do not matter: on an individual basis we should write an email to the professor to set a date and time.

Material There is a dropbox folder with notes by a student from the previous years [[Pac18](#)] and handwritten notes by the professor.

There are many good textbooks, for example “*Cosmology*” by Lucchin and Coles [[LC02](#)].

Chapter 1

Cosmography

Thursday
2019-10-3

1.1 The cosmological principle

The basis for the modern treatment of cosmology is the **cosmological principle** (or Copernican principle): roughly stated, it is “*we do not occupy a special, atypical position in the universe*”. We will discuss the validity of this later in this section. It is extremely useful to make such an assumption since it endows our model of the universe with a great deal of symmetry, which makes its mathematical treatment manageable.

A more formal statement is:

Proposition 1.1.1 (Cosmological principle). *Every comoving observer observes the Universe around them at a fixed time in their reference frame as being homogeneous and isotropic.*

Comoving means moving coherently with the absolute reference frame, which is defined as the rest frame of the cosmic fluid, which determines the geometry of the universe.

When we observe the Cosmic Microwave Background (CMB) we see that we are surrounded by radiation distributed like a blackbody of temperature $T_{\text{CMB}} \approx 2.725 \text{ K}$ [Fix09].

This radiation is not uniformly distributed in the sky: we see a *dipole modulation* of around $\Delta T \approx 3.4 \text{ mK}$ [Col18]. This is due to the Doppler effect: the Solar System is *not* comoving with respect to the CMB.

In fact, we can measure the *peculiar velocity* of the Solar System this way: it comes out to be around $c\Delta T/T_{\text{CMB}} \approx 370 \text{ km/s}$. This cannot be explained by the movement of the Sun through the galaxy, nor by the movement of the galaxy through the Local Group: the Local group is actually moving with respect to the absolute reference. In fact, the velocities of the Sun with respect to the Local Group and the velocity of the Local Group with respect to the CMB are almost directed in opposite directions, so the velocity of the LG with respect to the CMB can be measured to be $\approx 620 \text{ km/s}$ [Col18, Table 3].

So, the absolute reference frame can be experimentally defined as the frame of the observer who sees the CMB with zero dipole moment.

The CMB has anisotropies of the order of $20 \mu\text{K}$ (root mean square) [Wri03] at higher order multipole moments: it is uniform to approximately 1 part in 10^5 .

The word “time” in **fixed time** refers to the proper time of a comoving observer, which is called *cosmic time*.

Homogeneity means that the characteristics of the universe as observed from a point are the same as they would be as observed from any other point.

Isotropy means that the characteristics of the universe as observed in a certain direction are the same as they would be if they were observed in another direction.

On the validity of the cosmological principle The principle is expected to hold only on very large¹ scales: at small scales we see structures, such as galaxies or our Solar System, so we surely do not have homogeneity.

How can we talk about homogeneity if we can only look at the universe from a single point? We must *assume* that any other observer would also see isotropy as we do.

Isotropy around every point implies homogeneity. We observe isotropy, we assume homogeneity. So, we can go from isotropy to homogeneity under the assumption that we are typical observers.

In the end, this assumption has to be made before any cosmological study starts: it might not be completely correct, but it allows us to make falsifiable predictions, so we shall keep it until the models it allows us to create do not match observations anymore.

1.2 The geometry of spacetime

The best description for gravity we have so far is given by the general theory of relativity (GR). In it, spacetime is modelled as a 3+1-dimensional semi-Riemannian manifold with a line element which is generally given by the expression $ds^2 = g_{ab} dx^a dx^b$ and which.

Latin indices can take values from 0 to 3, and we adopt the “mostly minus” metric signature.

Such a spacetime can have up to $4(4 + 1)/2 = 10$ global continuous symmetries, which can be classified into:

- (a) 1 time translation;
- (b) 3 Lorentz boosts;
- (c) 3 spatial translations;
- (d) 3 spatial rotations.

The metric of Minkowski spacetime, which has all of these 10 symmetries, reads:

$$ds^2 = c^2 dt^2 - d\vec{x}^2 \quad (\text{cartesian coordinates}) \quad (1.1a)$$

$$= c^2 dt^2 - dr^2 - r^2 d\Omega^2 \quad (\text{spherical coordinates}), \quad (1.1b)$$

where $r = |\vec{x}|$, while θ and φ are the spherical angles and $d\Omega^2 = d\theta^2 + \sin^2 \theta d\varphi^2$.

In cosmology we lose time translation symmetry, since the universe is expanding, and Lorentz boost symmetry, since as we saw in the previous section we can tell at which speed we are moving with respect to the CMB.

¹Larger than $260h^{-1}\text{Mpc} \approx 380\text{Mpc}$ [YBK10], where $h \approx 0.68$ [Col16]: this will be made clearer in later sections, but is meant to give an idea of the length scales involved.

We keep the purely spatial symmetries: so, our description of spacetime will be as a 3+1-dimensional manifold which, if we fix the temporal coordinate for a comoving observer, reduces to a *maximally symmetric* 3-dimensional space — for 3 dimensions the maximum number of symmetries is $3(3+1)/2 = 6$, which correspond to (c) and (d).

It can be shown that the most general form of a metric satisfying these conditions is the **Friedmann-Lemaître-Robertson-Walker line element**,² which, in the comoving frame, reads

$$ds^2 = c^2 dt^2 - a^2(t) \left(\frac{dr^2}{1 - kr^2} + r^2 d\Omega^2 \right). \quad (1.2)$$

The coordinate r does not have the dimensions of a length: we choose our variables so that $a(t)$ is a length, while r is adimensional.³

The parameter $a(t)$ is called the *scale factor*: varying it amounts to rescaling all of space. It has the dimensions of a length, and it depends on the cosmic time t .

The parameter k is a constant describing the spatial curvature, which can always be normalized to ± 1 or 0. Universes with:

- $k = 1$ are called closed universes;
- $k = -1$ are called open universes;
- $k = 0$ are called flat universes.

We can choose the normalization of k , but its sign is a constant. Positive values correspond to $k = 1$, negative values to $k = -1$. When computing probability distributions for k , we must not consider it to be a discrete variable but a continuous variable instead: so, the set of flat universes with $k = 0$ has zero measure, and thus zero probability with any probability density function.

1.2.1 A bidimensional example

We consider surfaces: these are the simplest manifolds which can have intrinsic curvature.

Intrinsic curvature is described by the Riemann tensor, which has 20 independent components in 4D and 6 in 3D, so it is difficult to visualize.

On the other hand, in 2D has only 1 independent component: $R_{1212} = R \det g_{ab}$, where R is the scalar curvature.

R has an immediate geometric interpretation: it is equal to $2/(R_1 R_2)$, where R_i are the radii of the osculating circles at the point; it is positive if the circles are in the same direction, as they are for a sphere, and negative if they are in different directions, as they

²This is sometimes called just “Robertson-Walker”, or RW, or FLRW.

³This is a matter of convention: if we normalize $|k| = 1$ then we will have $0 \leq r < 1$; we could also let k be arbitrarily large; then r would also be.

are for a hyperboloid.

If we have a direction in which the surface is flat, such as in a cylinder, then $R = 0$: it is intrinsically flat.

The metric for a cartesian flat 2D plane is $dl^2 = a^2(dr^2 + r^2 d\theta^2)$. The constant a is included since we want r to be adimensional.

The metric for the surface of a sphere is: $dl^2 = a^2(d\theta^2 + \sin^2 \theta d\varphi^2)$, where $a^2 = R^2$, the square radius of the sphere.

The metric for the surface of a hyperboloid is: $dl^2 = a^2(d\theta^2 + \sinh^2 \theta d\varphi^2)$, therefore the only difference is that trigonometric functions become hyperbolic ones.

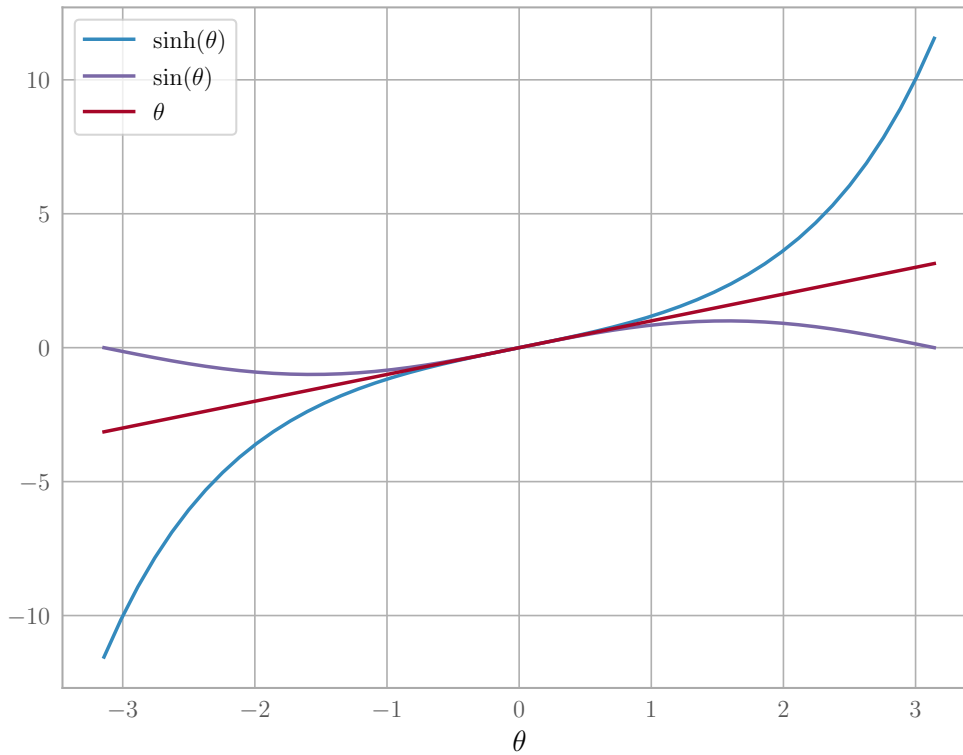


Figure 1.1: Plot of the functions $\sin(\theta)$ and $\sinh(\theta)$ in the interval $-\pi \leq \theta \leq \pi$.

Do note that for the sphere $\theta \in [-\pi, \pi]$ while for the hyperboloid we can in principle have $\theta \in \mathbb{R}$: this is indicative of the fact that the sphere is bounded, while the hyperboloid is infinite,

For both of these, let us define the variable: $r = \sin \theta$ in the spherical case, and $r = \sinh \theta$ in the hyperbolic case.

As we change variable we do the following manipulation for the sphere:

$$dr^2 = \left(\frac{dr}{d\theta} \right)^2 d\theta^2 = \cos^2 \theta d\theta^2 \quad (1.3a)$$

$$\implies d\theta^2 = \frac{dr^2}{\cos^2 \theta} = \frac{dr^2}{1 - r^2}, \quad (1.3b)$$

and similarly for the hyperboloid, except for the fact that in that case $\cosh^2 \theta = 1 + \sinh^2 \theta = 1 + r^2$.

So, the line elements become respectively:

$$dl^2_{\text{sphere}} = a^2 \left(\frac{dr^2}{1 - r^2} + r^2 d\varphi^2 \right) \quad (1.4a)$$

$$dl^2_{\text{hyperboloid}} = a^2 \left(\frac{dr^2}{1 + r^2} + r^2 d\varphi^2 \right). \quad (1.4b)$$

We have a striking similarity to the Robertson-Walker metric: we only need to make the substitution $d\varphi \rightarrow d\Omega$ in order to recover it.

We can also work backwards and rewrite the RW line element in sphere- or hyperboloid-like coordinates:

$$dl^2 = c^2 dt^2 - a^2 \begin{cases} d\chi^2 + \sin^2 \chi d\Omega^2 \\ d\chi^2 + \chi^2 d\Omega^2 \\ d\chi^2 + \sinh^2 \chi d\Omega^2 \end{cases} \quad (1.5)$$

where we introduce a variable χ defined so that if $k = +1$ then $r = \sin \chi$, if $k = 0$ then $r = \chi$, and if $k = -1$ then $r = \sinh \chi$.

The properties of the sphere and of the hyperboloid actually carry over to the 3D case: a spacetime with positive curvature is bounded, while if it is flat or hyperbolic it is unbounded.

1.2.2 Other forms of the RW metric

If we wish to use cartesian coordinates the RW metric takes the following expression:

$$ds^2 = c^2 dt^2 - a^2(t) \left(1 + \frac{k|x|^2}{4} \right)^{-2} (dx^2 + dy^2 + dz^2). \quad (1.6)$$

Universes in which a is a constant are called *Einstein spaces*.

We can also change time variable: the *conformal time* η is such that $dt = a(\eta) d\eta$, where $a(\eta) \stackrel{\text{def}}{=} a(t(\eta))$: so, we will have

$$ds^2 = a^2(\eta) \left(c^2 d\eta^2 - \left(\frac{dr^2}{1 - kr^2} + r^2 d\Omega \right) \right). \quad (1.7)$$

Conformal transformations are defined to be those which preserve angles locally; since angles are defined through the metric but the angle between two vectors does not change if they are rescaled, this can be translated into the condition that the metric is rescaled by a generic function:

$$g_{ab} \rightarrow a^2(x^i)g_{ab}. \quad (1.8)$$

In our case, the dependence on the point in spacetime is reduced to a dependence on the cosmic time only since we have symmetry with respect to spatial translations.

In general our spacetime does not have conformal symmetry: a generic massive particle in it has a Compton wavelength $\lambda = h/(mc)$ which defines its interaction cross section, so if the spacetime expands an ensemble of these particles will have different dynamics.

However, conformal geometry is useful for the description of particles which have no characteristic length, such as photons. The photons of the CMB look like they are thermal: they were thermal in the early universe, since they were in thermal equilibrium with matter (photons and matter particles were constantly Compton-scattering off each other); when matter and radiation decoupled the photons scattered for the last time and then kept travelling. The universe has since expanded by a factor ≈ 1090 , but because of the fact that the photons do not have a characteristic length their distribution can still be modelled as a blackbody distribution for an appropriately rescaled temperature.

1.3 The energy budget of the universe

Up until now we did not consider any dynamics in our spacetime. We will discuss this topic in more detail in later sections, but for now we give the result: the dynamics of the universe are described by the Friedmann equations:

$$\dot{a}^2 = \frac{8\pi G_N}{3}\rho a^2 - kc^2 \quad (1.9a)$$

$$\ddot{a} = -\frac{4\pi G_N}{3}a\left(\rho + \frac{3P}{c^2}\right) \quad (1.9b)$$

$$\dot{\rho} = -3\frac{\dot{a}}{a}\left(\rho + \frac{P}{c^2}\right) \quad (1.9c)$$

where dots denote differentiation with respect to the *cosmic time* t , G_N is Newton's gravitational constant, $\rho = \rho(t)$ is the energy density, and $P = P(t)$ is the isotropic pressure.

The curvature k appears in the first equation: so we can try to measure it by comparing the other two terms in the equations. This way, we can determine whether the universe is flat or curved.

In order to discuss this problem, let us establish some notation: an important parameter is $H(t) \stackrel{\text{def}}{=} \dot{a}/a$, the *Hubble parameter*. We can write an equation for it from the first Friedmann one:

$$H^2 = \frac{8\pi G}{3}\rho - \frac{kc^2}{a^2} \quad (1.10)$$

If $k = 0$, then there we must have a critical energy density

$$\rho_C(t) = 3H^2(t)/(8\pi G), \quad (1.11)$$

and we define $\Omega(t) = \rho(t)/\rho_C(t)$. For a flat universe, $\Omega = 1$, and so we can determine $k = \text{sign}(\Omega - 1)$, since:

$$1 = \frac{8\pi G}{3} \frac{\rho}{H^2} - k \frac{c^2}{a^2 H^2} \quad (1.12) \quad \begin{array}{l} \text{Dividing equation} \\ (1.10) \text{ through by } H^2 \end{array}$$

$$\frac{\rho}{\rho_C} - 1 = k \frac{c^2}{a^2 H^2} \quad (1.13)$$

$$\text{sign}(\Omega - 1) = \text{sign}\left(k \frac{c^2}{a^2 H^2}\right) = \text{sign}(k) = k. \quad (1.14) \quad \begin{array}{l} \text{Took sign of both} \\ \text{sides} \end{array}$$

The alternative is measuring k directly, say, by drawing really large triangles and measuring the sum of their internal angles. The latter is a “Newtonian” measurement, while the former is a “GR” measurement. The Newtonian measurement is much harder to perform.

Further, we define the **Hubble parameter**:

$$H_0 = H(t_0) = 100h \times \text{km s}^{-1} \text{Mpc}^{-1} \approx 70 \text{ km s}^{-1} \text{Mpc}^{-1} \quad (1.15)$$

where $h \approx 0.7$ is a number, and t_0 just means *now*.⁴

The reason for this peculiar way to write the constant is that historically it has been difficult to determine the value of H_0 precisely, and it affects many astronomical conversions: keeping it indeterminate in this way allows us to quickly update our old estimates if we measure H_0 more precisely later. Historically, in the American school, the pupils of Hubble thought $h \sim 0.5$, while the French school thought $h \sim 1$. Now, a great issue in cosmology is the disagreement between the measurements of H_0 obtained from the cosmic distance ladder and those obtained from the CMB [Won+19].

1.3.1 Energy density

So, in order to determine the spatial curvature of the universe we need to look at $\Omega \propto \rho$: so, we need to measure the energy density of the universe.

How do we do it?

Let us start by considering the energy density *today* (index 0) due to *galaxies* (index g): ρ_{0g} . We do not directly observe the mass⁵ of galaxies: we can only measure their luminosities.

⁴Also, pc means “parsec”:

$$1 \text{ pc} = \frac{1 \text{ AU}}{1 \text{ arcsec}} \approx 3.085 \times 10^{16} \text{ m} \approx 3.26 \text{ ly}, \quad (1.16)$$

where the angle is to be interpreted as adimensional (in radians), and AU is an astronomical unit, the Earth-Sun average distance. The definition is as such because of a way we have to measure the distances to nearby objects, by measuring their parallax between winter and summer: if they are close enough, they will have apparently moved with respect to further objects, because of the movement of the Earth around the Sun.

⁵Here the terms mass and energy are used equivalently: the velocity of galaxies with respect to the CMB is nonrelativistic, so we approximate $E = \gamma mc^2 \approx mc^2$.

So, we do the following: we compute the mean value of ρ , the mass per unit volume, with the aid of the galaxy luminosity per unit volume ℓ : we find

$$\langle \rho \rangle = \left\langle \ell \frac{\rho}{\ell} \right\rangle = \langle \ell \rangle \left\langle \frac{\rho}{\ell} \right\rangle = \langle \ell \rangle \left\langle \frac{M}{L} \right\rangle, \quad (1.17)$$

where $\langle M/L \rangle$ is the average ratio of mass over luminosity per galaxy: we had a ratio of densities, but since we are considering averages we can integrate above and below with respect to the spatial volume.

In order to simplify the expression, we assume that the ρ/ℓ ratio is uncorrelated to the luminosity per unit volume ℓ : is there a justification for this, beyond saying that it is true by homogeneity?

It is measured in units of solar mass over solar luminosity: M_{\odot}/L_{\odot} . Reference values for these are $M_{\odot} \sim 1.99 \times 10^{33} \text{ g}$, while $L_{\odot} \sim 3.9 \times 10^{33} \text{ erg s}^{-1}$.

We denote $\langle \ell \rangle \stackrel{\text{def}}{=} \mathcal{L}_g$: it is the mean (intrinsic, bolometric⁶) luminosity of galaxies per unit volume.

By definition, it is given by

$$\mathcal{L}_g = \int_0^{\infty} L \Phi(L) dL, \quad (1.18)$$

where $\Phi(L)$ is the number density of galaxies per unit volume and unit luminosity: the *luminosity function*.

The Schechter function is an empirical estimate for the shape of this distribution:

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$$\Phi(L) = \frac{\Phi^*}{L^*} \left(\frac{L}{L^*} \right)^{-\alpha} \exp\left(-\frac{L}{L^*}\right), \quad (1.19)$$

where Φ_* , L_* and α are parameters, with dimensions of respectively a number density, a luminosity and a pure number.

These can be fit by observation: we find $\Phi^* \approx 10^{-2} h^3 \text{ Mpc}^{-3}$, $L^* \approx 10^{10} h^{-2} L_{\odot}$ and $\alpha \approx 1$.

The integral for \mathcal{L}_g converges despite the divergence of $\Phi(L)$ as $L \rightarrow 0$, since it is multiplied by L : so we do not need to really worry about the low-luminosity divergence of the distribution.

The result of the integral for a generic value of α is $\mathcal{L}_g = \Phi^* L^* \Gamma(2 - \alpha)$, where Γ is the Euler gamma function; for the $\alpha = 1$ case we get a factor $\Gamma(2 - 1) = 1$.

Numerically, we get $\mathcal{L}_g \approx (2.0 \pm 0.7) \times 10^8 h L_{\odot} \text{ Mpc}^{-3}$.

This is reported incorrectly in [Pac18] as $\sim 10^{18}$! The correct figure can be found in [LC02, eq. 4.5.14]. I do not think there is much of a point in reporting the specific value, since the parameters are already given as order-of-magnitude estimates.

Now, we must estimate $\langle M/L \rangle$. The luminosity of galaxies can be measured readily, the great difficulty lies in estimating their mass.

⁶Bolometric means “total, over all wavelengths”, as opposed to the luminosity in a certain wavelength band, which is easier to measure in astronomy.

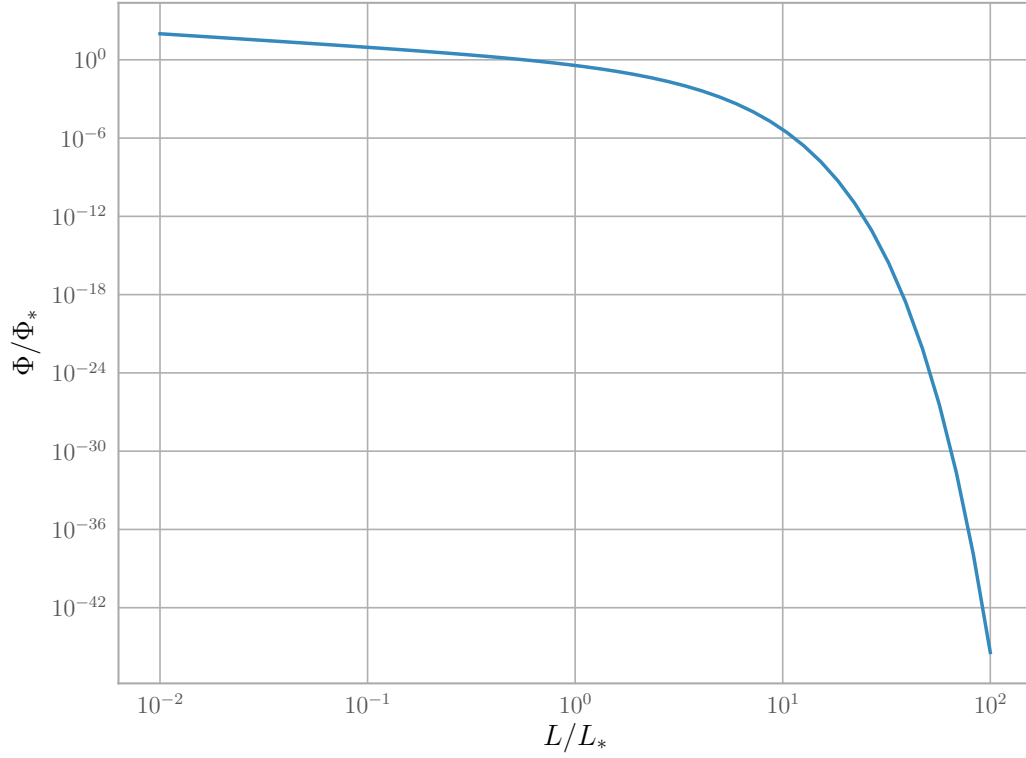


Figure 1.2: A rough plot of the Schechter function for $\alpha = 1$.

1.3.2 Estimating the masses of galaxies

We must distinguish between the different shapes of the galaxies: spiral galaxies are characterized by rotation of the stars about the galactic center, while in elliptical galaxies the stars' motion is disordered.

Spiral galaxies

If we see a spiral galaxy edge-on, we will have a side of it coming in our direction, and the other side moving away from us (after correcting for other sources of Doppler shift, such as the velocity of the whole galaxy). So, using the Doppler effect we can measure the distribution of the velocity in the galaxy as a function of the radius.

In order to get a theoretical model, we can approximate the galaxy as a sphere: this is very rough (spiral galaxies are closer to being disk-like), but it gives the same qualitative result, so there is no need for a more precise model in this context.

We model the galaxy velocity distribution using Newtonian mechanics: the GR corrections are negligible at these scales, galaxies are much larger than their Schwarzschild radii.

Equating the gravitational force GM/R^2 to the centripetal acceleration v^2/R we find:

$$v = \sqrt{\frac{GM}{R}}, \quad (1.20)$$

where v and M are functions of R : M is the mass contained in the spherical shell of radius R , and v is the orbital velocity at the boundary of the shell.

In the inside regions of the galaxy, where $M(R) \propto R^3$ since the density is approximately constant we will have $v \propto R$, while in the outskirts of the galaxy $M(R)$ will not change much, since all the mass is inside, so we will have $v \propto R^{-1/2}$.

Our prediction is then a roughly linear region, and then a region with $v \sim R^{-1/2}$. This is shown as the “Predicted” curve in figure 1.3.

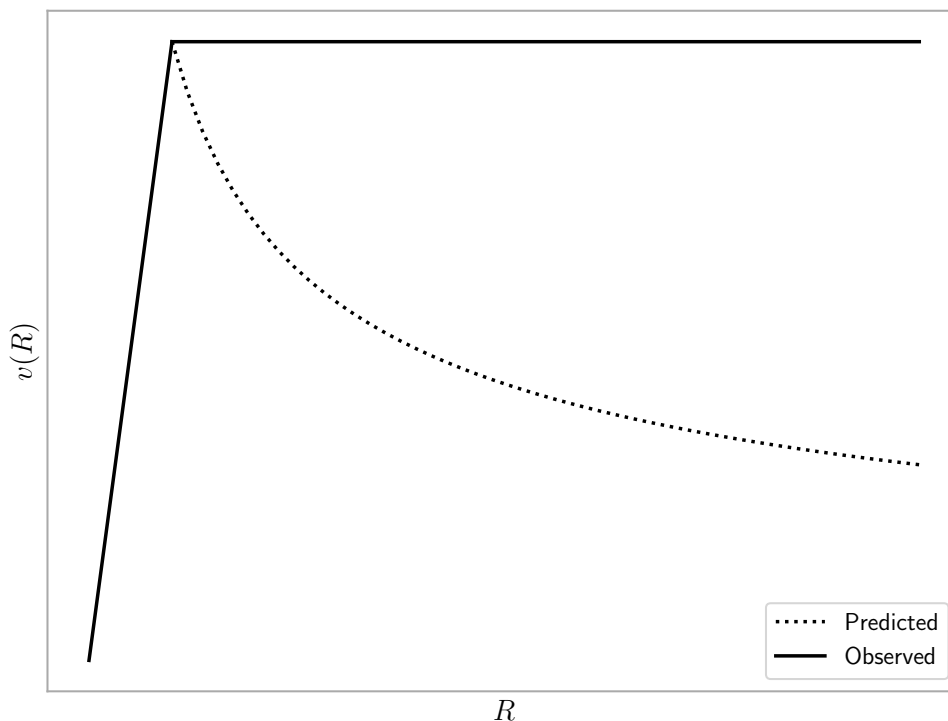


Figure 1.3: Predicted and observed velocity distribution for galaxies. The point at which they start to diverge is approximately the radius of the bulk of the galaxy. This plot is approximate, realistic models do not have such sharp corners.

Instead of this, when in the 1980s people started to be able to measure this curve accurately they saw that, after the linear region, $v(R)$ was approximately constant. So, is Newtonian gravity wrong?

An option to solve this problem was proposed by Milgrom and collaborators: it is called MOfified Newtonian Dynamics, or MOND: they propose that gravity is not actually always described by a r^{-1} potential, but instead at low accelerations it behaves differently. Specifically, they posit that the gravitational acceleration g should be modulated by a factor $\mu(g/a_0)$, where $a_0 \approx 8 \times 10^{-10} \text{ m/s}^2$ is a characteristic acceleration while $\mu(x)$ is an adimensional function such that $\mu \rightarrow 1$ for $x \gg 1$ and $\mu \rightarrow x$ for $x \ll 1$, such as $\mu = x/(1+x)$ [BM84].

This approach is Newtonian but there are also relativistic MOND variants. They do not match observation as well as the alternative approach does.⁷ The heaviest thing weighing against MOND is the fact that even using it we still need dark matter in order to fully explain observations.

Dark matter

The alternative option is that Newtonian mechanics describes galactic mechanics well, but the galaxy's matter distribution inferred from our observations is actually smaller than the real distribution, which extends outward further than the matter we can see: this is **dark matter**.

We'd need mass obeying $M(R) \propto R$ in order to have a constant value of v : since $M(R) = 4\pi \int_0^{R_{\max}} R^2 \rho(R) dR$, we need the density profile to decay like $\rho(R) \propto R^{-2}$. This is called a *thermal* density profile: we call it the *dark matter halo*, which surrounds all spiral galaxies.

We do not know what dark matter is: we can say that it interacts gravitationally but not electromagnetically. People tend to believe that it is made up of beyond-the-standard-model particles, like a *neutralino* or an *axion*. Historically, people thought the effect could be due to massive neutrinos; however their mass would need to be around 30 eV, and the analysis of the CMB data showed that the sum of the masses of the neutrino species must be $\sum m_\nu < 0.120$ [Col18, Table 7].

The total density of matter (dark+regular) is ~ 6 times more than that of regular matter alone: this must be accounted for in our estimate of the M/L ratio for spiral galaxies (since we have additional mass but not additional luminosity): with this correction, we find that for spiral galaxies

$$\left\langle \frac{M}{L} \right\rangle \approx 300h \frac{M_\odot}{L_\odot}, \quad (1.21)$$

Elliptical galaxies

If galaxies are not spiral-shaped, we have to weigh them in a different way: the Doppler broadening of spectral lines gives us a measure of the root-mean-square velocity.

⁷MOND would be compatible with the speed of gravity being less than the speed of light, which is equivalent to the graviton being massive. Recent measurements of gravitational waves seem to agree with the general-relativistic prediction that it is massless.

Later in the course we will obtain the (nonrelativistic) *virial theorem*, now we just state it: if T is the kinetic energy of a gravitationally bound system, U is the potential energy, then

$$2T + U = 0 \quad (1.22)$$

holds when the inertia tensor stabilizes, that is, when we have dynamical equilibrium.

The kinetic energy is $T = \frac{1}{2} \sum_i m_i v_i^2 = \frac{1}{2} M \langle v_r^2 \rangle$, where v_r is the radial⁸ component of the velocity, which we expect to account for one third of total energy by the equipartition theorem. M is the total mass of the galaxy.

The potential energy, instead, is $U = -GM^2/R$. Substituting these expressions into the virial theorem we get

$$2 \times \frac{3}{2} M \langle v_r^2 \rangle - \frac{GM^2}{R} = 0 \quad (1.23)$$

$$M = \frac{3R}{G} \langle v_r^2 \rangle, \quad (1.24)$$

so if we can measure $\langle v_r^2 \rangle$ through Doppler broadening and we can give a reasonable estimate for the radius R of the galaxy we can give an estimate for M .

Global matter contributions

Accounting for the dark matter mass, we get $\langle M/L \rangle \approx 300hM_\odot/L_\odot$.

To date, the best estimate for ρ_{0c} is

$$\rho_{0c} \approx 1.88 \times 10^{-29} h^2 \text{g/cm}^3, \quad (1.25)$$

so, in order to have $\Omega = 1$, we'd need $\langle M/L \rangle$ to be equal to:

$$\left\langle \frac{M}{L} \right\rangle = \frac{\rho_{0c}}{\mathcal{L}_g} \approx \frac{1.88 \times 10^{-29} h^2 \text{g/cm}^3}{2 \times 10^8 h L_\odot / \text{Mpc}^3} \approx 1390 h M_\odot / L_\odot. \quad (1.26)$$

We can define quantities of the form $\Omega_{0i} = \rho_{0i}/\rho_{0c}$, where i is a type of matter, such as baryonic matter, dark matter, dark energy, radiation and so on, whose density is represented as ρ_{0i} .

These variables quantify how much, at the present, time, of the cosmic energy budget is accounted for by that type of matter.

So, only $\Omega_{0b} \approx 5\%$ of the energy budget is given by baryonic matter (not all of which is visible), while around $\Omega_{0DM} \approx 27\%$ is dark matter. Together, they are just denoted as “matter”, and $\Omega_{0m} \approx 30\%$.

We can ask ourselves: is dark matter actually baryonic matter which for some reason we cannot see, such as black holes or brown dwarfs?

This cannot be the case: our observations, combined with models for primordial nucleosynthesis, gives the following bounds for the baryonic energy density:

$$0.013 \leq \Omega_{0b} h^2 \leq 0.025. \quad (1.27)$$

⁸By “radial” we mean directed towards us, not towards the center of the galaxy.

The upper bound for Ω_{0b} is around $2.5\%/h^2 \approx 5.4\%$.

This would seem to indicate that $\Omega_0 \approx 0.3 \ll 1$: however we are failing to consider a crucial contribution. Consider the second Friedmann equation (1.9b): in the Newtonian limit $P \sim 0$ while $\rho > 0$, so we get $\ddot{a} < 0$: the universe contracts. This is not what is observed: we actually see it in accelerated expansion.

Dark energy

The measurements leading to this conclusion are performed by estimating the distance and redshift of far-away objects whose intrinsic luminosity is well known, called *standard candles*: the most commonly used are type Ia supernovae and Cepheid variables.

So, if the expansion is accelerated then $\ddot{a} > 0$: this means, again from the second Friedmann equation (1.9b), that $P < -\rho c^2/3$. This is commonly expressed by defining $w = P/\rho c^2$. In order to have accelerated expansion we need $w < -1/3$; what is observed is closer to $w \sim -1$.

This negative pressure has the effect of a *tension*, pulling the universe apart. This contribution is also called a “cosmological constant term”, for reasons which will be expanded upon in later sections, and denoted with the letter Λ .

We cannot see directly neither dark matter nor dark energy: how do we distinguish the two? Dark matter tends to cluster, while dark energy is uniformly distributed.

From observations of both the anisotropies in the CMB and the distribution of galaxies we can determine that $\Omega_\Lambda \approx 0.7$.

Radiation

We still need to compute the contribution of the energy of electromagnetic and neutrino radiation to the total energy balance. Let us start with radiation: the greatest fraction of the radiation energy density is contained in the CMB, which is extremely close to a Planckian distribution:

$$B(\nu, T) = \frac{2h}{c^2} \frac{\nu^3}{\exp\left(\frac{h\nu}{k_B T}\right) - 1}, \quad (1.28)$$

with $T = T_{0\gamma} \approx (2.725 \pm 0.001)$ K. B is a measure of spectral intensity: it is measured in units of energy per unit second, area, solid angle and frequency. This is a well-known distribution, whose integral is given by

$$\rho_{0\gamma} = \frac{\sigma_r T_{0\gamma}^4}{c^2} = 4.8 \times 10^{-34} \text{ g cm}^{-3} \quad (1.29)$$

where $\sigma_r = \pi^2 k_B^4 / (15 \hbar^3 c^3)$, while $\sigma_{SB} = \sigma_r c / 4$.

Doing the computation, by

```
1 (4 * cosmo.Tcmb0**4 * ac.sigma_sb / ac.c**3 ).cgs
```

gives $4.6 \times 10^{-34} \text{ g/cm}^3$: it is very close to the figure of 4.8 given in [Pac18], but I do not see where the discrepancy would come from. (Confermo il tuo conto, magari ha cacciato un typo)

So, we have that the radiation contribution to the global energy balance is $\Omega_{0\gamma} \approx 2.5 \times 10^{-5} h^{-2} < 0.01 \%$, definitely negligible.

Typo in [Pac18]: exponent of h is -2 , not $+2$.

We are going to show in later sections that if neutrinos were massless, their temperature would be $T_\nu = (4/11)^{1/3} T_\gamma < T_\gamma$.

They might not be massless, and if they are not the main contribution to the energy density they will give will be from their masses. However, recent observations (for example, by the Planck satellite) are bounding the mass of the neutrinos, $\sum m_\nu \leq 0.12 \text{ eV}$: we have

$$\rho_\nu = 3N_\nu \frac{\langle m_\nu \rangle}{10 \text{ eV}} 10^{-30} \text{ g cm}^{-3}, \quad (1.30)$$

where N_ν is the number of neutrino species. Even if we assume the upper bound, $N_\nu \langle m_\nu \rangle = 0.12 \text{ eV}$, we get $\Omega_{0\nu} < 0.5 \%$.

Conclusions

If we add up all the contributions to $\Omega = \Omega_b + \Omega_{DM} + \Omega_\Lambda$ (neglecting, as we said, the contributions by EM radiation and neutrinos) we find experimentally $\Omega_k = \Omega - 1 \approx (5_{-40}^{+38}) \times 10^{-4}$ [Pla+19, Table A.2].

So, with the data we have currently we cannot determine the sign of the universe's spatial curvature.

1.3.3 The Hubble law

At the end of the 1920s Edwin Hubble compared the estimates for the distances of far-away galaxies (obtained through standard candles and other types of estimates) to their velocities relative to us, as measured through their redshift. His results are shown in figure 1.4: he obtained a roughly linear relation in the form:

$$v = H_0 d, \quad (1.31)$$

where v is the velocity of the galaxies, and d is their distance from us, while H_0 is a constant of proportionality. Hubble's measurements suggested $H_0 \sim 500 \text{ km/s/Mpc}$, more refined modern ones using techniques similar to those used by Hubble yield a value of $(73.24 \pm 1.74) \text{ km/s/Mpc}$ [Rie+16].

Measurements of H_0 through analysis of the CMB gives an incompatible value: $H_0 = (67.8 \pm 0.9) \text{ km/s/Mpc}$ [Col16].

Let us now show that this H_0 is actually the same one we defined before, $H_0 = \dot{a}/a$. H_0 is called the **Hubble constant**, since it is constant with respect to the direction we look in the sky.

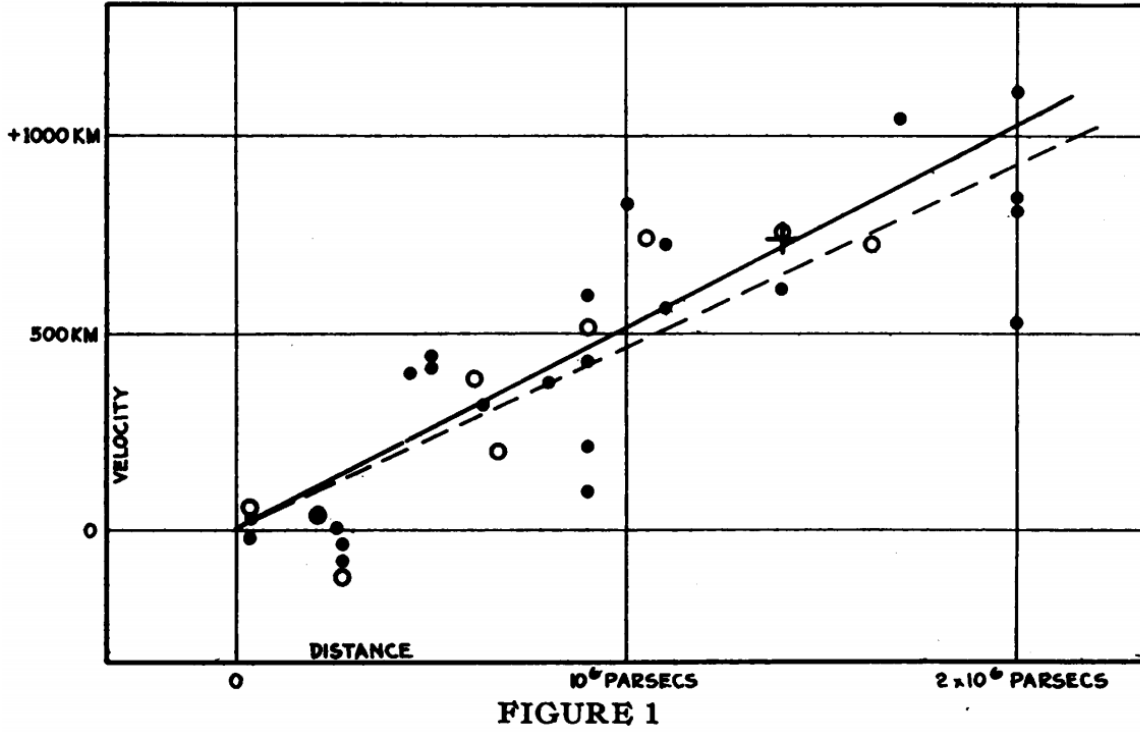


Figure 1.4: Original velocity versus distance data from Edwin Hubble's paper in 1929 [Hub29].

We are considering the distance connecting us (the center of our reference frame) to a distance galaxy: so we drop the angular part in the flat ($k = 0$) Friedmann-Lemaître-Robertson-Walker (FLRW) line element:

$$ds^2 = c^2 dt^2 - a^2(t) dr^2 \quad (1.32)$$

So, at a fixed time the physical distance is given by $d = a(t)r$: therefore $v = \dot{d} = \dot{a}r = \frac{\dot{a}}{a}d = H_0 d$.

Do note that we neglected the temporal part of the metric: this is equivalent to assuming that photons travel instantaneously. Assuming the universe is spatially flat is correct up to second order: for a general value of k we have

$$d = a \int_0^r \frac{d\tilde{r}}{\sqrt{1 - k\tilde{r}^2}} = a \left(r + \mathcal{O}(r^3) \right), \quad (1.33)$$

since the integral gives either r , $\arcsin(r)$ or $\operatorname{arcsinh}(r)$, depending on k and all three of these equal r up to second order.

So, this is Newtonian and rough, but it gives us the correct intuitive idea. We now wish to make this reasoning more precise.

The first step, since we want to discuss our observations of light, is to define the redshift.

Definition 1.3.1 (Redshift). *The redshift z of a photon is defined by*

$$z = \frac{\lambda_0 - \lambda_e}{\lambda_e} = \frac{\lambda_0}{\lambda_e} - 1 = \frac{\nu_e}{\nu_0} - 1, \quad (1.34)$$

where λ_0 and λ_e are the observed and emission wavelengths respectively, while ν are frequencies with the same notation.

We will show that the redshift can be found from the ratio of the scale factors now and at emission: $1 + z = a_0/a_e$. Therefore, $\nu_0/\nu_e = a_e/a_0$.

We wish to study the distribution of light from an astronomical source: in Minkowski spacetime the apparent luminosity ℓ decreases like r^{-2} if r is the distance from the object. In a generic spacetime this will not be the case: however, we can define a measure of spatial distance d_L such that $\ell = L/(4\pi d_L^2)$, where L is the intrinsic luminosity of the object.

Do note that ℓ is dimensionally a luminosity flux: it is measured in units of energy per unit time per unit area.

Definition 1.3.2. *The luminosity distance d_L is defined as:*

$$d_L = \sqrt{\frac{L}{4\pi\ell}}. \quad (1.35)$$

Since $L \propto \ell$, this is a well-defined measure of distance between two generic points in spacetime, regardless of the presence of a source of light there.

How do we relate the luminosity distance and the scale factor? The radiation from our source is spread on a sphere: we integrate the angular part of the FLRW metric over a sphere of fixed comoving radius r to find its area.

The metric restricted to the angular coordinates at fixed r is given by

$$ds^2 = a^2 r^2 (d\theta^2 + \sin^2 \theta d\varphi^2), \quad (1.36)$$

so the area form on the surface of the sphere is

$$dA = \sqrt{\det g} d\theta \wedge d\varphi, \quad (1.37)$$

where $\sqrt{\det g} = \sqrt{a^4 r^4 \sin^2 \theta} = a^2 r^2 \sin \theta$. So,

$$A = \int_{S^2} dA = a^2 r^2 \int_0^\pi d\theta \int_0^{2\pi} d\varphi \sin \theta = 4\pi a^2 r^2, \quad (1.38)$$

where we substituted the wedge product for the regular tensor product of the differentials since our axes are orthogonal. Now, at which time do we compute the scale factor? We are measuring the flux at the surface of the sphere, at the time at which we are observing: therefore we need to compute it at observation time. This gives us $A = 4\pi r^2 a_0^2$.

Now, the emitted luminosity is in the form:

$$L = \frac{dN}{dt_e} \langle h\nu_e \rangle, \quad (1.39)$$

where dN/dt is the number of photons emitted per unit time, whose average energy is $\langle h\nu \rangle$. From the point of view of the observer, the number of photons is the same, while the frequency of the observed photon and the time interval dt_e change: specifically, $\nu_e = \nu_0 a_0/a_e$ and $dt_e = dt_0 a_e/a_0$.⁹ Therefore, the observed absolute luminosity obeys the relation $L = L_0(a_0/a_e)^2$.

So, putting everything together we get:

$$\ell = \frac{L_0}{A} = \frac{L}{4\pi r^2 a_0^2} \left(\frac{a_e}{a_0} \right)^2, \quad (1.40)$$

note that the value of k does not enter into the equation.

Therefore:

$$d_L = \sqrt{\frac{L}{4\pi\ell}} = \sqrt{\frac{4\pi r^2 a_0^2}{4\pi} \left(\frac{a_0}{a_e} \right)^2} = \frac{a_0^2}{a_e} r = a_0(1+z)r. \quad (1.41)$$

Another way of defining a distance is the one we get by directly integrating the radial part of the line element: this way, we are effectively using a space-like measuring stick, working at fixed cosmic time and fixing both of the angles θ and φ .

Thursday
2019-10-10

Definition 1.3.3 (Proper distance). *The proper distance d_P at a fixed cosmic time t to an object at a comoving radial coordinate r is given by:*

$$d_P = a(t) \int_0^r \frac{d\tilde{r}}{\sqrt{1 - k\tilde{r}^2}}. \quad (1.42)$$

A derivation of the Hubble law

We want to derive the Hubble law ($v = H_0 d$) mathematically. It can also be restated as $cz = H_0 d$: the observed velocity of recession of the objects is measured through redshift, which is described (in the nonrelativistic limit¹⁰) by the formula

$$\lambda_0 = \lambda_e \left(1 \pm \frac{v}{c} \right), \quad (1.44)$$

where the sign in the \pm is a plus if the object is receding from us. We also defined z through $\lambda_0/\lambda_e = 1 + z$: this means that we can identify $z = v/c$.

Now, we are going to “move away from the current epoch by Taylor expanding”: the scale factor at a time t , $a(t)$, can be written as

$$a(t) = a_0 + \dot{a}_0(t - t_0) + \frac{1}{2}\ddot{a}_0(t - t_0)^2 + \mathcal{O}(|t - t_0|^3) \quad (1.45a)$$

We drop the error
term

⁹This is the case even though the temporal component of the metric does not change, since we are not considering a fixed instance of cosmic time (which is unphysical): instead, we are considering the time intervals that the emitter and observed measure between the crests of a light wave sent from one to the other.

¹⁰The relativistic expression is

$$1 + z = \sqrt{\frac{1 + v/c}{1 - v/c}} \gtrsim 1 + \frac{v}{c}. \quad (1.43)$$

$$\approx a_0 \left(1 + H_0(t - t_0) - \frac{1}{2} q_0 H_0^2 (t - t_0)^2 \right), \quad (1.45b) \quad \begin{array}{l} \text{Substituted} \\ H_0 = \dot{a}_0/a_0. \end{array}$$

where $q_0 \stackrel{\text{def}}{=} -\ddot{a}_0 a_0 / (\dot{a}_0)^2$ is called the *deceleration parameter* by historical reasons: people thought they would see deceleration $\ddot{a} < 0$ when first writing this, so a positive q_0 , but the deceleration parameter is instead measured to be negative.

Now, $1 + z = a_0/a$ can be expressed as:

$$1 + z \simeq \left(1 + H_0(t - t_0) - \frac{1}{2} q_0 H_0^2 (t - t_0)^2 \right)^{-1}. \quad (1.46)$$

Do note that this is derived starting with a formula which is correct up to second order in the time interval $\Delta t = t - t_0$: when expanding we cannot trust terms of order higher than second. Expanding with this in mind we get:¹¹

$$1 + z \simeq 1 - H_0 \Delta t + \frac{q_0}{2} H_0^2 \Delta t^2 + H_0^2 \Delta t^2, \quad (1.51)$$

therefore

$$z = H_0(t_0 - t) + \left(1 + \frac{q_0}{2} \right) H_0^2 (t_0 - t)^2 \quad (1.52)$$

$$= (t_0 - t) \left[H_0 + \left(1 + \frac{q_0}{2} \right) H_0^2 (t_0 - t) \right]. \quad (1.53)$$

Changed the time intervals to $t_0 - t = -\Delta t$, the square of which is the same as before.

Bringing the bracket to the other side we get

$$t_0 - t = z \left[H_0 + \left(1 + \frac{q_0}{2} \right) H_0^2 (t_0 - t) \right]^{-1} \quad (1.54)$$

$$= z \left[H_0 + \left(1 + \frac{q_0}{2} \right) H_0 z \right]^{-1} \quad (1.55)$$

$$= \frac{z}{H_0} \left[1 + \left(1 + \frac{q_0}{2} \right) z \right]^{-1}, \quad (1.56)$$

¹¹We need to compute the first and second derivatives of

$$\left(1 + H_0 \Delta t - \frac{q_0}{2} H_0^2 \Delta t^2 \right)^{-1} = (1 + a \Delta t + b \Delta t^2)^{-1} = f(\Delta t) : \quad (1.47)$$

they are

$$\frac{df}{d\Delta t} = -(1 + a \Delta t + b \Delta t^2)^{-2} (a + 2b \Delta t) \quad (1.48)$$

$$\frac{d^2 f}{d\Delta t^2} = 2(a + 2b \Delta t)^2 (1 + a \Delta t + b \Delta t^2) - (1 + a \Delta t + b \Delta t^2)^{-2} (2b), \quad (1.49)$$

so we have

$$\left. \frac{df}{d\Delta t} \right|_{\Delta t=0} = -a \quad \text{and} \quad \left. \frac{1}{2!} \frac{d^2 f}{d\Delta t^2} \right|_{\Delta t=0} = \frac{1}{2!} (2a^2 - 2b) = a^2 - b. \quad (1.50)$$

where we substituted the first order expression $t_0 - t = z/H_0$: we are allowed to make this substitution since the expression is multiplied by z , which has the same asymptotic order as $t_0 - t$ (since H_0 is finite): working to first order inside the brackets is equivalent to working to second order in the global expression.

By the same reasoning, we can expand the inverse bracket to first order:

$$t_0 - t = \frac{z}{H_0} \left[1 - \left(1 + \frac{q_0}{2} z \right) \right] = \frac{z}{H_0} - \left(1 + \frac{q_0}{2} \right) \frac{z^2}{H_0}, \quad (1.57)$$

We would like the time interval to disappear: we want a distance, not a time, so we should seek an expression for r instead of Δt . We are observing photons, for which $ds^2 = 0$, which is equivalent $c^2 dt^2 = a^2(t) dr^2 / (1 - kr^2)$. Taking a square root and integrating we get:

$$\int_t^{t_0} \frac{c dt}{a(t)} = \pm \int_r^0 \frac{d\tilde{r}}{\sqrt{1 - k\tilde{r}^2}}, \quad (1.58)$$

where we should select the negative sign since we want positive quantities on both sides. The other choice would correspond to the photon being emitted from the Earth and received at the comoving radius of the source.

The integral on the right hand side can be solved analytically: it is

$$- \int_r^0 \frac{d\tilde{r}}{\sqrt{1 - k\tilde{r}^2}} = \begin{cases} \arcsin r = r + \mathcal{O}(r^3) & k = 1 \\ r & k = 0 \\ \operatorname{arcsinh} r = r + \mathcal{O}(r^3) & k = -1 \end{cases} \quad (1.59)$$

in all cases, it is just r up to *second* order (since the next term in the expansion of an arcsine or hyperbolic arcsine is of third order).

On the left hand side, we can substitute in $a(t)$ from equation (1.45b):

$$\int_t^{t_0} \frac{c dt}{a(t)} = \frac{c}{a_0} \int_t^{t_0} d\tilde{t} \left[1 + H_0(\tilde{t} - t_0) + \mathcal{O}(\Delta \tilde{t}^2) \right]^{-1} \quad (1.60)$$

$$= \frac{c}{a_0} \left[t_0 - t + \frac{H_0}{2} (t - t_0)^2 \right] + \mathcal{O}(\Delta t^3) \quad (1.61)$$

where we used the expression for the scale factor to first order only since the integration raised the order of the estimate by one. we have:

$$\frac{c}{a_0} \left((t_0 - t) + \frac{1}{2} H_0 (t_0 - t)^2 + \mathcal{O}(\Delta t^3) \right) = r, \quad (1.62)$$

since the term proportional to q_0 only gives a third order contribution. We can now substitute the expression for the time difference with respect to the redshift (1.57), only computed to second order:

$$r = \frac{c}{a_0} \left[\frac{z}{H_0} \left(1 - \left(1 + \frac{q_0}{2} \right) z \right) + \frac{H_0}{2} \left(\frac{z}{H_0} \left(1 - \left(1 + \frac{q_0}{2} \right) z \right) \right)^2 \right] \quad (1.63)$$

$$= \frac{c}{a_0} \left[\frac{z}{H_0} - \left(1 + \frac{q_0}{2}\right) \frac{z^2}{H_0} + \frac{H_0}{2} \frac{z^2}{H_0^2} \right] \quad (1.64)$$

Ignored the third and higher order terms in the square.

$$= \frac{c}{a_0 H_0} \left[z - \frac{1}{2}(1 + q_0)z \right]. \quad (1.65)$$

Now, we can insert this expression for r into the formula for the luminosity distance (1.41):

$$d_L = a_0^2 \frac{r}{a} = a_0(1+z)r \quad (1.66)$$

$$= a_0(1+z) \frac{c}{a_0 H_0} \left(z - \frac{1}{2}(1 + q_0)z^2 \right). \quad (1.67)$$

As we expect, the term a_0 disappears: it is a bookkeeping parameter, the physical properties of a universe described by a FLRW metric are invariant under a global rescaling of the scale factor.

Our expression also contains cubic terms in z : removing these to get back to second order we find

$$d_L = \frac{cz}{H_0} (1+z) \left(1 - \frac{1}{2}(1 + q_0)z \right) \quad (1.68)$$

$$= \frac{cz}{H_0} \left(1 + \left(1 - \frac{1}{2} - \frac{1}{2}q_0 \right) z \right) \quad (1.69)$$

$$= \frac{cz}{H_0} \left(1 + \frac{1}{2}(1 - q_0)z \right). \quad (1.70)$$

We can turn this into a relation for cz in terms of d_L by substituting in the first order expression $d_L H_0 = cz$ into the second order term:

$$cz = d_L H_0 + \frac{q_0 - 1}{2} cz^2 \quad (1.71)$$

$$cz = H_0 \left(d_L + \frac{1}{2}(q_0 - 1) \frac{H_0}{c} d_L^2 \right), \quad (1.72)$$

and we can notice that the relation is approximately linear and independent of acceleration for low redshift, but we can detect the acceleration at higher redshift. Typically we need to measure galaxies at least 10 Mpc away in order to detect these second order effects. As we mentioned in the beginning, the data show the parameter q_0 to be negative.

This effect is similar to a Doppler effect, but the analogy is not perfect: the redshift is caused by the expansion of space itself, and the apparent velocities of the galaxies at high redshift would be superluminal.

The reasoning in [Pac18] is not completely convincing: extrapolating a superluminal velocity from $cz \sim 10c$ is meaningless, since these formulas were derived in the nonrelativistic approximation. If we use the correct relativistic formula for the Doppler shift this issue does not arise. We might be able to show that the apparent speed of high-redshift galaxies is superluminal but we need another way to do it. (Decisamente hai ragione, aggiungi poi che la relazione vale per z piccolo e 10 non è piccolo solitamente. Ad ogni modo ha ragione a dire che il rate di espansione potrebbe essere maggiore di c senza creare problemi alla relatività)

Definition 1.3.4 (Angular diameter distance). *The angular diameter distance is defined as the ratio of the object's physical transverse size L to its angular size in radians $\Delta\theta$:*

$$d_A = \frac{L}{\Delta\theta} = a(t)r, \quad (1.73)$$

which is peculiar in that it is not monotonic in z [Hog00]: at $z \gtrsim 1$ it starts decreasing.

Redshift-scale factor relation

Let us prove the statement from before, $\lambda_0/\lambda = a_0/a$: photons are emitted with a certain wavelength λ_e , at a comoving radius r from us, and detected at λ_o .

The line element for the photon is $ds^2 = 0$, therefore $c dt / a(t) = \pm dr / \sqrt{1 - kr^2}$.

As before, we can integrate this relation from the emission to the absorption: we call it $f(r)$ (it can be any of the functions shown in equation (1.59)):

$$\int_t^{t_0} = \frac{c d\tilde{t}}{a(\tilde{t})} = f(r) \quad (1.74)$$

If we map $t \rightarrow t + \delta t$ and $t_0 \rightarrow \delta t_0$ in the integration limits, the integral must be constant since it only depends on r — do note that all the expansion of the universe is accounted for by the increasing scale factor, objects are stationary with respect to the comoving radial coordinate. We are computing the integral for two successive wavefront of the light. Then, we equate the two:

$$f(r) = \int_t^{t_0} \frac{c d\tilde{t}}{a(\tilde{t})} = \int_{t+\delta t}^{t_0+\delta t_0} \frac{c d\tilde{t}}{a(\tilde{t})}, \quad (1.75)$$

which we can split into:

$$\left[\int_{t+\delta t}^t + \int_t^{t_0} + \int_{t_0}^{t_0+\delta t_0} - \int_t^{t_0} \right] \frac{c d\tilde{t}}{a(\tilde{t})} = 0, \quad (1.76)$$

where, since all the integrals have the same argument, we collect it at the end for clarity. We simplify the original integral and swap the integration limits to get:

$$\int_t^{t+\delta t} \frac{c d\tilde{t}}{a(\tilde{t})} = \int_{t_0}^{t_0+\delta t_0} \frac{c d\tilde{t}}{a(\tilde{t})}, \quad (1.77)$$

which can we approximate by

$$\frac{c\delta t}{a(t)} = \frac{c\delta t_0}{a(t_0)}, \quad (1.78)$$

since the periods of the photons we are considering are generally much smaller than the cosmic timescales.

Since the frequency of the emitted and observed photons must be proportional to the inverse of the time intervals δt or δt_0 , we have

$$\nu_e a(t_e) = \nu_o a(t_o), \quad (1.79)$$

therefore

$$1 + z = \frac{\lambda_o}{\lambda_e} = \frac{a_0}{a}. \quad (1.80)$$

Chapter 2

Friedmann models

The Friedmann equations describe the dynamical evolution of the universe, as opposed to the static description given by the FLRW metric.

2.1 A Newtonian derivation of the Friedmann equations

The Friedmann equations are derived starting from the Einstein Field Equations for the FLRW metric, however we can derive them using an almost purely Newtonian argument.

We will not be able to recover the full equations, since a Newtonian fluid's pressure is $P \ll \rho c^2$: its contribution to the stress-energy-momentum tensor is negligible. So, through our argument we will recover the equations we would have with $P = 0$.

The only non-Newtonian step in our derivation is the justification of the Newtonian approximation: we wish to make use of the theorem attributed to Birkhoff, but first derived by the Norwegian physicist J.T. Jebsen [JR05].

Proposition 2.1.1 (Jebsen-Birkhoff). *The only solution to the vacuum Einstein Field Equations which is spherically symmetric is given by the Schwarzschild metric [MTW73, sec. 32.2]:¹*

$$ds^2 = \left(1 - \frac{2GM}{c^2 r}\right) c^2 dt^2 - \left(1 - \frac{2GM}{c^2 r}\right)^{-1} - r^2 d\Omega^2. \quad (2.1)$$

With this in mind, let us take a spacetime with uniform density ρ . We consider a sphere, and imagine taking all the mass inside the sphere away.

It is claimed in [Pac18] that it is a corollary of the JB theorem that the metric inside a spherical shell is Minkowski, this was not obvious to me so I looked for a reference. An article by Zhuang and Yi [ZY12] actually claims that this is false, and that the metric inside the shell is indeed a slice of Schwarzschild spacetime but not necessarily with $M = 0$, and their reasoning seems to make sense.

¹This is not relevant for what we will discuss here, but do note that there is no requirement of the generating (internal) mass distribution to be static: this is, in fact, the reason why spherically symmetric collapsing or pulsating stars cannot emit gravitational waves.

This can also be applied to a hole: in this case, it tells us that the metric inside the cavity is the Minkowski metric.

The mass taken away will be $M(\ell) = \frac{4\pi}{3}\rho\ell^3$, where $\vec{l} = a(t)\vec{r}$ is the radius of the sphere. We suppose that the gravitational field is *weak*:

$$\frac{GM(\ell)}{\ell c^2} \ll 1. \quad (2.2)$$

We put a test mass on the surface of the sphere. What is the motion of the mass due to the gravitational field from the center? It will surely be radial, therefore

$$\ddot{\ell} = -\frac{GM(\ell)}{\ell^2} = -\frac{4\pi G}{3}\rho\ell. \quad (2.3)$$

This seems to give us a net force even though by our hypotheses there should be none, this is actually not an issue since the unit vectors in our equations will go away in the end.

Then, we have

$$\ddot{a}r = -\frac{4\pi G}{3}\rho a r \rightarrow \ddot{a} = -\frac{4\pi G}{3}\rho a. \quad (2.4)$$

This is part of a Friedmann equation: the isotropic pressure term cannot be recovered, it's like the speed of light is infinite now.

Now consider

$$\dot{\ell}\ddot{\ell} = -\frac{GM}{\ell^2}\dot{\ell}, \quad (2.5)$$

therefore

$$\frac{1}{2}\dot{\ell}^2 = \frac{4\pi}{3}G\rho\ell^3 + \text{const}, \quad (2.6)$$

and integrating we get

$$\dot{a}^2 r^2 = \frac{8\pi G}{3}\rho a^2 r^2 + \text{const} \quad (2.7)$$

or, removing the r^2 term, which is a constant,

$$\dot{a}^2 = \frac{8\pi G}{3}\rho a^2 + \text{const}. \quad (2.8)$$

We know that this new constant is related to the energy per particle.

A universe with negative k expands forever and so on.

The number k was badly defined to be only ± 1 or 0 : its newtonian version is much better represented by $k_N = kc^2$.

Fri Oct 11 2019

On Oct 31 Marco Peloso will not give his lecture, so Sabino will do both lectures.

Recalling last lecture: we consider a universe with constant density $\rho \dots$

We can also recover the third Friedmann equation

$$\dot{\rho} = -3\frac{\dot{a}}{a}\left(\rho + \frac{P}{c^2}\right), \quad (2.9)$$

without the last term, which yet again will come from the relativistic consideration. We will consider *ideal fluids*. From thermodynamics we have:

$$dE + p dV = 0. \quad (2.10)$$

We can write $E = \rho c^2 a^3$. Then, this becomes $d(\rho c^2 a^3) + p da^3 = 0$; expanding:

$$c^2 d\rho a^3 + c^2 \rho da^3 + p da^3 = 0. \quad (2.11)$$

This amounts to the third Friedmann equation, exactly.

$$\dot{a}^2 = \frac{8\pi G}{3} \rho a^2 - kc^2 \quad (2.12a)$$

$$\ddot{a} = -\frac{4\pi G}{3} a \left(\rho + \frac{3P}{c^2} \right) \quad (2.12b)$$

$$\dot{\rho} = -\frac{3\dot{a}}{a} \left(\rho + \frac{P}{c^2} \right) \quad (2.12c)$$

The equations are in terms of the three parameters a , ρ and p , which are all functions of time.

The third equation comes from the Bianchi identities $\nabla_\mu G^{\mu\nu} = 0$.

Rewriting the first equation:

$$\ddot{a}^2 = \frac{8\pi G}{3} \rho a^2 - kc^2 \quad (2.13)$$

and

$$2\dot{a}\ddot{a} = \frac{8\pi G}{3} \dot{\rho} a^2 + \frac{16\pi G}{3} \rho \dot{a} a \quad (2.14)$$

We then substitute in the expression we have for $\dot{\rho}$ from the third FE. Everything is multiplied by \dot{a} : if it is not zero we have

$$\frac{\ddot{a}}{a} = -4\pi G \rho - \frac{4\pi G p}{c^2} + \frac{8\pi G}{3} \rho = -4\pi G \left(\frac{\rho}{3} + \frac{P}{c^2} \right). \quad (2.15)$$

The equation system is underdetermined. We have to make an assumption: we will assume our fluid is a *barotropic* perfect fluid: $p \stackrel{!}{=} p(\rho)$.

Very often this equation of state will look like $p = w\rho c^2$, for a constant w (not ω !). This is related to the adiabatic constant γ by $\gamma = w + 1$.

We are going to assume homogeneity and isotropicity.

Some possible equations of state are $p \equiv 0$, or $w = 0$: this means $\gamma = 1$. This is a *dust pressureless fluid*.

What is the speed of sound of our fluid? We only have the adiabatic speed of sound $c_s^2 = \partial p / \partial \rho$, where the derivative is to be taken at constant entropy and is just a total derivative in the barotropic case.

Also in the baryonic case $p = 0$ is a good approximation.

If $p = 0$ we can simplify:

$$\dot{\rho} = -3\frac{\dot{a}}{a}\rho \implies \rho \propto a^{-3}. \quad (2.16)$$

More generally, not assuming $w = 0$ we get:

$$\dot{\rho} = -3(1+w)\rho\frac{\dot{a}}{a} \implies \rho \propto a^{-3(1+w)} = a^{-3-3\gamma} \quad (2.17)$$

Another case is a *gas of photons*: in that case $p = \rho c^2/3$, so $w = \frac{1}{3}$, $\gamma = \frac{1}{3}$, $c_s^2 = c^2/3$: the speed of sound is $c/\sqrt{3}$. These photons are thermal: perturbations can propagate (even without interactions with matter. . .).

In this case we get $\rho \propto p \propto a^{-4}$.

Stiff matter is $p = \rho c^2$, $w = 1$, $\gamma = 2$ and $c_s = c$. This is an incompressible fluids: it is so difficult to set this matter in motion that once one does it travels at the speed of light. Now, $\rho \propto p \propto a^{-6}$.

A possible case is $p = -\rho c^2$: $w = -1$ and $\gamma = 0$: we cannot compute a speed of sound. Now ρ and p are constants. This is the case of dark energy (?).

This can be interpreted as an interpretation of the cosmological constant Λ .

Now we relace the last FE with $w = \text{const}$, $\rho(t) = \rho_*(a(t)/a_*)^{-3(1+w)}$.

Now, if we substitute into the second FE we get that gravity is attractive ($\ddot{a} < 0$) iff $w > -1/3$.

[Plot: ρ vs a : the cosmological constant is constant, matter is decreasing, radiation is decreasing faster].

In this plot, a can be interpreted as the time. We can insert the spatial curvature in the plot: it decreases, but slower than matter. Now, the dark energy in the universe is more important than the curvature.

Let us solve the first FE: inserting the third one we get

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho_*\left(\frac{a}{a_*}\right)^{-3(1+w)} - \frac{kc^2}{a^2}. \quad (2.18)$$

We defined the parameter $\Omega = \frac{8\pi G\rho}{3H^2} = \rho/\rho_C$. Experimentally this is very close to 1. The Einstein-de Sitter model is one where we take $\Omega \equiv 1$: negligible spatial curvature. This amounts to making $k = 0$.

$$\dot{a}^2 = \frac{8\pi G}{3}\rho_*a_*^{3(1+w)}a^{-(1+3w)} \quad (2.19)$$

therefore $\dot{a} = \pm Aa^{\frac{1+3w}{2}}$, or $a^{\frac{1+3w}{2}} da = A dt$. A solution is:

$$a(t) = a_* \left(1 + \frac{3}{2}(1+w)H_*(t-t_*)\right)^{\frac{2}{3(1+w)}} \quad (2.20)$$

where $H_*^2 = \frac{8\pi G}{3}\rho_*$, coupled to:

$$\rho(t) = \rho_* \left(1 + \frac{3}{2}(1+w)H_*(t-t_*)\right)^{-2} \quad (2.21)$$

There is a time where the bracket in $a(t)$ is zero: we call it as t_{BB} define it by

$$1 + \frac{3}{2}(1+w)H_*(t_{\text{BB}} - t_*) = 0. \quad (2.22)$$

Since the curvature scalar is $R \propto H^2$, at t_{BB} the curvature is diverges.

Hakwing & Ellis proved that if $w > -1/3$ we unavoidably must have a Big Bang.

We can define a new time variable by $t_{\text{new}} \equiv (t - t_*) + 2H_*^{-1}/(3(1+w))$. Then, we can just write:

$$a \propto t_{\text{new}}^{\frac{2}{3(1+w)}} \quad (2.23)$$

and this allows us to get rid of t_* .

Inserting this new time variable, we get

$$\rho(t) = \frac{1}{6(1+w)^2\pi 4t^2} \quad (2.24)$$

and the Hubble parameter is:

$$H(t) = \frac{2}{3(1+w)t} \quad (2.25)$$

Some cases are:

$$\begin{cases} w = 0 \implies a \propto t^{2/3} \\ w = 1/3 \implies a \propto t^{1/2} \\ w = 1 \implies a \propto t^{1/3} \end{cases} \quad (2.26)$$

The *De Sitter* universe is one where $w \rightarrow 1$: $a(t) \propto \exp(Ht)$ and $H = \text{const.}$ (CHECK stuff)

Thu Oct 17 2019

If we neglect spatial curvature, which is small, we can write the luminosity distance as an integral which we can compute:

$$d_L \equiv \left(\frac{L}{4\pi\ell} \right)^{1/2}. \quad (2.27)$$

Our metric is the FLRW line element. Then, we can write d_L as:

$$d_L = a_0(1+z)r(z). \quad (2.28)$$

We now define the *conformal time* τ : we want to impose $a^2(\tau) d\tau^2 = dt^2$. Then, the RW line element becomes:

$$ds^2 = a^2(\tau) \left(c^2 d\tau^2 - \frac{dr^2}{1 - kr^2} - r^2 d\Omega^2 \right). \quad (2.29)$$

This is very important when we talk about zero-mass particles, with no intrinsic length scale. Using the variable χ , we have:

$$ds^2 = a^2(\tau) \left(c^2 d\tau^2 - d\chi^2 - f_k^2(\chi) d\Omega^2 \right), \quad (2.30)$$

where $f_k(\chi) = r$ is equal to $\sin(\chi)$, χ or $\sinh(\chi)$ if k is equal to 1, 0 or -1 .

If we look at photons moving radially, we get

$$ds^2 = 0 = a^2(\tau) \left(c^2 d\tau^2 - d\chi^2 \right), \quad (2.31)$$

therefore $c^2 d\tau^2 = d\chi^2$: setting $c = 1$, we get $\tau(t_0) - \tau(t_e) = \chi(r_e) - \chi(0)$, where a subscript e means “emission”.

$$d\tau = \frac{dt}{a} = \frac{da}{a\dot{a}}, \quad (2.32)$$

and now recall $a = a_0/(1+z)$: differentiating this we get

$$da = -\frac{a_0}{(1+z)^2} dz, \quad (2.33)$$

$$\frac{da}{a^2} = \frac{da (1+z)^2}{a^2} = -\frac{dz}{a_0}, \quad (2.34)$$

which means

$$de = -\frac{dz}{a_0 H(z)}. \quad (2.35)$$

??? probably there is wrong stuff here

The Hubble parameter is given by

$$H^2 = \frac{8\pi G}{3} \rho - \frac{kc^2}{a^2}, \quad (2.36)$$

with density $\rho(t) = \rho_r(t) + \rho_m(t) + \rho_\Lambda$, where the first term scales like a^{-4} , the second a^{-3} , the third is constant. Thus they scale like $(1+z)^4$, $(1+z)^3$ and so on.

Then we can write a law for the evolution of $H(z) = H_0 E(z)$. Recall the definition of $\Omega(t)$: we can look at the $\Omega_i(t)$ for i corresponding to matter, radiation and so on:

$$\Omega_i(z) = \frac{8\pi G \rho_i(t)}{3H^2(z)} = \frac{8\pi G}{3H^2} \frac{\rho_i(z)}{E^2(z)} \stackrel{\text{def}}{=} \Omega_{i,0} \frac{(1+z)^\alpha}{E^2(z)}. \quad (2.37)$$

In the case of radiation, $p = \rho c^2/3$, and then $\alpha = 4$.

For matter $p = 0$: $\alpha = 3$.

For the cosmological constant Λ : $\alpha = 0$.

For spatial curvature we have $\alpha = 2$.

For the Ω corresponding to the curvature we define: $\Omega_k = -tc^2/(a^2 H^2)$.

We must have

$$1 = \Omega_r + \Omega_m + \Omega_\Lambda + \Omega_k. \quad (2.38)$$

Recall the definition of $E^2(z)$:

$$E^2(z) = \frac{H^2}{H_0^2} = \Omega_{k,0} + \Omega_{r,0}(1+z)^4 + \Omega_{m,0}(q+z)^3. \quad (2.39)$$

and to get E we just take the square root. We have $\tau(t_0) - \tau(t_e) = \chi(r_e)$. Integrating:

$$\chi(r) = c \int_{a_t}^a \frac{da}{a\dot{a}} = \int \frac{dz'}{E(z')}, \quad (2.40)$$

therefore

$$r = f_k \left(\frac{c}{a_0 H_0} \int_0^z \frac{dz'}{E(z')} \right). \quad (2.41)$$

Two weeks ago we defined the luminosity distance: now we can compute it.

How do we decide which k to use?

Now, suppose we are looking at a certain far-away object with angular size $\Delta\theta$: we fix r in the RW line element, and look at a constant time: then we get a linear size corresponding to the angular one of

$$ds = a(t)r\Delta\theta, \quad (2.42)$$

which, when divided by $\Delta\theta$, is called angular diameter $D_A = ar = a_0 r(z)/(1+z)$. This changes with distance... (?)

$$d_L = a_0(1+z)r, \quad (2.43)$$

then

$$\frac{d_L}{D_A} = (1+z)^2. \quad (2.44)$$

Einstein thought that the universe had to be static.

Recall the Friedmann equations (1.9). Now, if we look at static solutions for matter ($p = 0$): the third equation becomes the identity, the derivatives of a are zero: therefore the second equation gives us $\rho \equiv 0$: there cannot be matter.

Now we know that the universe is neither static nor stationary.

Einstein modified his equations in order to get a static non-empty solution.

The Einstein equations read

$$G_{\mu\nu} = 8\pi G T_{\mu\nu}, \quad (2.45)$$

when $c = 1$, where the Einstein tensor $G_{\mu\nu}$ can be defined with

$$G_{\mu\nu} = R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R. \quad (2.46)$$

Einstein added a term $-\Lambda g_{\mu\nu}$ to the LHS of the Einstein equations. This is allowed since

1. it is symmetric;
2. it has zero covariant divergence, since Λ is constant while $\nabla_\mu g^{\mu\nu} = 0$.

Then, we can rewrite the EE with a modified stress-energy tensor, to which we add $\Lambda g_{\mu\nu}/8\pi G$. Comparing this to an ideal fluid tensor

$$T_{\mu\nu} = \begin{bmatrix} \rho & 0 & 0 & 0 \\ 0 & p & 0 & 0 \\ 0 & 0 & p & 0 \\ 0 & 0 & 0 & p \end{bmatrix}, \quad (2.47)$$

we get $\rho \rightarrow \rho + \Lambda/8\pi G$ and $p \rightarrow p + \Lambda/8\pi G$.

Inserting this into the Friedmann equations we get:

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho + \frac{\Lambda}{3} - \frac{k}{a^2}, \quad (2.48)$$

$$\frac{\ddot{a}}{a} = -\frac{4\pi G}{3}\rho + \Lambda, \quad (2.49)$$

while in the third equation, in the contribution $\tilde{\rho} + \tilde{p}$ the two Λ terms cancel. Then for the first equation we get

$$\frac{8\pi G}{3}\rho + \frac{\Lambda}{3} = \frac{k}{a^2}, \quad (2.50)$$

and for the second:

$$4\pi G\rho = \Lambda. \quad (2.51)$$

Then,

$$\Lambda\left(\frac{1}{3} + \frac{2}{3}\right) = \Lambda = \frac{k}{a^2}, \quad (2.52)$$

and we want a solution with $k = 1$, $\Lambda > 0$.

A candidate for the cosmological constant term is the vacuum energy in QFT: however the estimate given there is around 10^{120} times off.

The next topic is a solution of the Friedmann equations. We will try to do it with $p = 0$: $\rho \propto a^{-3}$.

Fri Oct 18 2019

From yesterday: other consequences of inserting Λ . Now we have a different approach from Einstein's: we insert the constant not in order to get a static universe but just as a measurable parameter of our theory.

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho + \frac{\Lambda}{3} - \frac{k}{a^2}, \quad (2.53)$$

with $\rho = \rho_0(a_0/a)^3$. The more the universe expands, the more the cosmological constant term dominates, since the other terms are inversely dependent on a .

The asymptotic state, neglecting spatial curvature, gives us a steady-state solution:

$$a(t) \propto \exp\left(\sqrt{\frac{\Lambda}{3}}t\right). \quad (2.54)$$

This is called a *de Sitter* solution. An alternative is found if we *do* consider the spatial curvature, with $k = \pm 1$. The two other solutions are also called *de Sitter* and look like cosines (?): they can be mapped into each other.

How?

These belong to the maximally symmetric solutions: Minkowski, dS and Anti de Sitter: the latter has $\Lambda < 0$.

In AdS, do we have $a \propto \exp(ikt)$?

The *no-hair cosmic theorem* is actually a conjecture: it states that asymptotically only the dark energy contribution is relevant: all the matter and everything else is forgotten.

$$a \propto (\sinh(At))^{2/3}, \quad (2.55)$$

where we define $2A/3 = \sqrt{\Lambda/3}$ is a solution which *interpolates* between the current — matter dominated — universe and the asymptotic one, since the hyperbolic sine is locally a simple exponential.

We can rewrite the Friedmann equation as

$$\dot{a}^2 = \frac{8\pi G}{3}\rho a^2 - k, \quad (2.56)$$

and now we will solve it with $k = \pm 1$. In general, for an ODE like $y = f(y')$ with f' continuous we introduce $y' \equiv p$ with $p \neq 0$: then $y = f(p)$, which implies

$$y' = \frac{df}{dp}p', \quad (2.57)$$

and then

$$p = \frac{df}{dp}p' \implies \frac{dx}{dp} = \frac{1}{p} \frac{df}{dp}, \quad (2.58)$$

which gives by integration the solution:

$$x = \int dp \frac{1}{p} \frac{df}{dp} \quad \text{with} \quad y = f(p). \quad (2.59)$$

Using this for our problem, we get $\dot{a}^2 = Aa^{-1} - k$, where $A \equiv 8\pi G a_0^3 \rho_0 / 3$. We can rewrite this as

$$a = \frac{A}{p^2 + k} \quad \text{where} \quad p = \dot{a}. \quad (2.60)$$

Then we get:

$$\dot{a} = p = -2Ap\dot{p} \frac{1}{(p^2 + k)^2}, \quad (2.61)$$

and using our formula

$$t = -2A \int \frac{dp}{(p^2 + k)^2}. \quad (2.62)$$

In order to go forward, we distinguish the cases: if $k = +1$, then $p = \tan(\theta)$, therefore $1 + p^2 = \sec^2 \theta$ which implies $dp = d\theta \sec^2 \theta$.

For the time:

$$t = \int -2A d\theta \cos^2(\theta) = -A(\theta + \sin(\theta) \cos(\theta)) + \text{const}, \quad (2.63)$$

and we apply the trigonometric identity $\sin(\theta) \cos(\theta) = \sin(2\theta)/2$:

$$t = -\frac{A}{2}(2\theta + \sin(2\theta)) + \text{const}, \quad (2.64)$$

now we can define $2\theta = \pi - \alpha$, therefore $\sin(2\theta) = \sin(\alpha)$: this gives us $t = A/2(\alpha - \sin(\alpha))$ and $p = 1/\tan(\alpha/2) = \tan(\pi/2 - \alpha/2)$.

So we almost have our solution:

$$a = \frac{A}{1 + \tan^2(\pi/2 - \alpha/2)} = A \cos^2(\frac{\pi}{2} - \frac{\alpha}{2}) = \frac{A}{2}(1 + \cos(\pi - \alpha)) = \frac{A}{2}(1 - \cos(\alpha)), \quad (2.65)$$

which should be complemented with the equation we found for t .

Reinserting the constants we have:

$$a = a_0 \frac{\Omega_0}{2(\Omega_0 - 1)} (1 - \cos(\alpha)), \quad (2.66)$$

and

$$t = \frac{1}{H_0} \frac{\Omega_0}{2(\Omega_0 - 1)^{3/2}} (\alpha - \sin(\alpha)), \quad (2.67)$$

but now we switch from α to θ for historical reasons.

[Plot: $a(\theta)$ vs θ].

We have $\dot{a} > 0$ when $0 \leq \theta \leq \theta_m = \pi$, while $\dot{a} < 0$ when $\theta_m \leq \theta \leq 2\pi$. This is the *turn-around* angle. The angles 0 and 2π correspond to the Big Bang and the Big Crunch.

At θ_m we have:

$$a = a_0 \frac{\Omega_0}{\Omega_0 - 1}, \quad (2.68)$$

and

$$t = \frac{\pi}{2H_0} \frac{\Omega_0}{(\Omega_0 - 1)^{3/2}}, \quad (2.69)$$

therefore if we set $a = 1$ at the current time we get

$$t_0 = \frac{1}{2H_0} \frac{\Omega_0}{(\Omega_0 - 1)^{3/2}} \left(\arccos\left(\frac{2}{\Omega_0} - 1\right) - \frac{2}{\Omega_0}(\Omega_0 - 1)^{1/2} \right), \quad (2.70)$$

In a pure matter model we would have $a \propto t^{2/3}$, which would imply $H_0 = 2/(3t_0)$ or $t_0 = 2/(3H_0)$: this does not make sense! It would give us an age of the universe of the order 10^9 yr, while the measured age of the universe is 1.4×10^{10} yr. This is for a flat universe.

For a given H_0 , we would have a smaller t_0 with a closed universe.

For $k = -1$ we do exactly the same steps with hyperbolic functions instead of trigonometric ones: we get

$$a(\psi) = a \frac{\Omega_0}{2(1 - \Omega_0)} (\cosh \psi - 1), \quad (2.71)$$

with $\cosh \psi = 2/\Omega_0 - 1$ and

$$t(\psi) = \frac{1}{H_0} \frac{\Omega_0}{2(\Omega_0 - 1)^{3/2}} (\sinh \psi - \psi). \quad (2.72)$$

The same reasoning as before gives us a $t_0 > 2/(3H_0)!$ This is then more attractive.

Radiation's energy density $\rho_r(a)$ can be seen as a function of the redshift: $\rho_r(z) = \rho_{0r}(1+z)^4$, since $(1+z) = a_0/a$.

For matter $\rho_m(z) = \rho_{0m}(1+z)^3$, for Λ instead $\rho_\Lambda(z) = \rho_{0\Lambda}$.

Let us define a moment called the *equality redshift* z_{eq} . This is when $\rho_r(z_{\text{eq}}) = \rho_m(z_{\text{eq}})$. This means that

$$(1 + z_{\text{eq}}) = \frac{\rho_{0,m}}{\rho_{0,r}} = \frac{\Omega_{0,m}}{\Omega_{0,r}}, \quad (2.73)$$

where we divided and multiplied by the critical density.

We know that $\Omega_{0,m}$ is around 0.3, while for the radiation it would be easier to measure the density. Neutrinos are matter now but they were radiation at the time of equality (???).

Accounting for everything, we think that

$$1 + z_{\text{eq}} \simeq 2.3 \times 10^4 \Omega_{0,m} h^2. \quad (2.74)$$

This means that the recombination of electrons and protons into Hydrogen happened when the universe was already *matter dominated*.

Another time is z_Λ , defined by: $\rho_m(z_\Lambda) = \rho_\Lambda(z_\Lambda)$.

$$1 + z_\Lambda = \left(\frac{\rho_{0,\Lambda}}{\rho_{0,m}} \right)^{1/3} \simeq \left(\frac{0.7}{0.3} \right)^{1/3}, \quad (2.75)$$

which implies that $z_\Lambda \approx 0.33$.

2.2 The thermal history of the universe

The model which won out is the *hot Big Bang* model.

Consider radiation: we know from Stefan's law that $\rho_r \propto T^4$, while $\rho_r \propto a^{-4}$. Therefore we would expect *Tolman's law* to hold: $T \propto 1/a$.

By *total number of galaxies* we mean the galaxies in our past light cone.

When (in natural units) we get temperature of the order of a certain elementary particle, then statistically that type of particle will usually be ultra-relativistic.

The number density of particles is:

$$n = \frac{g}{(2\pi)^3} \int d^3q f(\vec{q}), \quad (2.76)$$

in units where $c = \hbar = k_B = 1$. The parameter g is the number of helicity states, q is the three-momentum. The function $f(\vec{q})$ is a pdf in phase space. We will have to integrate over position and momentum: the metric we will get in the end must not depend on anything but time, by the cosmological principle.

There is no general rule for g : for photons, we only have two spin states; the “rule” $g = 2s + 1$ is not actually applied, for photons $\vec{s} = 0$ is unphysical, for gravitons $|\vec{s}| \leq 1$ is unphysical. g accounts for all internal degrees of freedom: for atoms we also have vibration, rotation...

The energy density is

$$\rho = \frac{g}{(2\pi)^3} \int d^3\vec{q} E(q) f(\vec{q}), \quad (2.77)$$

where $E^2 = q^2 + m^2$. For photons $E = q$, for nonrelativistic particles $E \approx m + q^2/2m$.

The adiabatic pressure is

$$P = \frac{g}{(2\pi)^3} \int d^3\vec{q} \frac{q^2 f(\vec{q})}{3E}, \quad (2.78)$$

which comes from a consideration of the diagonal components T_{ii} of the stress energy tensor of particles. Alternatively, we can say that this comes from imposing $dE = P dV + dQ$.

This definition gives us $P = \rho/3$ for photons directly.

In full generality the distribution is

$$f(\vec{q}) = \left(\exp\left(\frac{E - \mu}{T}\right) \pm 1 \right)^{-1}, \quad (2.79)$$

where we have a plus for fermions, and a minus for bosons. Here, μ is the chemical potential: it becomes relevant when the gas becomes hot and dense. It can be introduced as a Lagrange multiplier for changes in number of particles.

The Planck distribution is given by:

$$f_k(\vec{q}) = \left(\exp\left(\frac{q}{T}\right) - 1 \right)^{-1}, \quad (2.80)$$

since they are bosons with no chemical potential:

In general we can say that if for some species we have the reaction $i + j \leftrightarrow k + l$, then $\mu_i + \mu_j = \mu_k + \mu_l$. We can deduce them by the known relations: for example, from the annihilation of electron and positron we can derive $\mu_{e^+} = -\mu_{e^-}$. This rule is not trivial; it is an ansatz of thermodynamical equilibrium to get a solution of the Boltzmann equation which allows us to write the Saha equations.

Fri Oct 25 2019

Do we wish to have a part on stellar astrophysics right now, before going on with cosmology? Let him know.

Are we in thermal equilibrium? In general, no. This must be considered when dealing with CMB anisotropies.

QFT must be dealt with not only at zero temperature, but also at finite temperatures. People started doing this in the seventies.

The time variable is then periodic, with period $2\pi\beta$, where $\beta = 1/(k_B T)$.

Is this connected to imaginary time?

This is “in-in” instead of “in-out”: there are no equilibrium states before or after the interaction.

We will simplify: we assume thermal equilibrium at any time in the evolution. A key point in cosmology is the *absence* of time translation invariance, therefore energy is not conserved.

This part will come from Weinberg’s book.

We use units where $c = k_B = \hbar = 1$. The number density of a certain species is

$$n(\dots) = \frac{g}{(2\pi)^3} \int d^3q f(\vec{q}), \quad (2.81)$$

where $f(\vec{q})$ is the phase space distribution of the particles, while g is the number of helicity states of that species. The dots will be explained later: they are about the parametrization of the phase space distribution.

For any species, $E = \sqrt{m^2 + |\vec{q}|^2}$: this applies *on shell*, for the classical equations of motion: when there are quantum fluctuations it does not hold. The energy density is

$$\rho(\dots) = \frac{g}{(2\pi)^3} \int d^3q E(\vec{q}) f(\vec{q}), \quad (2.82)$$

while the pressure is

$$P(\dots) = \frac{g}{(2\pi)^3} \int d^3q f(\vec{q}) \frac{q^2}{3E}. \quad (2.83)$$

Because of isotropy, in the cases we need to consider, the dependence on \vec{q} is actually a dependence on $|\vec{q}| \equiv q$.

What is our ansatz for the phase space distribution?

$$f(q) = \frac{1}{\exp\left(\frac{E-\mu}{T}\right) \mp 1}, \quad (2.84)$$

where the sign is $-$ for bosons, $+$ for fermions.

The chemical potential deals with the flux of particles. We recover the Planck distribution when $E = q$ and $\mu = 0$: this tells us that photons do not have any chemical potential. If we measured $\mu \neq 0$ that would be called a *spectral distortion*, but it seems like that is not the case.

The rule for the sum of the chemical potentials only holds at equilibrium.

We can relate some chemical potentials by reactions, and this tells us about the conserved quantities which follow from the symmetry group of our theory.

If there was a global charge, phenomenologically we’d expect global magnetic fields, but we see them only of the order of the $10 \mu\text{G}$.

So we have an upper bound on the global charge density (?)

We can estimate the orders of magnitude for the various species in the universe.

In the end: very often, $\mu/T \ll 1$. Therefore, the dependence of the number density, energy density and pressure would be also on μ but we forget it, and keep only the temperature dependence: we write $n(T)$, $\rho(T)$ and $P(T)$.

From the second principle of thermodynamics we know that the entropy in a certain volume V at temperature T , $S(V, T)$ is given by:

$$dS = \frac{1}{T}(d\rho(T, V) + P(T) dV). \quad (2.85)$$

Then we have

$$\frac{\partial S}{\partial V} = \frac{1}{T}(\rho(T) + P(T)), \quad (2.86)$$

and

$$\frac{\partial S}{\partial T} = \frac{V}{T} \frac{d\rho(T)}{dT}. \quad (2.87)$$

Recall the Pfaff relations: in order for the differential to be exact it needs to be closed, which means that the second partial derivatives need to commute:

$$\frac{\partial}{\partial T} \left(\frac{1}{T}(\rho(T) + P(T)) \right) = \frac{\partial}{\partial V} \left(\frac{V}{T} \frac{d\rho(T)}{dT} \right). \quad (2.88)$$

We can proceed with the calculation, to get:

$$\frac{dP}{dT} = \frac{1}{T}(\rho(T) + P(T)). \quad (2.89)$$

Cosmology has not entered into the picture yet, but it can by the third Friedmann equation, which can be rewritten as

$$a^3 \frac{dP}{dt} = \frac{d}{dt} (a^3(P + \rho)), \quad (2.90)$$

and these two, when put together, are equivalent (check!) to

$$\frac{d}{dt} \left(\frac{a^3}{T}(\rho + P) \right) = 0, \quad (2.91)$$

therefore this quantity is a constant of motion.

(For the RW line element, the square of the determinant $\sqrt{-g} = a^3$, so we get

$$\frac{d}{dt} \left(\sqrt{-g} \frac{\rho + P}{T} \right) = 0, \quad (2.92)$$

)

So the quantity which is differentiated is constant. If we plug this back into the differential expression for the entropy, we get:

$$dS = d \left(\frac{(\rho + P)V}{T} \right), \quad (2.93)$$

therefore the differentiated quantities are equal up to a constant, but since $V \propto a^3$ we get that the *entropy is constant*.

$$S \equiv S(a^3, T) = \frac{a^3}{T} (\rho + P) = \text{const.} \quad (2.94)$$

Since $\rho \propto T^4$ and $\rho \propto a^{-4}$, we expect Tolman's law: $Ta \sim 1$.

Also, for photons $P = \rho/3$: so we get $S \propto \frac{1}{T} 43(a^3/T) T^4 \propto T^3 a^3 = \text{const.}$

We only consider photons since they have a much larger number density.

Is this because the suppression is exponential while the ratio of energies is somehow polynomial?

Now we wish to do the integrals for the thermodynamical quantities:

$$\int d^3\vec{q} = 4\pi \int_0^\infty dq q^2, \quad (2.95)$$

so we get

$$n(T) = \frac{g}{(2\pi^2)} \int dq q^2 f(q) \quad (2.96a)$$

$$\rho(T) = \frac{g}{2\pi^2} \int dq q^2 f(q) E(q) \quad (2.96b)$$

$$P(T) = \frac{g}{6\pi^2} \int dq q^2 f(q) \frac{q^2}{E(q)}, \quad (2.96c)$$

which can be solved analytically in the ultrarelativistic and nonrelativistic approximations.

Ultrarelativistic means $q \gg m$: actually we compare these to the temperature: we define $x = q/T$, and then the energy is $\sqrt{x^2 + m^2}/T^2$: we suppose $m/T \ll 1$.

The momentum will not always be large, but the cases in which it is large give a much greater contribution if the temperature is large.

Ultrarelativistic We approximate $E(q) \sim q$, so

$$n(T) = \frac{g}{2\pi^2} \int_{\mathbb{R}^+} dq q^2 \left(\exp(q/T) \mp 1 \right)^{-1}, \quad (2.97)$$

$$\rho(T) = \frac{g}{2\pi^2} \int_{\mathbb{R}^+} dq q^3 \left(\exp(q/T) \mp 1 \right)^{-1}, \quad (2.98)$$

$$P(T) = \frac{g}{6\pi^2} \int_{\mathbb{R}^+} dq q^3 \left(\exp(q/T) \mp 1 \right)^{-1}, \quad (2.99)$$

so we can see that in this approximation, which is equivalent to $m \approx 0$, we get matter behaving like radiation: $P = \rho/3$. The result of the integrals depends on the type of the particles, and it is

$$n(T) = \begin{cases} \frac{\zeta(3)}{\pi^2} g T^3 & BE \\ \frac{3}{4} \frac{\zeta(3)}{\pi^2} g T^3 & FD \end{cases}, \quad (2.100)$$

and we get the proportionality to T^3 . For the energy density:

$$\rho(T) = \begin{cases} \frac{\pi^2}{30} g T^4 & BE \\ \frac{7}{8} \frac{\pi^2}{30} g T^4 & FD \end{cases}, \quad (2.101)$$

and for the pressure we just divide by 3. For photons, $g = 2$ and they are bosons, so we get

$$\rho_\gamma = \frac{\pi^2}{15} T^4. \quad (2.102)$$

Nonrelativistic Now $m \gg T$, so we get $E \approx m + \frac{q^2}{2m} + O((q/\sqrt{m})^4)$ by expanding.

So, can we just substitute this? It would seem like this makes the exponential very large, so the difference between bosons and fermions becomes negligible. So, to zeroth order in (q^2/m) we get

$$f \approx \exp\left(-\frac{m - \mu}{T}\right), \quad (2.103)$$

and then

$$n = \frac{g}{2\pi^2} \exp\left(-\frac{m - \mu}{T}\right) \int_{\mathbb{R}^+} dq q^2, \quad (2.104)$$

which diverges. This is the ultraviolet catastrophe: we need the *Planckian*, by considering the first order in q^2/m : then

$$n = \frac{g}{2\pi^2} \exp\left(-\frac{m - \mu}{T}\right) \int_{\mathbb{R}^+} dq q^2 \exp\left(-\frac{q^2}{2mT}\right) = g \left(\frac{mT}{2\pi}\right)^{3/2} \exp\left(\frac{\mu - m}{T}\right), \quad (2.105)$$

where we applied the identity

$$\int_{\mathbb{R}} dx x^2 \exp(-\alpha x^2) = \frac{\sqrt{\pi}}{2\alpha^{3/2}}. \quad (2.106)$$

To recover the energy density and pressure we could do $\rho = mn$ and $P = nT$.

Do these come out of doing the other two integrals?

Therefore, $P = T\rho/m$, so $P \ll \rho$: for nonrelativistic particles, we can deal with them as *noninteracting dust*.

$P = nT$ is just the ideal gas law.

If we compare relativistic particles to nonrelativistic ones, the former dominate the latter in terms of all of these three quantities.

Thu Oct 31 2019

Neutrinos have very low mass: therefore they become relativistic very quickly.

We saw last time the case of either bosons' or fermions' number density, energy density and pressure.

In the nonrelativistic case instead we must consider the second order in q/\sqrt{m} , and we get the Boltzmann suppression factor out of the integral: $\exp(-m/T)$.

In the nonrelativistic case there is no distinction between bosons and fermions.

What happens in the part of the history of the universe where all the particles were in thermal equilibrium?

If the timescales of the interactions are much larger than the cosmological events timescales (such as the energy of the universe), then those interactions do not happen. These are called *decoupled*.

Let us consider ultrarelativistic particles which are not *decoupled*, in the early universe which is radiation dominated (here “radiation” refers to all kinds of ultrarelativistic particles).

In this case, from the “conservation of the stress-energy tensor” we know that $\rho \propto a^{-4}$.

Our equation is

$$H^2 = \frac{8\pi G}{3}\rho - \frac{k}{a^2}, \quad (2.107)$$

and we want to neglect the last term. The equation becomes

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho_* \left(\frac{a}{a_*}\right)^{-4} - \frac{k}{a^2}, \quad (2.108a)$$

but it is not enough to look at the slopes: we can, however, have information about the normalization as well from present-day observations. If the second term is much smaller than the first today, then it was even more so in the far past. We can rewrite the equation as

$$1 = \Omega_{\text{tot}} - \frac{k}{a^2 H^2} \equiv \Omega_{\text{tot}} + \Omega_{\text{curvature}}, \quad (2.109)$$

where $\Omega_{\text{curvature}} \equiv -k/(a^2 H^2)$.

So we neglect the second term: approximately we then have

$$\left(\frac{\dot{a}}{a}\right)^2 = \frac{8\pi G}{3}\rho_{\text{rad}}. \quad (2.110)$$

We know that $a(t) \propto t^{1/2}$, therefore $\dot{a}/a = H = 1/2t$.

Wait, how does this work?

So, we have

$$\frac{1}{4t^2} = \frac{8\pi G}{3}g_* \frac{\pi^2}{30}T^4, \quad (2.111)$$

where

$$g_* = g_*(T) = \sum_{i \in BE} g_i \left(\frac{T_i}{T}\right)^4 + \frac{7}{8} \sum_{i \in FD} g_i \left(\frac{T_i}{T}\right)^4, \quad (2.112)$$

where we insert a correction factor in order to not consider particles which are nonrelativistic (?). *BE* and *FD* denote the bosonic and fermionic degrees of freedom respectively.

The i spans the different types of particles in the SM. The T_i are the thermalization temperatures of the various species.

The plot of g_* in terms of the temperature goes “down in steps”, as more and more species become nonrelativistic.

Then we have a formula for temperature in terms of time:

$$\frac{1}{2t} = \left(\frac{8\pi G}{3}\right)^{1/2} g_*^{1/2} \left(\frac{\pi^2}{30}\right)^{1/2} T^2, \quad (2.113)$$

therefore we get:

$$t \approx 0.301 g_*^{-1/2} \frac{m_P}{T^2} \approx \left(\frac{T}{\text{MeV}}\right)^{-2} s, \quad (2.114)$$

where $m_P = G^{-1/2}$ is the Planck mass. Beware: there are different conventions for this! The factor $g_*^{-1/2}$ is of order 1 around $T \approx 1 \text{ MeV}$, so we ignore it.

The number is found by:

$$0.301 \approx \frac{1}{2\sqrt{\frac{8\pi\pi^2}{3 \times 30}}}. \quad (2.115)$$

The Planck mass is approximately $1.2 \times 10^{19} \text{ GeV}$. It is the scale at which we need to account for quantum gravitational effects.

When the universe is 1 second old, weak interactions decouple.

Why are we allowing different thermalization temperatures?

A hypothesis is that the particles exit the Planck epoch, $t \approx m_P$, thermalized. When decoupling occurs, they can stop being thermalized.

Another hypothesis is that they stop being thermalized at a certain temperature during inflation.

Some time ago we discussed entropy conservation.

The entropy density is given by

$$s \equiv S/V = (P + \rho)/T = \frac{4}{3} \frac{\rho}{T} = (2\pi^2/45) g_{*s} T^3; \quad (2.116)$$

since $V \propto a^3$ we have $sa^3 = \text{const}$.

We defined a new g_* to account for the different species:

$$g_{*s} \equiv \sum_{i \in BE} g_i \left(\frac{T_i}{T}\right)^3 + \frac{7}{8} \sum_{i \in FD} g_i \left(\frac{T_i}{T}\right)^3, \quad (2.117)$$

so our new Tolman's law is $Ta g_{*s}^{1/3} = \text{const}$.

When the neutrinos decouple, their temperature keeps scaling like $1/a$: they keep being thermalized with the photons even though they do not react.

When we reach $T \approx 0.5 \text{ MeV}$ electrons and positrons become thermalized: they decay into photons: the photon temperature T increases with respect to what would happen without this process (it actually just decreases slower).

The neutrinos are not updated when this process happens: their temperature then becomes consistently less than the one of the photons, but they keep scaling the same: one of their temperatures is a constant multiple of the other after this event.

Now we will calculate this multiple.

First, at 1 MeV, neutrinos decouple. Then, at 0.5 MeV, electrons and positrons decouple.

We require continuity at this transition: $T_>$ and $T_<$ must both be T_e .

$$sa^3 = \frac{2\pi^2}{45} g_{*s>} T_>^3 a_>^3 = \frac{2\pi^2}{45} g_{*s<} T_<^3 a_<^3, \quad (2.118)$$

but we can drop all the factors we know to be continuous on the boundary:

isn't a different on either side of the boundary, if it is not instant?

we are left with

$$T_< = \left(\frac{g_{*s>}}{g_{*s<}} \right)^{1/3} T_<, \quad (2.119)$$

where we can compute

$$g_{*s>} = 2 + \frac{7}{8}(4) = \frac{11}{2}, \quad (2.120)$$

where we considered photons, electrons but not neutrinos (since they are separated). On the other hand,

$$g_{*s<} = 2, \quad (2.121)$$

since electrons are not thermalized anymore. Therefore

$$T_< = \left(\frac{11}{4} \right)^{1/3} T_>, \quad (2.122)$$

which allows us to compute the neutrino temperature at any time:

$$T_\nu = T_\gamma \left(\frac{4}{11} \right)^{1/3}, \quad (2.123)$$

since they scale the same.

Now let us compute $\rho_r(T = 0.1 \text{ MeV})$. Let us assume that the global temperature is the one of the photons: we get

$$g_* = \sum_{i \in BE} g_i \left(\frac{T_i}{T} \right)^4 + \frac{7}{8} \sum_{i \in FD} g_i \left(\frac{T_i}{T} \right)^4 \quad (2.124a)$$

$$= 2 + \frac{7}{8} \left(3 \times 2 \left(\frac{T_\nu}{T_\gamma} \right)^4 \right) \quad (2.124b)$$

$$= 2 + \frac{21}{4} \left(\frac{4}{11} \right)^{4/3} \approx 3.36, \quad (2.124c)$$

since we need to consider neutrinos, which contribute to the total energy density, but not electrons which are not relativistic.

Last hour we discussed the Planck mass $m_P \approx 1.2 \times 10^{19}$ GeV: it also defines a wavelength, $\lambda_P \approx 10^{-33}$ cm and a timescale $t_P \approx 10^{-43}$ s (both of these are $1/m_P$ in natural units).

When the age of the universe is of this order, our theories are not guaranteed to work.

2.2.1 The cosmological horizon

Let us consider null geodesics in a De Sitter universe with a RW metric: we can take them as radial, and find

$$c^2 dt^2 = a^2(t) \frac{1}{1 - kr^2} dr^2, \quad (2.125)$$

therefore we get:

$$\int^t \frac{c dt}{a(t)} = \int^r \frac{d\tilde{r}}{\sqrt{1 - k\tilde{r}^2}} = f(r), \quad (2.126)$$

but in order to get something which has the dimensions of a length we need to multiply by a :

$$d_{\text{Hor}}(t) = a(t) \int_0^t \frac{c d\tilde{t}}{a(\tilde{t})}. \quad (2.127)$$

“If this integral is convergent, we should be worried”.

If we integrate from the beginning of time to now, we get the spatial (current) distance elapsed by a photon which started at the start of time. This is the radius of the largest region we could in principle observe. It is of the order of 3 Gpc.

This is not the case, for example, in Minkowski spacetime.

If there is an end of time, there is also a future horizon.

The quantity $d_{\text{Hor}}(t)$ is increasing in time: this causes the Cosmological Horizon problem.

What is the problem?

An important moment is the *recombination of hydrogen*: the formation of the first Hydrogen atoms, so the first moment at which Compton scattering can occur. This causes the decoupling of radiation and baryonic matter. This is the moment at which the radiation in the CMB was emitted.

The CMB is very close to being uniform. It was emitted something like $t = 3.8 \times 10^5$ yr after the BB.

Let us consider the past light cones from two points diametrically opposite with respect to us, at this point in time: they do not overlap, so the CMB cannot be causally correlated. However, it seems like it is!

What is the angular scale at which the light cones overlap?

We have a classification by Bianchi of non-isotropic universes (in 9 classes).

There is the Mixed Master Universe.

Let us consider

$$R_{\text{Hor}} = \int_0^t \frac{c \, d\tilde{t}}{a(\tilde{t})}. \quad (2.128)$$

We can hypothesize that there was a period where the comoving radius was decreasing with time. This would solve the problem. We can approximate

$$d_{\text{Hor}}(t) = a(t) \int_0^t \frac{c \, d\tilde{t}}{a(\tilde{t})} \approx ct \sim \frac{c}{H} \equiv d_H, \quad (2.129)$$

where we defined the new *Hubble distance*. We also define the comoving Hubble radius: $r_H = c/(Ha) = c/\dot{a}$.

For it to be decreasing, the condition is

$$\dot{r}_H = -\frac{\ddot{a}}{\dot{a}^2} < 0, \quad (2.130)$$

therefore we need $\ddot{a} > 0$ for at least some time. We know that

$$\ddot{a} = -\frac{4\pi G}{3} \left(\rho + \frac{3P}{c^2} \right), \quad (2.131)$$

therefore we need to have $\rho + 3P/c^2 < 0$. So, since the energy density is positive, the condition is $P < -\rho/3$.

Let us consider the parameter \ddot{a} : it is

$$\ddot{a} = \dot{a}H + a\dot{H} = a(H^2 + \dot{H}) > 0, \quad (2.132)$$

so the condition is $H^2 + \dot{H} > 0$. In terms of $P/\rho = w$ we have:

1. $\dot{H} < 0$ while $H^2 + \dot{H} > 0$: this corresponds to $-1 < w < -1/3$;
2. $\dot{H} = 0$: this is De Sitter: $a(t) = \exp(Ht)$, corresponding to $w = -1$;
3. $\dot{H} > 0$: here the solution is

$$a(t) = a_* \left(1 + \frac{3}{2} ((1+w)H_*(t-t_*)) \right)^{2/(3(1+w))}, \quad (2.133)$$

which corresponds to $w < -1$ and means $a(t) \propto t^p$ for some $p > 1$. This is called power-law inflation.

Big Rip singularity: with $a(t) = |t - t_{\text{as}}|^{-\alpha}$ with $\alpha > 0$. This is called poli-inflation.

What is this about?

The third condition is very hard to achieve. The boundary at $w = -1$ is called the *phantom divide*.

Now, let us consider the *flatness problem*: this was first proposed by Dick and Peebles in 1986.

The parameter $\Omega = \frac{8\pi G\rho(z)}{3H^2(z)}$ diverges from 1 as time increases. Measurements of Ω_{tot} gave approximately 0.1.

How is it possible that the universe is still so flat, even when the universe is so old? Oldness and flatness seem incompatible.

This is a type of *fine-tuning* problem. Typically, if there is a fine-tuning problem then it signals that we should improve our theory.

Let us assume that w is constant. Then, $\rho(z) = \rho_0(1+z)^{3(1+w)}$. Recall that

$$H^2(z) = \frac{8\pi G}{3}\rho(z) - \frac{k}{a^2} \quad (2.134a)$$

$$H_0^2 = \frac{8\pi G}{3}\rho_0 - \frac{k}{a_0^2}, \quad (2.134b)$$

the latter of which implies $1 = \Omega - k/(a_0^2 H_0^2)$. So,

$$H^2(z) = H_0^2 \left(\frac{8\pi G}{3H_0^2}\rho(z) - \frac{k}{a^2 H_0^2} \right) \quad (2.135a)$$

$$= H_0^2 \left(\frac{\rho(z)}{\rho_0} \Omega_0 \frac{a_0^2}{a} (1 - \Omega_0) \right) \quad (2.135b)$$

$$= H^2(z) = H_0^2(1+z)^2 \left(\Omega_0(1+z)^{1+3w} + (1 - \Omega_0) \right), \quad (2.135c)$$

so in the end

$$\Omega(z) = \frac{8\pi G\rho_0}{3H_0^2} \frac{\rho(z)}{\rho_0} \frac{H_0^2}{H^2(z)} \quad (2.136a)$$

$$= \Omega_0(1+z)^{1+3w} \left(1 - \Omega_0 + \Omega_0(1+z)^{1+3w} \right), \quad (2.136b)$$

and since $\rho(z) = \rho_0(1+z)^{3(1+w)}$ we get

$$\Omega^{-1}(z) - 1 = \frac{\Omega_0 \left((1 - \Omega_0) + \Omega_0(1+z)^{1+3w} \right) - (1+z)^{1+3w}}{(1+z)^{1+3w}} \quad (2.137a)$$

$$= (\Omega_0^{-1} - 1)(1+z)^{-(1+3w)}. \quad (2.137b)$$

If we assume $w = 1/3$ for all times (which is false, but we do it to get a result that is close enough) we get

$$\Omega^{-1}(z) - 1 = (\Omega_0^{-1} - 1) \left(\frac{T_0}{T(z)} \right)^2. \quad (2.138)$$

If we compute T_{Planck}/T_0 we get approximately 4.5×10^{32} . Then when squaring we get 10^{-64} , without the approximation $w = 1/3$ we get 10^{-60} .

Then, we see that there is something deeply unnatural in the Friedmann model.

Thu Nov 07 2019

We talk about inflation again. The comoving horizon increases with time:

$$d_H(t) = a(t) \int_0^t \frac{c \, d\tilde{t}}{a(\tilde{t})} \sim ct \sim \frac{c}{H} \equiv \text{Hubble horizon}, \quad (2.139)$$

it is also called the *past event horizon*. The reason why we can plot the history of the universe in a single spacetime diagram is because there is a well-defined transformation which brings an infinite interval to a finite one, using a hyperbolic arctangent: this gives us a *Penrose diagram*.

The comoving Hubble radius is $r_H = c/aH = c/\dot{a}$. This can grow.

If we have positive pressure $p > 0$, then the scale factor goes like $a \propto t^{2/(3(1+w))}$ with $w > 0$, then the comoving radius increases with time.

We can actually still get increasing comoving radii with a weaker condition: $p > -\frac{1}{3}\rho c^2$. This is the actual boundary (it can be checked looking at the derivative of a).

This is directly connected to the sign of the acceleration in the Friedmann equation.

If these conditions are always met, Hawking and Ellis proved that a Big Bang is inevitable.

The inflation hypothesis is that, even though now we have $\ddot{a} < 0$ now (or, it was so for some time: now the expansion seems to be accelerating), there was a period in which we had $\ddot{a} > 0$.

These are drawn as straight lines, but it is only qualitative: we are looking at the sign of the slope.

Then, there was a time in the past at which the comoving horizon was as large as it is now: now we will see how much inflation there must have been in order to solve the horizon problem up to now, ignoring the fact that now the universe's expansion is accelerating.

The inequality we want to impose is

$$r_H(t_i) \geq r_H(t_0), \quad (2.140)$$

where t_0 is now while t_i is the beginning of inflation.

A sphere with comoving radius $d_H(t_i)$ will expand after inflation up to

$$d_H(t_i) \frac{a(t_f)}{a(t_i)}, \quad (2.141)$$

where t_f is the time of the end of inflation.

$$d_H(t_i) \frac{a(t_f)}{a(t_i)} \geq d_H(t_0) \frac{a(t_f)}{a(t_0)}. \quad (2.142)$$

We want to see what the limiting condition is.

$$Z_{\min} = \frac{d_H(t_0)}{d_H(t_i)} \frac{a_f}{a_0} = \frac{H_i}{H_0} \frac{a_f}{a_0}, \quad (2.143)$$

is Z a redshift?

$$z_{\min} \frac{H_i a_f}{H_0 a_0} = \frac{H_i}{H_f} \frac{H_f a_f}{H_0 a_0}, \quad (2.144)$$

or

$$\frac{H_f}{H_i} z_{\min} = \frac{H_f a_f}{H_0 a_0}, \quad (2.145)$$

in which we can insert our solution to the Friedmann equations, for the scale factor and Hubble parameter in function of time:

$$H(t) = H_* \left(\frac{a(t)}{a_*} \right)^{-\frac{3(1+w)}{2}}. \quad (2.146)$$

This can be found using the results we found some time ago: the expressions for a and H were equal up to a different thing multiplying the parenthesis, and a different exponent.

This is of course an approximation, but it works. A better number can be found by integrating numerically over different more realistic equations of state. We find:

$$z = \frac{a_f}{a_i} =, \quad (2.147)$$

Put earlier

$$\frac{H_f}{H_i} = z_{\min}^{-\frac{3(1+w)}{2}}, \quad (2.148)$$

therefore we get

$$z_{\min} = \left(\frac{H_f a_f}{H_0 a_0} \right)^{-\frac{2}{(1+3w_{\inf})}}, \quad (2.149)$$

where w_{\inf} is calculated at the time of matter-radiation equality.

So we get:

$$\frac{H_f}{H_0} = \frac{H_f}{H_{\text{eq}}} \frac{H_{\text{eq}}}{H_0} = \left(\frac{a_f}{a_{\text{eq}}} \right)^{-2} \left(\frac{a_{\text{eq}}}{a_0} \right)^{-3/2} = \left(\frac{a_f}{a_0} \right)^{-2} \left(\frac{a_0}{a_{\text{eq}}} \right)^{-1/2}, \quad (2.150)$$

which means that the minimum inflation redshift must be

$$z_{\min} = \left(\left(\frac{a_f}{a_0} \right)^{-1} \left(\frac{a_0}{a_{\text{eq}}} \right)^{1/2} \right)^{\frac{-2}{1+3w_{\inf}}}, \quad (2.151)$$

so the result can be expressed in terms of temperatures:

$$\frac{a_0}{a_f} = \frac{T_f}{T_0} = \frac{T_f}{T_{\text{pl}}} \frac{T_{\text{pl}}}{T_0}, \quad (2.152)$$

where the T_{pl} is the Planck temperature, and $a_0/a_{\text{eq}} = 1 + z_{\text{eq}}$: in the end our result is

$$z_{\text{min}} = \left(\frac{T_{\text{pl}}}{T_0} (1 + z_{\text{eq}}) \frac{T_f}{T_{\text{pl}}} \right)^{-\frac{2}{1+3w_{\text{inf}}}}. \quad (2.153)$$

Recall that 1 GeV is equal to 10^{13} K, and $T_{\text{pl}} = 10^{19}$ GeV. Also, $1 + z_{\text{eq}} = 2.3 \times 10^4 \Omega h^2$

What are these units?

We get

$$z_{\text{min}} \approx 10^{30} \frac{T_f}{T_{\text{pl}}}, \quad (2.154)$$

but what is the early universe temperature at the end of inflation? It must allow baryogenesis, but will still be less than one, but there is an upper bound based on the fact that we have not observed primordial gravitational waves from this time: it must be at most 10^{-3} , so we find that the minimum redshift is of the order $z_{\text{min}} \lesssim 10^{30} \sim e^{60}$, or 60 e -folds.

This is an order of order of magnitude estimate.

From the Friedmann equation we get

$$1 = \Omega(t) - \frac{kc^2}{a^2 H^2}, \quad (2.155)$$

so $\Omega(t) - 1 = kr_H^2$. Now, consider the Ω_i of inflation: we get

$$\frac{\Omega - 1}{\Omega_i - 1} = \left(\frac{r_{H0}}{r_{Hi}} \right)^2 < 1. \quad (2.156)$$

What we discuss now might be outside of our possibilities of comprehension.

Inflation is equivalent to $\ddot{a} > 0$, which is equivalent to $p < -\frac{1}{3}\rho c^2$:

$$\ddot{a} = -\frac{8\pi G}{3} \left(\rho + \frac{3p}{c^2} \right) a. \quad (2.157)$$

A quantum Hamiltonian for a harmonic oscillator is

$$H = \frac{1}{2} \sum \omega (a^\dagger a + a a^\dagger), \quad (2.158)$$

but this might give infinite energy for the ground state. In nonrelativistic QM we know that the energy is defined up to a constant, but in GR this is not the case: this energy gravitates!

Another way this comes up is the Casimir effect: we have virtual particles, which can pop up for times satisfying $\Delta E \Delta t \sim \hbar$. If we put two metallic plates close to each other, we get a force.

In QFT, we either have scalars, vectors or spinors. Can a scalar field have a nonzero expectation value, while respecting the Robertson-Walker symmetries? Yes, we just take a function of time.

For a vector, we cannot have nonzero expectation: a nonzero expectation value gives us a preferred direction. For a spinor, the same holds.

However, an object like $\bar{\psi}\psi$ behaves like a scalar, even though it comes from a vector.

There are almost no scalar particles in nature! The only one is the Higgs field.

The action for GR is given by

$$S = S_{\Phi} + S_{g_{\mu\nu}} + S_{\text{world}}, \quad (2.159)$$

where

$$S = \int d^4x \sqrt{-g} \mathcal{L}, \quad (2.160)$$

and the gravitational Lagrangian is $\mathcal{L}_g = R/16\pi G$. A kinetic Lagrangian is

$$\mathcal{L} = \frac{m}{2} \dot{q}^2 - V(q), \quad (2.161)$$

for a scalar field its equivalent would be

$$\mathcal{L} = \frac{1}{2} \left(\partial_{\mu} \Phi \right)^2 - V(\Phi), \quad (2.162)$$

in Minkowski spacetime. It then becomes:

$$\mathcal{L} = \frac{1}{2} g^{\mu\nu} \nabla_{\mu} \Phi \nabla_{\nu} \Phi - V(\Phi), \quad (2.163)$$

but the covariant derivative of a scalar is just its partial derivative.

If we add a massive term, proportional to $R\Phi^2$, we get that adding it to the global action looks like gravity.

Fri Nov 08 2019

We continue the discussion from yesterday on the dynamics of inflation.

The Lagrangian for a scalar field in GR is

$$\mathcal{L} = \frac{1}{2} g^{\mu\nu} \nabla_{\mu} \Phi \nabla_{\nu} \Phi - V(\Phi). \quad (2.164)$$

The “contravariant derivative” does not exist.

Why?

We can add a term $\zeta R\Phi^2$, which has the right dimensions. Actions are dimensionless since $\hbar = 1$, and since $ds^2 = g_{\mu\nu} dx^{\mu} dx^{\nu}$ the metric is also dimensionless. Therefore the dimensional analysis of $\int d^4x \sqrt{-g} \mathcal{L}$ gives us that \mathcal{L} must have dimensions of m^{-4} .

The field Φ has the dimensions of a mass, which is an inverse length. The coupling constants are conventionally taken to be dimensionless: therefore if we are to add a term to the Lagrangian, it must be $\zeta\Phi^2$ times an inverse square length: so we can insert R .

The value of ζ is undetermined: [16] gives us conformal symmetry, while in other cases we get [14]. A Weyl transformation allows us to remove the additional term: we move from the Jordan frame (where we *do* have coupling between our scalar field and the curvature) to the Einstein frame, in which we do not.

This is a prototype for modified GR theories.

Varying the Einstein-Hilbert action with respect to the metric gives the LHS of the Einstein equations. If we vary with respect to something else we get the equations of motion of that other thing.

This means that the stress-energy tensor is just the functional variation of everything but the action of the metric in the global action, with respect to the metric.

We get for our scalar field:

$$T_{\mu\nu}(\Phi) = \Phi_{,\mu}\Phi_{,\nu} - g_{\mu\nu}\left(\frac{1}{2}g^{\rho\sigma}\Phi_{,\rho}\Phi_{,\sigma} - V(\Phi)\right), \quad (2.165)$$

and then $G_{\mu\nu} = 8\pi GT_{\mu\nu}$.

We can get an explicit solution by using the symmetries of our spacetime: we assume that $\Phi(x^\mu) = \varphi(t)$. This however in QFT is an operator, but we cannot have an operator on the RHS of the EFE: so we do a semiclassical theory, an equivalent of Hartree-Fock: we take an average in the vacuum state, and our equations become

$$G_{\mu\nu} = 8\pi G \left\langle \hat{T}_{\mu\nu} \right\rangle_0, \quad (2.166)$$

where we define the ground state as that one with the most symmetry allowed.

The symmetries we must consider are only rotations and translations. There are no issues of commutation, since we do not quantize space unlike the quantum loop gravity people.

If we perturb, we get $\Phi = \varphi + \delta\Phi$: so $\langle \Phi^2 \rangle = \varphi^2 + 2\langle \varphi\delta\Phi \rangle + \langle \delta\Phi^2 \rangle$, but the second term is zero since $\langle \delta\Phi \rangle = 0$ and φ is constant. The last term in these diverges. We do not know how to deal with it. We therefore assume that it is small.

So: when computing the stress energy tensor we get only diagonal terms: a perfect fluid!

The energy density is the Hamiltonian:

$$\rho = T_{00} = \frac{1}{2}\dot{\varphi}^2 + V(\varphi) = H, \quad (2.167)$$

while the pressure is the Lagrangian:

$$P = \frac{1}{2}\dot{\varphi}^2 - V(\varphi) = \mathcal{L}. \quad (2.168)$$

We call the stuff in the universe which is not our field “radiation”, with energy density ρ_r .

$$H^2 = \frac{8\pi G}{3} \left(\frac{1}{2}\dot{\varphi}^2 + V + \rho_r \right) \quad (2.169a)$$

$$\frac{\ddot{a}}{a} = -\frac{8\pi G}{3}(\dot{\phi}^2 - V + \rho_r) \quad (2.169b)$$

$$\dot{\rho}_{\text{tot}} = -3\frac{\dot{a}}{a}(\rho_{\text{tot}} + P_{\text{tot}}), \quad (2.169c)$$

but in the continuity equation we can split the contributions by inserting an unknown factor Γ , the transfer of energy between the field and radiation.

$$\dot{\rho}_\phi = -3\frac{\dot{a}}{a}\dot{\phi}^2 + \Gamma \quad (2.170a)$$

$$\dot{\rho}_r = -4\frac{\dot{a}}{a}\rho_r - \Gamma, \quad (2.170b)$$

which, denoting $' = \partial_\phi$:

$$\dot{\rho}_\phi = \dot{\phi}\ddot{\phi} + V'\dot{\phi} \quad (2.171a)$$

$$\ddot{\phi}\dot{\phi} + V'\dot{\phi} = -3\frac{\dot{a}}{a} + \Gamma, \quad (2.171b)$$

but we drop Γ since we assume there is little radiation.

One solution is $\dot{\phi} = 0$, if not:

$$\ddot{\phi} + 3\frac{\dot{a}}{a} = -V', \quad (2.172)$$

recall the definition of

$$w = -\frac{1}{3} = \frac{P}{\rho} = \frac{\frac{1}{2}\dot{\phi}^2 - V}{\frac{1}{2}\dot{\phi}^2 + V}, \quad (2.173)$$

so one possibility we have is

$$\dot{\phi}^2 \gg 2|V| \implies w = 1, \quad (2.174)$$

or else

$$\dot{\phi}^2 \ll 2|V| \implies w = -1. \quad (2.175)$$

The continuity equation gives us the Klein-Gordon equation again: it is tautological.

$\phi = \text{const}$ was one of the first solutions proposed. This model seems so fit the data.

Several proposals were made in the late seventies, early eighties.

A very simple model for a symmetry-breaking potential is the Ginzburg-Landau:

$$V \propto (\Phi^2 - \sigma^2)^2, \quad (2.176)$$

which gives a seeming “mass term” $-2\phi^2\sigma^2$, which has the wrong sign: it is “tachyonic”!

The configuration at $\Phi = 0$ is unstable. The one at $|\Phi| = \sigma$ is not symmetric under $\Phi \rightarrow -\Phi$.

People realized that QFT is a subcase of a condensed matter approach in which we have a thermal bath, an *environment*. This is *finite temperature QFT*.

We consider then an *effective potential* for the temperature: $V_T(\Phi) = V(\Phi) +$ functions of T . This might be $V(\Phi) + \alpha\phi^2 T^2 + \gamma T^4$, with positive α . The quadratic term then gives us a *positive* mass term: at temperatures larger than some critical temperature we get stability at $\Phi = 0$, but what happens if we lower the temperature?

Then, there is symmetry breaking.

Let us see how our Friedmann equations account for this situation.

The temperature of radiation is $\rho_r = \frac{\pi^2}{30} g_*(r) T^4$. If we start with a universe which is radiation dominate, then it ends up to be De Sitter.

This is a consequence of the *No Hair Cosmic Theorem*.

There is a potential barrier between the metastable “ $\Phi = 0$ ” state, and the symmetry breaking other ones. (even though it does not show in the fourth degree potential model).

This can happen through quantum tunneling, but there is a delay: a *first order phase transition with supercooling* (by “super” what is meant is just that the temperature goes below T_C even though we still are in the center symmetric state).

We get bubbles of symmetry broken by fluctuation, expanding through the universe but never meeting because of the expansion.

This is the “old inflation model”.

A new inflation model involves “slow rolling”.

The equation $\ddot{\phi} + 3H\dot{\phi} = -V'$ looks like a regular equation of motion: after a time $1/H$ the “friction” velocity-dependent term dominates.

Then we get a slow-roll regime: $H^2 = \frac{8\pi G}{3} V$ and $\dot{\phi} \approx -V'/3H$.

We exploit the flatness of the potential. There are quantum fluctuations during inflation.

The solution is *chaotic inflation*, by Linde 1984: Since $\Delta E \Delta t \approx \hbar$, we do not know at which state we are actually. As time passes, the energy uncertainty decreases. The initial condition for the distribution of the universe is then determined by the uncertainty principle.

An alternative is *eternal* chaotic inflation. If a fluctuation increases the potential universe, then H^2 increases, then the region feels a larger volume. The case where the field goes towards the minimum is unlikely. Why did it happen? This can only be answered with the anthropic principle.

Thu Nov 14 2019

During inflation, the comoving Hubble radius ($r_H = 1/\dot{a}$) decreases.

t_i is the beginning time of inflation, t_f is its end, t_Λ is the time when the cosmological constant became dominant, and then we get to now: t_0 .

What happened before inflation?

The cosmic no-hair conjecture is what allows inflation to delete inhomogeneties. So, there might have been perturbations before inflation: we cannot know. Up to which scale? Perturbations on scales larger than the cosmological horizon are not perceivable as perturbations: we only perceive our local mean value.

Below the largest inflation scale, the perturbations are erased by inflation: we see perturbations on these scales which are produced during inflation.

The energy density of radiation scales as $\rho_r \propto a^{-4}$, while the one of matter instead it scales as $\rho_m \propto a^{-3} \propto e^{-3Ht}$ since $a \propto e^{Ht}$.

The maximum temperature of radiation after inflation is given by the one which corresponds to the maximum latent energy released by our scalar field due to its coupling to the rest of the universe, which acts as a sort of viscous force.

So we get that this latent energy ΔV is of the order of T_{rad}^4 .

What is the typical temperature needed to produce baryon symmetry?

What is baryon symmetry?

These questions are hard to answer without a grand unification theory.

How much antimatter is there in the universe? A long time ago, it was thought that there might have been regions in the universe which were filled with antimatter by looking for a γ -ray background, but this was not found. We did not find them.

The baryon number is the difference between the number of baryons and antibaryons: $(n_b - n_a)/(n_b + n_a)$. This is actually computed with quark numbers, so it can be computed even when protons and neutrons have not yet formed.

It seems like the only antimatter known is the one which was formed by us, or cosmic rays.

The value $\Omega_{0b} \approx 0.04$ is defined as

$$\rho_{0b} = \Omega_{0b} \rho_c, \quad (2.177)$$

where $\rho_c = 3H_0^2/(8\pi G)$.

When we define the number of protons in the universe we also fix the number of neutrons, since the universe is globally neutral. So we can estimate:

$$n_{0b} = \frac{\rho_{0b}}{m_p} \approx 1.12 \times 10^{-5} \text{ cm}^{-3} \times h^2 \Omega_{0b}, \quad (2.178)$$

where $h \sim 0.7$, and similarly we can compute

$$n_{0\gamma} \approx 420 \text{ cm}^{-3}, \quad (2.179)$$

so we can look at the baryon to photon number ratio:

$$\eta_0 = \frac{n_{0b}}{n_{0\gamma}} \approx 3 \times 10^{-8} \Omega_{0b} h^2, \quad (2.180)$$

why does this number have this value?

The denominator in $(n_b - n_a)/(n_b + n_a)$ is approximately $2n_\gamma$ then, and both the difference in the numerator and the numerator scale like a^{-3} , so this value is a constant. We do not see antimatter, therefore $n_a = 0$. So, we get

$$\frac{n_b - n_a}{n_b + n_a} \approx \frac{n_b}{2n_\gamma} = \frac{1}{2} \eta_0. \quad (2.181)$$

This must then have been the case also in the matter dominated epoch, during which there was a very slight imbalance in baryons vs antibaryons.

In order to generate a baryon-antibaryon asymmetry we need

1. ?
2. C and CP violation (while CPT symmetry must hold for any well-behaved QFT);
3. out-of-equilibrium processes.

These were proposed by a famous Soviet scientist.

We define the interaction rate Γ : it is the number of interactions per unit time.

The Hubble rate $H = \dot{a}/a$ is also an inverse time: baryons and antibaryons are practically speaking *decoupled* if $\Gamma \leq H$: this is equivalent to saying that the time of interaction is larger than the age of the universe.

When are particles actually coupled or decoupled? Γ can be calculated as $\Gamma = n \langle \sigma v \rangle$, where n is the number density, v is the velocity of the particles, and σ is the cross section of the interaction.

We need to distinguish the types of interactions we are dealing with. In general interactions are carried by gauge bosons. Either they are massless (like the photon) or they are massive (like the W^\pm and Z bosons, as long we are below the scale of electroweak symmetry breaking $\sim 10^2$ GeV).

At larger energies than those, the weak interaction also becomes long-range.

In the massless case, the cross-section is $\sigma \sim \alpha^2/T^2$, where $g = \sqrt{4\pi\alpha}$.

In the massive case, for temperatures $T \leq m_x$, the cross section is of the order $\sigma \sim G_x^2 T^2$.

There is an inversion in the T dependence between long and short range interactions.

Typically $G_x = \alpha/m_x^2$. Then, $\sigma \sim \alpha^2 T^2/m_x^4$.

This difference is because in general the formula is like

$$\sigma \sim \frac{T^2}{E^4}, \quad (2.182)$$

and we either have $E \sim m$ in the low-speed case, or $E \sim T$ in the high-speed case.

So, we know that $\Gamma = n \langle \sigma v \rangle$. In the massless case this is something like $\Gamma \sim T^3 \sigma \sim T^3 \sim \alpha^2 T$ since $T \sim 1/a$.

Then, we get

$$H = \sqrt{\frac{8\pi G}{3}} g_*^{1/2} \left(\frac{\pi^2}{30} \right)^{1/2} T^2 \sim \frac{T^2}{m_{\text{pl}}}, \quad (2.183)$$

since $G \sim 1/m_{\text{pl}}^2$. Then,

$$\frac{\Gamma}{H} \sim \frac{\alpha^2 T m_{\text{pl}}}{T^2} \sim \alpha^2 m_{\text{pl}} \frac{1}{T}, \quad (2.184)$$

so we have decoupling when $T > \alpha^2 m_{\text{pl}}$. Essentially, at temperatures larger than $T = 10^{16}$ GeV this massless photon is decoupled.

“Above the Planck epoch, even gravitational interactions are decoupled”.

What about the massive interactions? We have $\Gamma \sim T^2 G_x^2 T^2$, to compare with $H \sim T^2/m_{\text{pl}}$: we get

$$\frac{\Gamma}{H} \sim \frac{T^3 G_x^2}{T^2/m_{\text{pl}}} \sim G_x^2 m_{\text{pl}} T^3 \leq 1, \quad (2.185)$$

equivalently, $T < m_{\text{pl}}^{-1/3} G_x^{-2/3}$.

Suppose that we are considering the gravitational interaction: in that case, we get $T < m_{\text{pl}}$ since G_x is related to G_N

What is the relation?

For the weak interaction, we have

$$T < \left(\frac{m_x}{100 \text{ GeV}} \right)^{4/3} \text{ MeV}, \quad (2.186)$$

which is why below 1 MeV neutrinos are decoupled.

Let us now consider the consequences of these decoupling conditions. First of all we look at the recombination of hydrogen. At very high temperatures, there are free electrons and free protons. Protons first appeared in the universe as non-relativistic, at $T \sim 1 \text{ MeV}$ while $m_p \sim 1 \text{ GeV}$.

At a certain point, it becomes possible to create neutral H atoms from these free particles. We will use a special case of the Boltzmann formula, which governs this and many other phenomena: the Saha equation.

The reaction is $e + p \leftrightarrow H + \gamma$. We want to look at a density in phase space. We'd need all the scattering matrices, and all the phase space densities of the particles. The Saha equation is basically an ansatz at thermal and chemical equilibrium: $\mu_e + \mu_p = \mu_H + \mu_\gamma$. The chemical potential μ entered in the exponent of the FD and BE expressions.

We know that $\mu_\gamma = 0$. At thermal equilibrium the number density of the electrons is:

$$n_e = g_e \left(\frac{m_e T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_e - m_e}{T} \right), \quad (2.187)$$

and an exactly analogous formula holds for n_p and n_H : for protons we have

$$n_p = g_p \left(\frac{m_p T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_p - m_p}{T} \right). \quad (2.188)$$

Degeneracy is not an issue, since we are talking about cosmology.

The number density of photons is given by

$$n_\gamma = \frac{2\zeta(3)T^3}{\pi^2}, \quad (2.189)$$

The binding energy of the hydrogen is $B = m_p + m_e - m_H = 13.6 \text{ eV}$. Instead of m_H we write $+m_p + m_e - B$. So the number density of hydrogen atoms is given by:

$$n_H = g_H \left(\frac{m_H T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_p - \mu_e - m_p - m_e + B}{T} \right), \quad (2.190)$$

and we can simplify things since $m_p, m_H \gg m_e, B$, and substitute in the number densities for electrons and protons. We get approximately

$$n_H = \frac{g_H}{g_e g_p} n_e n_p \left(\frac{m_e T}{2\pi} \right)^{-3/2} \exp(B/T), \quad (2.191)$$

but the universe is locally and globally neutral: $n_e = n_p$, and $n_b = n_p + n_H$.

We will see that most of the hydrogen we produce will not be when the temperature is of the order of the binding energy, but much later.

Fri Nov 15 2019

We start again from where we left off, with hydrogen recombination.

We want to estimate the moment at which hydrogen first formed, which marks the point at which electrons and photons interact efficiently: before, they interacted with Compton scattering which is very efficient; after they interact with hydrogen atoms in a way that is very inefficient.

After this, then, we say that photons and matter are *decoupled*.

The scattering cross section (for Compton?) goes like the inverse square of the mass.

This means that the universe is not only *globally*, but also *locally* neutral.

This decoupling is what allows for star formation. Also, this is when the CMB starts. It is made of microwaves now, but it was higher earlier.

The phase space distribution of photons is scale-invariant since they have zero mass: so we can say that the photons' distribution *looks* thermal, but it is actually not technically since there are no interactions anymore.

However the photons travel freely and are perceived as thermal, and they give an almost perfect blackbody! The errorbars in a plot for it must be magnified by 10^4 in order to be seen.

We have $\Gamma_\gamma = n_e \sigma_T$, where σ_T is the Thompson cross section. Neutrality implies $n_e = n_p$. In principle we should account for Helium: a couple minutes after the BB He-4 nuclei started to form, but they made up only something like 25% of the mass, which means 6% of the number density: so we say that the number density of baryons is

$$n_b = n_p + n_H. \quad (2.192)$$

Our ansatz for the Boltzmann equation is $\mu_e + \mu_p = \mu_H$ since photons have no chemical potential.

i denotes a generic one in e , p and H . Then

$$n_i = g_i \left(\frac{m_i T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_i - m_i}{T} \right), \quad (2.193)$$

we need to account for the chemical potential since it is the driver of this process. How much is n_e/n_b ? the same as n_p/n_b . We call this quantity X_e , the ionization number. We expect $X_e = 1$ in the early universe, and at the end of the process it will diminish: naively we'd expect it to get to 0, but actually there remains some residual ionization, some free protons and electrons. A proper calculation would account for the non-equilibrium contributions. However, we estimate the process as being in equilibrium: this will underestimate the number of electrons.

Today, most of the hydrogen is ionized (there is a ***-Peterson effect which shows this): this means that matter and radiation interact again.

This is the second important time in the history of the universe.

Can we see the early stars? Not really, we see galaxies only up to $z \sim 10$, these stars would be at something like $z \sim 30$. . . There might be more to this.

Let us come back to the calculation:

$$n_H = g_H \left(\frac{m_H T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_H - m_H}{T} \right) \quad (2.194a)$$

$$= g_H \left(\frac{m_H T}{2\pi} \right)^{3/2} \exp \left(\frac{\mu_e + \mu_p - m_e - m_p + B}{T} \right) \quad (2.194b)$$

$$= \frac{g_H}{g_e g_p} \left(\frac{m_H T}{2\pi} \right)^{3/2} \left(\frac{m_e T}{2\pi} \right)^{-3/2} \left(\frac{m_p T}{2\pi} \right)^{-3/2} n_e n_p \exp \left(\frac{B}{T} \right) \quad (2.194c)$$

$$\frac{n_H}{n_e n_p} = \left(\frac{m_e T}{2\pi} \right)^{-3/2} \exp \left(\frac{B}{T} \right) \quad (2.194d)$$

$$\frac{n_b - n_p}{n_p^2} = \left(\frac{m_H T}{2\pi} \right)^{3/2} \exp \left(\frac{B}{T} \right), \quad (2.194e)$$

which we can manipulate, using the following identity:

$$\frac{n_b - n_p}{n_p^2} = \frac{n_b (1 - n_p/n_b)}{n_b^2 X_e^2} = \frac{1}{n_b} \frac{1 - X_e}{X_e}, \quad (2.195)$$

where we use: $n_e = n_p$, and the definition of $X_e = n_p/n_b$. Then, we bring the n_b to the other side of the equation: we get

$$\frac{1 - X_e}{X_e^2} = \underbrace{\frac{n_b}{n_p}}_{\eta_p} \left(\frac{m_e T}{2\pi} \right)^{-3/2} \exp \left(\frac{B}{T} \right) n_\gamma \quad (2.196a)$$

$$= \frac{4\sqrt{2}\zeta(3)T^3}{\pi^2} \left(\frac{m_e T}{2\pi} \right)^{-3/2} \exp \left(\frac{B}{T} \right) \eta_0 \quad (2.196b)$$

$$= \frac{4\sqrt{2}\zeta(3)}{\sqrt{\pi}} \eta_0 \left(\frac{T}{m_e} \right)^{3/2} \exp \left(\frac{B}{T} \right), \quad (2.196c)$$

which means $T \sim 0.3 \text{ eV}$, much lower than the ionization energy of Hydrogen.

This does depend on the value we assign to Ω_0 and h .

Approximately, it occurred somewhere around $z \sim 1100$ (we conventionally say that recombination happened when $X_e = 0.1$).

After recombination, we have the *last scattering*: the moment at which the CMB was formed.

One Nobel prize this year was awarded to Jim Peebles, a friend of Sabino's: together with his PhD supervisor Dicke, he was the first to calculate this stuff.

Peebles in 1964 (?) did this calculation both in GR and in Brahms-Dicke theory, a modified gravity theory.

Let us describe the early universe, before the first nucleosynthesis. Important papers in this topic are by G. Gamow, and by Alpher, Bethe and Gamow.

Our hypotheses are:

1. the universe passed through a very high temperature phase, with $T > 10^{12}$ K;
2. the universe is described by GR and SM;
3. the chemical potentials for the neutrinos μ_ν have certain upper bounds;
4. there is no matter-antimatter separation (as in, “bubbles”);
5. there are no strong magnetic fields;
6. the number of exotic particles has a certain upper bound.

There are magnetic fields in the universe, but they are not homogeneous and relatively weak. Exotic particles are predicted by certain unification theories, they are generically defined as ones which we have not observed yet.

We have to explain the fact that we observe an excess of He-4 in the early universe: we define the yield

$$y \equiv \frac{m_{\text{He-4}}}{m_b} > 0.25. \quad (2.197)$$

In terms of particle number, the ratio is more like 0.06. We do not produce Carbon or anything higher than it: the process which forms it is inefficient at high temperature, low density like the early universe. Higher Z elements are only produced in stars.

Helium-4 is produced but also destroyed by stars.

The main channels in the early universe are:

1. $n + p \leftrightarrow d + \gamma$ (d denotes deuterium);
2. $d + d \leftrightarrow {}^3\text{He} + n$;
3. ${}^3\text{He} + d \leftrightarrow {}^4\text{He} + p$.

There is no weak interaction here, unlike what happens in stars. In stars, there are no free neutrons so this process is not possible.

We cannot treat this properly, we only give a story. The slowest process of the three is the first, since it is heavily affected by photons: they destroy deuterium.

The binding energy of deuterium is around 2.2 MeV. After we have produced deuterium, Helium-4 is readily produced.

First of all, we need the neutron to proton ratio. We are working at energies of around 1 MeV, so protons and neutrons are not relativistic anymore. This process takes place around three minutes after the beginning. For $i = n, p$:

$$n_i = g_i \left(\frac{m_i T}{2\pi} \right)^{3/2} \exp\left(\frac{\mu_i m_i}{T} \right), \quad (2.198)$$

so their number ratio is around:

$$\frac{n}{p} \sim \exp\left(\frac{m_p - m_n}{T}\right), \quad (2.199)$$

where $m_p - m_n \approx 1.3 \text{ MeV} \approx 1.5 \times 10^{10} \text{ K}$.

We have the processes

1. $n + \gamma_e \leftrightarrow p + e^-$;
2. $n + e^+ \leftrightarrow p + \bar{\nu}_e$;
3. $n \rightarrow p + e^- + \bar{\nu}_e$.

We can replace the temperature in the exponential by $T \rightarrow T_{d_v}$, the decoupling temperature of the neutrinos, since that is the moment around which this happens.

So, we get around $\exp(-1.5)$. We define:

$$X_n(t) \equiv \frac{n}{n+p} \sim 0.17. \quad (2.200)$$

Later it will change because of β decay, going like:

$$X_n(t) = X_n(t_{d_v}) \exp\left(-\frac{t - t_{d_v}}{\tau_n}\right), \quad (2.201)$$

where $\tau_n = \log 2 \tau_{1/2}$, and this half-life is $\tau_{1/2} \approx 10.5 \pm 0.2$.

The binding energy of deuterium is around 2.2 MeV. Then, we proceed exactly like we did with hydrogen, and finally get

$$X_d = \frac{3}{4} n_b X_n X_p \left(\frac{m_d}{m_n m_p}\right)^{3/2} \left(\frac{T}{2\pi}\right)^{-3/2} \exp\left(\frac{B}{T}\right) \quad (2.202a)$$

$$= \frac{3}{4} \eta_0 X_n X_p \left(\frac{m_d}{m_n m_p}\right)^{3/2} \frac{2\zeta(3)}{\pi^2} (2\pi)^{3/2} T^{3/2} \exp\left(\frac{B}{T}\right) \quad (2.202b)$$

$$\approx \frac{3}{4} \eta_0 X_n (1 - X_n) \left(\frac{m_d}{m_n m_p}\right)^{3/2} \frac{2}{\pi^2} (2\pi)^{3/2} \zeta(3) T^{3/2} \exp\left(\frac{B_d}{T}\right), \quad (2.202c)$$

Check calculation.

which describes the *deuterium bottleneck*, which is what impedes this process until photons are very diluted. As soon as photons are diluted enough, they stop bottlenecking.

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We exponentiate the equation from before: we get

$$X_d = X_n X_p \exp\left(-29.33 + \frac{25.82}{T_\rho} - \frac{3}{2} \log T_\rho + \log(\Omega_0 h)\right), \quad (2.203)$$

where $T_\rho =$

What is going on? What is T_ρ ?

We want to understand why there is so much He-4 in the universe, since it is destroyed in stars!

This model fits observation as long as $0.011h^{-2} \leq \Omega_0 \leq 0.25h^{-2}$. Most people agree that we are around the upper bound.

This is indirect evidence for dark energy: why?

The lifetime of the neutron, $\tau_{1/2}$ is something that is also relevant, since it affects the baryon ratios.

The Gamow factor Γ is proportional to the Fermi coupling constant G_F^2 , which is connected to $\tau_{1/2}$.

Let us suppose we increase the lifetime of neutrons, $\tau_{1/2}$. This changes the moment at which we reach equation $\Gamma \sim H$.

Increasing the lifetime of neutrons decreases the amount of He-4 in the universe: less is produced.

We know that

$$H^2 = \frac{8\pi G}{3}\rho_r, \quad (2.204)$$

where $G = 1/m_p^2$ and $\rho = \frac{\pi^2}{30}g_*(T)T^4$.

If we fix the temperature, and change g_* (by adding degrees of freedom), then we get more He-4.

This gives us observational constraints on the additions of exotic particles to our theory, since that would change g_* . This bounds the number of neutrino families by 3.0 something, so there cannot be more than 3 families of light neutrinos.

If gravitons are thermal, then they also contribute to radiation.

There is also another parameter in the Friedmann equation; it is m_P : modified gravity theories often predict variations of the gravitational constant with time.

Dark matter has no relevant electromagnetic interactions: it only interacts gravitationally, and is able to cluster; dark energy, instead, is uniformly distributed.

We divide it into Hot and Cold dark matter: HDM and CDM. There is also something called *Warm* dark matter, which has intermediate properties.

In HDM, particles have very high thermal motion. They move fast, and tend to destroy gravitational potential wells in which they might settle by moving out of them, and thus decreasing the quantity of matter there.

They do this on scales comparable to the maximum distance travelled by them: this is calculated as vt , where v is their average thermal velocity, (and t is the age of the universe?).

The structures formed by these are of scales similar to or larger than $10^{15}M_\odot$, but we observe smaller structures also! They were formed later, by fragmentation: this is the top-down approach.

We also have a bottom-up approach, which is compatible with CDM.

Neutrinos were thought to be Dark Matter, and would have been hot.

The top-down approach, however, is falsified by the observation of high-redshift quasars combined with the scale of the anisotropies of the CMB: in order to account for high-redshift small-scale structures (we have seen stars at $z \sim 20$!) we would have to increase the amplitude of the anisotropies to a scale which is not compatible to the anisotropies we see in the CMB.

Now, neutrinos are not useful for cosmology.

We have $\Gamma \sim T^5$, and $\Gamma = H$: $T^5/\tau_{1/2} \sim T^2$ implies that $\tau_{1/2} \sim T^3$.

The decoupling temperature for HDM is larger than the temperature at which they become nonrelativistic, which is of the order of the mass.

For CDM, instead, the decoupling temperature is *smaller* than the temperature at which they become relativistic.

Now, we discuss the Boltzmann equation: there is an operator acting in phase space, the Liouville operator, which is equal to the collision operator.

$$\mathbb{L}[f] = \mathbb{C}[f], \quad (2.205)$$

where all the scattering operators live on the RHS.

We start with a Newtonian description: the phase space has position, momentum and time as coordinates, and on it we define a density function $f(\vec{q}, \vec{p}, t)$.

The Einstein equations are blind to the momentum distribution, since $T_{\mu\nu}$ does not depend on the momentum. We can say that all of the momentum has been marginalized.

We define the operator

$$\hat{\mathbb{L}} = \frac{D}{Dt} = \frac{\partial}{\partial t} + \frac{d\vec{x}}{dt} \cdot \nabla_x + \frac{d\vec{p}}{dt} \cdot \nabla_p \quad (2.206a)$$

$$= \frac{\partial}{\partial t} + \vec{v} \cdot \nabla_x + \frac{\vec{F}}{m} \cdot \nabla_p. \quad (2.206b)$$

In GR, we geometrize the gravitational force: it is included in the inertial motion of the particles.

The relativistic version of this is

$$\hat{\mathbb{L}} = p^\alpha \partial_\alpha - \Gamma_{\beta\gamma}^\alpha p^\beta p^\gamma \frac{\partial}{\partial p^\alpha}. \quad (2.207)$$

This is because $p^\alpha = dx^\alpha/d\lambda$ satisfies the geodesic equation:

$$\frac{dp^\alpha}{d\lambda} + \Gamma_{\beta\gamma}^\alpha p^\beta p^\gamma = 0, \quad (2.208)$$

or $\frac{Dp^\alpha}{Dt} = 0$. We can then write down the Christoffel bit of the geodesic equation instead of the derivative $dp^\alpha/d\lambda$. This is because, *on shell*:

$$g^{\alpha\beta} p_\alpha p_\beta = m^2. \quad (2.209)$$

We want to describe the *abundance* of dark matter, in phase space: the number of DM particle per EM-interacting particle.

We then make the Liouville operator explicit, with the Christoffel symbols of the RW metric. Acting on f , we get

$$\hat{\mathbb{L}} = p^0 \partial_t f - \frac{\dot{a}}{a} |\vec{p}|^2 \frac{\partial f}{\partial E}, \quad (2.210)$$

which we will not prove.

The derivative with respect to the momentum must be ∂_{p^0} ,

because of isotropy?

the Christoffel symbols of RW are

$$\Gamma_{\beta\gamma}^0 = \frac{\dot{a}}{a} \delta_{\beta\gamma}. \quad (2.211)$$

Recall the definition of the number density

$$n(t) = \frac{g}{(2\pi)^3} \int d^3 \vec{p} f(|\vec{p}|, t), \quad (2.212)$$

so in the end the equation becomes:

$$\hat{\mathbb{L}} = \dot{n} + 3 \frac{\dot{a}}{a} n. \quad (2.213)$$

Conventionally we divide by p^0 : we get for the Boltzmann equation:

$$\frac{\partial f}{\partial t} - \frac{\dot{a}}{a} \frac{p^2}{E} \frac{\partial f}{\partial E} = \frac{1}{E} \hat{\mathbb{C}}[f]. \quad (2.214)$$

Now we integrate in $d^3 \vec{p}$. We get

$$\frac{\partial n}{\partial t} - \frac{\dot{a}}{a} \int d^3 \vec{p} \frac{p^2}{E} \frac{df}{dE} = \frac{g}{(2\pi)^3} \int d^3 p \frac{1}{E} \mathbb{C}[f]. \quad (2.215)$$

More properly, we are doing $\mathbb{L} \langle f \rangle$.

We manipulate:

$$\int d^3 p \frac{p^2}{E} \frac{\partial f}{\partial E} = 2 \int d^3 p p^2 \frac{\partial f}{\partial E^2} \quad (2.216a)$$

$$= 2 \int d^3 p p^2 \frac{\partial f}{\partial p^2} = \int d^3 p p \frac{\partial f}{\partial p} \quad (2.216b)$$

$$= 2\pi \int_0^\infty dp p^3 \frac{\partial f}{\partial p} = -3 \int d^3 p p f. \quad (2.216c)$$

Here, we integrated by parts in the second to last step, and set to zero the boudary term $4\pi p^3 f$, calculated from 0 to infinity, since at 0 we have $p = 0$, and at (momentum) infinity we have $f = 0$.

So, we can see that the LHS is equal to $\dot{n} + 3\dot{a}n/a$: so we find

$$\dot{n} + 3 \frac{\dot{a}}{a} n = \frac{g}{(2\pi)^3} \int d^3 p \frac{1}{E} \hat{\mathbb{C}}[f]. \quad (2.217)$$

This is the cosmological version of the Boltzmann equation.
We model the RHS as something like

$$\hat{\mathcal{C}}[f] = \Psi - \langle \sigma v \rangle n^2, \quad (2.218)$$

where we have $\Gamma_A = \langle \sigma v \rangle n$ times n , where Γ is the rate of annihilation.

At equilibrium the LHS is equal to zero: so we can write $\Psi = \langle \sigma v \rangle n_{\text{eq}}^2$, and the RHS becomes $\langle \sigma v \rangle (n_{\text{eq}}^2 - n^2)$.

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Yesterday we arrived at the following equation:

$$\dot{n} + 3\frac{\dot{a}}{a}n = \Psi - \langle \sigma_A v \rangle n^2, \quad (2.219)$$

where the term Ψ is the source of new particles, while the next term accounts for the annihilation of particles.

Under decoupling, the equation reduces to

$$\dot{n} + 3\frac{\dot{a}}{a}n = 0 \implies \frac{d}{dt}(na^3) = 0. \quad (2.220)$$

Under equilibrium, $\Gamma > H$.

Therefore $\Psi = \langle \sigma_A v \rangle n_{\text{eq}}^2$. Recall that the Γ of annihilation is equal to $\Gamma = \langle \sigma_A v \rangle n$.

If the timescale of the collisions is longer than the age of the universe, $\Gamma < H$, then once again we have $a^3 n = \text{const}$.

We can see thermal *relics*, thermal distributions of objects which are not coupled anymore.

We define

$$n_C = n \left(\frac{a}{a_0} \right)^3, \quad (2.221)$$

so we can simplify the expression:

$$\dot{n} + 3\frac{\dot{a}}{a}n = \dot{n} \left(\frac{a_0}{a} \right)^3. \quad (2.222)$$

We want to factor out a parameter H . We get:

$$\dot{n}_C = - \left(\frac{a}{a_0} \right)^3 \langle \sigma_A v \rangle \left(\frac{a_0}{a} \right)^6 (n_C^2 - n_{\text{eq}}^2), \quad (2.223)$$

which can be written as

$$\frac{a}{n_{C, \text{eq}}} \frac{dn_C}{da} = - \frac{\langle \sigma_A v \rangle n_{\text{eq}}}{\dot{a}/a} \left(\left(\frac{n}{n_{\text{eq}}} \right)^2 - 1 \right), \quad (2.224)$$

so, if we define τ_{coll} as $1/(\langle\sigma_A v\rangle n_{\text{eq}})$ and $\tau_{\text{exp}} = 1/H$, we get

$$n_{\text{C, eq}} a \frac{dn_{\text{C}}}{da} = -\frac{\tau_{\text{exp}}}{\tau_{\text{coll}}} \left(\left(\frac{n}{n_{\text{eq}}} \right)^2 - 1 \right), \quad (2.225)$$

so if $\Gamma \gg H$, then $\tau_{\text{exp}}/\tau_{\text{coll}} \gg 1$, therefore $n = n_{\text{eq}}$, which also implies $n_{\text{C}} = n_{\text{C, eq}}$. This is the equilibrium case.

On the other hand, if $\Gamma \ll H$, then $\tau_{\text{exp}}/\tau_{\text{coll}} \ll 1$ we have decoupling, therefore $n_{\text{C}} = \text{const.}$

We know that at temperatures below 1.5 MeV (after decoupling), the following holds:

$$T_\nu = \left(\frac{4}{11} \right)^{1/3} T_\gamma, \quad (2.226)$$

and we want to do the same for dark matter.

Neutrinos are nonrelativistic today, but they became so at a relatively low redshift, a short time ago. We can parametrize the number density of neutrinos by the temperature.

The formula for the temperature of neutrinos is a special case of the following formula:

$$T_{0\nu} = \left(\frac{g_{*0}}{g_{*d}} \right)^{1/3} T_{0\gamma}, \quad (2.227)$$

where 0 means *now*, while *d* means *decoupling*. This can be applied to any species.

Let us compute this for a generic species x : we find

$$n_{0x} = B g_* \frac{\zeta(3)}{\pi^2} T_{0x}^3, \quad (2.228)$$

where the factor B accounts for the statistics: it is 1 for bosons, 3/4 for fermions. It is important to note that the temperature is a parameter, but these particles are *not thermal* anymore!

So, we get, using equation (2.189):

$$n_{0x} = \frac{B}{2} n_{0\gamma} g_x \frac{g_{*0}}{g_{*dx}}. \quad (2.229)$$

The energy density in general is given by $\rho_{0x} = m_x n_{0x}$. We get:

$$\rho_{0x} = \frac{B}{2} m_x n_{0\gamma} g_x \frac{g_{*0}}{g_{*gx}}, \quad (2.230)$$

therefore

$$\Omega_{0x} h^2 = \frac{m_x n_{0x}}{\rho_{0x}} h^2 = 2B g_x \frac{g_{*0}}{g_{*dx}} \frac{m_x}{10^2 \text{ eV}}. \quad (2.231)$$

CDM is made of particles which were already nonrelativistic when they decoupled.

Then, in the formula

$$n_x(T_{dx}) = g_x \left(\frac{m_x T_{dx}}{2\pi} \right)^{3/2} \exp \left(-\frac{m_x}{T_{dx}} \right), \quad (2.232)$$

we can assume that $T_{dx} \ll m_x$. So,

$$n_{0x} = n_x(T_{dx}) \left(\frac{a(T_{dx})}{a_0} \right)^3 = n(T_{dx}) \frac{g_{*0}}{g_{*x}} \left(\frac{T_{0\gamma}}{T_{dx}} \right)^3, \quad (2.233)$$

but it is difficult to compute T_{dx} , which is the solution to the equation $\Gamma(T_{dx}) = H(T_{dx})$.

We know that

$$H^2(T_{dx}) = \frac{8\pi G}{3} g_{*x} \frac{\pi^2}{30} T_{dx}^4, \quad (2.234)$$

or, in terms of the quantity $\tau_{\text{exp}} = 1/H$:

$$\tau_{\text{exp}} = 0.3 g_*^{-1/2} \frac{m_{\text{pl}}}{T_{dx}^2}. \quad (2.235)$$

On the other hand, $\Gamma = n \langle \sigma_A v \rangle$, and $\tau_{\text{coll}}(T_{dx})$, and in terms of τ_{coll} we get

$$\tau_{\text{coll}}(T_{dx}) = \left(n(T_{dx}) \sigma_0 \left(\frac{T_{dx}}{m_x} \right)^N \right), \quad (2.236)$$

where $N = 0, 1$: it is a fact from particle physics that the average has this kind of dependence:

$$\langle \sigma_A v \rangle = \sigma_0 \left(\frac{T}{m_x} \right)^N. \quad (2.237)$$

So, equaling the two τ we get:

$$\left(n(T_{dx}) \sigma_0 \left(\frac{T_{dx}}{m_x} \right)^N \right) = 0.3 g_*^{-1/2} \frac{m_{\text{pl}}}{T_{dx}^2}, \quad (2.238)$$

which can be solved iteratively. We solve it in terms of the parameter $x_{dx} = m_x/T_{dx}$, which must be much larger than one: this allows us to select the physical solution to the equation among the nonphysical ones.

The solution is found to be something like:

$$x_{dx} = \log \left(0.038 \frac{g_x}{g_{*xd}^{1/2}} m_{\text{pl}} m_x \sigma_0 \right) - \left(N - \frac{1}{2} \right) \log \log (\dots). \quad (2.239)$$

Chapter 3

Stellar Astrophysics

Dark matter is collisionless, however it is not *really*: the censorship theorem says that it cannot actually collapse into a naked singularity.

A star is a sphere of matter, usually baryonic matter, characterized by a density $\rho(r)$: the mass enclosed in a radius r is

$$m(r) = \int_0^r 4\pi\tilde{r}^2 \rho(\tilde{r}) d\tilde{r} . \quad (3.1)$$

The gravitational acceleration can be calculated from Gauss' theorem:

$$g(r) = \frac{Gm(r)}{r^2} . \quad (3.2)$$

Let us consider a small volume, with its enclosed mass $\Delta M = \rho(r)\Delta A\Delta r$. It can contrast the inward force due to gravity if there is a differential pressure.

Let us denote P as the pressure at the inner surface, and $P + \Delta P$ the pressure at the outer surface. Then,

$$(P + \Delta P)\Delta A - P\Delta A = \left(P(r) + \frac{dP}{dr}\Delta r \right)\Delta A - P(r)\Delta A = \frac{dP}{dr}\Delta A\Delta r . \quad (3.3)$$

The minus sign in what follows is since the force is inward:

$$-\Delta M\ddot{r} = \Delta M g(r) + \frac{dP}{dr}\Delta r\Delta A . \quad (3.4)$$

The equation of motion is

$$-\ddot{r} = g(r) + \frac{1}{\rho(r)} \frac{dP}{dr} , \quad (3.5)$$

so we see that the pressure gradient must have a minus sign.

If the internal energy is used up to do internal (chemical, nuclear) work, then it cannot support the star anymore, and it then collapses.

Usually, stars start from the Jeans phenomenon: dark matter and baryons collapse under their own weight.

Let us start at the decoupling of matter and radiation, something like $z = 1100$.

The clouds of matter collapse freely up to the moment at which their internal pressure starts to slow them down. Then, we get the necessary conditions for the fusion of Hydrogen.

During freefall, we do not have a pressure gradient, so we have

$$-\ddot{r} = g(r). \quad (3.6)$$

It is not generally the case, but let us suppose that the collapse is *orderly*, the interior collapses before the outer layers. We have to account for energy conservation: the total energy when our test shell is at a radius r is conserved as r changes. If the energy (per unit mass) is zero at infinity we have

$$E_{\text{tot}}(r_0) - \frac{Gm_0}{r_0} = \frac{1}{2} \left(\frac{dr}{dt} \right)^2 - \frac{Gm_0}{r}, \quad (3.7)$$

therefore

$$\frac{1}{2} \left(\frac{dr}{dt} \right)^2 = Gm_0 (r^{-1} - r_0^{-1}). \quad (3.8)$$

What is the freefall time? we take dt/dr from the equation and integrate it from r_0 to 0

$$t_{\text{free fall}} = \int_{r_0}^0 dr \frac{dt}{dr} = - \int_{r_0}^0 dr \frac{1}{2GM} \left(\frac{1}{r} - \frac{1}{r_0} \right)^{-1/2}, \quad (3.9)$$

where we have a minus sign since $dr/dt < 0$. We get

$$t_{\text{free fall}} = \sqrt{\frac{r_0^3}{2Gm_0}} \int_0^1 dx \sqrt{\frac{x}{1-x}} = \frac{\pi}{2} \sqrt{\frac{r_0^3}{2Gm_0}}, \quad (3.10)$$

and if we define the average density $\bar{\rho} = m_0 / (4\pi r_0^3/3)$ we get the simpler expression

$$t_{\text{free fall}} = \sqrt{\frac{3\pi}{32G\bar{\rho}}}. \quad (3.11)$$

We might be tempted to ignore the expansion of the universe in these calculations: we know that

$$H^2 = \frac{8\pi G}{3} \bar{\rho}, \quad (3.12)$$

and in the purely Newtonian case we know that $a(t) \propto t^{2/3}$ and $H = 2/(3t)$.

Therefore

$$\frac{4}{3t^2} = \frac{8\pi G}{3} \bar{\rho} \implies t^2 = \frac{4}{9} \frac{3}{8\pi G} \bar{\rho}^{-1}, \quad (3.13)$$

so

$$t_{\text{exp}} \propto \bar{\rho}^{-1/2}. \quad (3.14)$$

what is the relation between the universe's density and the star's?

Is the idea: the matter in the universe does not disperse too fast, stars theoretically are allowed to form?

At equilibrium, $\ddot{r} = 0$: so

$$-\ddot{r} = 0 = g(r) + \frac{1}{\rho(r)} \frac{dP}{dr}, \quad (3.15)$$

so hydrostatic equilibrium implies that, locally,

$$\frac{dP}{dr} = -G \frac{m(r)\rho(r)}{r^2}. \quad (3.16)$$

We multiply both sides by $4\pi r^3$ and integrate in dr : we get

$$\int_0^R dr 4\pi r^3 \frac{dP}{dr} = -G \int_0^R \frac{m(r)\rho(r)4\pi r^2}{r} dr, \quad (3.17)$$

and we can change variables: $\rho(r)4\pi r^2 dr = dm$, so we can identify the LHS with the total gravitational energy:

$$E_{\text{grav}} = -G \int \frac{m(G) dm}{r}. \quad (3.18)$$

On the RHS, instead, we can integrate by parts:

$$\left[P(r)4\pi r^3 \right]_0^R - 3 \int_0^R dr 4\pi r^2 P(r), \quad (3.19)$$

where the boundary term is zero: at the origin $r = 0$, at the surface (by definition) $P = 0$.

So, inserting the volume $V(R) \equiv 4\pi R^3/3$, we have

$$E_{\text{grav}} = -3 \int_0^R \frac{dr 4\pi r^2 P(r)}{V(R)} V(R), \quad (3.20)$$

which gives us the virial theorem: we can interpret the integral as a weighted average, so we get

$$E_{\text{grav}} = -3 \langle P \rangle V \quad \text{or} \quad \langle P \rangle = -\frac{1}{3} \frac{E_{\text{grav}}}{V}. \quad (3.21)$$

This is Newtonian: above a certain limit, the relativistic corrections destabilize the system.

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We derived

$$\langle T \rangle = -\frac{1}{3} \frac{E_{\text{grav}}}{V}, \quad (3.22)$$

so now we proceed: we want a relation between kinetic and gravitational energy densities.

We use a statistical mechanics approach.

The rate of momentum transfer in the direction x is given by

$$\frac{N}{L^3} p_x v_x, \quad (3.23)$$

and this will be true for either direction by isotropicity: so

$$P = \frac{n}{3} \langle \vec{p} \cdot \vec{v} \rangle, \quad (3.24)$$

in full generality.

Let us consider the two cases of nonrelativistic and fully relativistic. In the nonrelativistic case we have

$$\epsilon_p = mc^2 + \frac{p^2}{2m}, \quad (3.25)$$

where $p = mv$. In the ultrarelativistic case we have

$$\epsilon_p = pc, \quad (3.26)$$

and the velocity is approximately the speed of light.

Then, for a gas of nonrelativistic particles, we have

$$P = \frac{1}{3} n m v^2 = \frac{2}{3} n \left\langle \frac{1}{2} m v^2 \right\rangle = \frac{2}{3} \times \text{translational KE density}, \quad (3.27)$$

while in the relativistic case we have:

$$P = \frac{1}{2} n \langle pc \rangle = \frac{1}{3} \times \text{translational KE density}. \quad (3.28)$$

We will show that, if a star is made of a gas of classical nonrelativistic particles it tends to be stable, if the particles are relativistic then it tends not to be stable.

The virial theorem tells us that

$$2E_K + E_{\text{grav}} = 0, \quad (3.29)$$

in the nonrelativistic approximation.

We define: $\Delta E_{\text{tot}} = -\Delta E_K = \frac{1}{2} \Delta E_{\text{grav}}$.

We know that

$$\langle P \rangle = \frac{1}{3} \frac{E_K}{V} = -\frac{1}{3} \frac{E_{\text{grav}}}{V} \quad (3.30)$$

by the virial theorem: so the total binding energy is equal to zero, since this gives us

$$E_{\text{grav}} + E_K = E_{\text{tot}} = 0. \quad (3.31)$$

Now we differentiate the law $d(PV^\gamma) = 0$, since it is a constant for an adiabatic transformation: it gives us, using logarithmic derivatives,

$$\gamma \frac{dV}{V} + \frac{dP}{P} = 0, \quad (3.32)$$

so

$$d(PV) = -(\gamma - 1)P dV, \quad (3.33)$$

and we know that for an adiabatic transformation

$$dE_{\text{in}} + P dV = 0, \quad (3.34)$$

which implies

$$dE_{\text{in}} = \frac{1}{\gamma - 1} d(PV), \quad (3.35)$$

and let us assume that γ is approximately constant in the transformation: this means

$$E_{\text{in}} = \frac{PV}{\gamma - 1}, \quad (3.36)$$

so

$$P = (\gamma - 1) \frac{E_{\text{in}}}{V}, \quad (3.37)$$

which justifies the relations we used in cosmology, $P = w\rho$ with $w = \gamma - 1$.

We can rewrite the equation from before as

$$3(\gamma - 1)E_{\text{in}} + E_{\text{gr}} = 0, \quad (3.38)$$

which, together with $E_{\text{tot}} = E_{\text{in}} + E_{\text{gr}}$, give us

$$E_{\text{tot}} = -(3\gamma - 4)E_{\text{in}}, \quad (3.39)$$

which means that $\gamma > 4/3$ characterizes a bound system, while $\gamma < 4/3$ characterizes a free system.

There are two dangers: one is the fight against the pressure forces, one is the fight against the quantum forces (the Pauli exclusion principle) which do not allow the compression to happen further.

Now we discuss Jeans instability:

$$E_{\text{grav}} = -f \frac{GM^2}{R}, \quad (3.40)$$

where f is a numerical factor of the order 1, depending on the mass distribution. If the distribution is uniform, it is $3/2$.

3/2?

The kinetic energy is

$$E_K = \frac{3}{2} N k_B T, \quad (3.41)$$

and we want to impose the condition

$$f \frac{GM^2}{R} > \frac{3}{2} N k_B T, \quad (3.42)$$

and the Jeans criterion is this boundary of the stability region:

$$f \frac{g M_J^2}{R} = \frac{3}{2} \frac{M_J}{\bar{m}} k_B T, \quad (3.43)$$

where $\bar{m} = M/N$. The J denotes the fact that we are considering the specific boundary mass on both sides. Simplifying the formula we find:

$$M_J = \frac{3}{2} \frac{k_B T}{G \bar{m}} R, \quad (3.44)$$

and we can reframe this in terms of the density, which is defined by

$$M_J = \frac{4\pi}{3} \rho_J R^3. \quad (3.45)$$

We cube and multiply on both sides:

$$\rho_J M_J^3 = \left(\frac{3 k_B T}{2 G \bar{m}} \right)^3 R^3 \rho_J, \quad (3.46)$$

so we get

$$\rho_J = \frac{1}{M_J^3} \left(\frac{3 k_B T}{G \bar{m}} \right)^3 \left(\frac{4\pi}{3} \rho_J R^3 \right) \frac{1}{4\pi}, \quad (3.47)$$

so

$$\rho_J = \frac{3}{4\pi M_J^2} \left(\frac{3 k_B T}{2 G \bar{m}} \right)^3, \quad (3.48)$$

and we will have stability if the density is larger than this. So, if we want a collapse, we must decrease the mass...

When the last scattering happens, the pions are decoupled from the photons. Dark matter behaves differently from conventional matter.

We have

$$\dot{\rho}_r = -3H(\rho_r + P_r), \quad (3.49)$$

and

$$\dot{\rho}_m = -3H(\rho_m + P_m), \quad (3.50)$$

and $P_r = \rho_r/3$, which scale like a^{-4} and also as T^4 , which means $T \sim 1/a$.

$$d(\rho_m c^2 a^3) + P_m da^3 = 0, \quad (3.51)$$

where we usually approximate $\rho_m c^2 = m_p n_b c^2$, but we can include more terms:

$$\rho_m c^2 = m_p n_b c^2 \left(1 + (\gamma - 1)^{-1} \frac{k_B T}{m_p c^2} \right), \quad (3.52)$$

while the pressure is given by $P = n_b k_B T$: so in the end we find

$$d \left(\left(m_p n_b c^2 + \frac{3}{2} m_p n_b \frac{k_B T}{m_p} \right) a^3 \right) = -n_b k_B T da^3, \quad (3.53)$$

which after some computation gives us

$$\frac{1}{2} dT = -T \frac{da}{a}, \quad (3.54)$$

which implies $T_m \propto a^{-2}$ after baryogenesis.

This is for monoatomic baryonic matter, right?

Let us start writing equations for the stellar interior. The continuity equation is

$$\partial_t \rho + \nabla \cdot (\rho \vec{v}) = 0, \quad (3.55)$$

and also we have the Euler equation

$$\partial_t \vec{v} + (\vec{v} \cdot \nabla) \vec{v} = -\frac{1}{\rho} \vec{\nabla} P - \vec{\nabla} \Phi, \quad (3.56)$$

which can be written using the convective time derivative:

$$\frac{D}{Dt} = \partial_t + \vec{v} \cdot \nabla_x = u^\mu \partial_\mu, \quad (3.57)$$

which allows us to write

$$\frac{D}{Dt} \rho + \rho \nabla \cdot \vec{v} = 0, \quad (3.58)$$

and

$$\frac{D}{Dt} \vec{v} = -\frac{1}{\rho} \nabla \Phi. \quad (3.59)$$

What happens to the entropy? we define the entropy density s by $S = s\rho$.

An isentropic process is one in which

$$\frac{DS}{Dt} = 0, \quad (3.60)$$

which in terms of the entropy density is

$$\partial_t s + \vec{v} \cdot \vec{\nabla} s = 0. \quad (3.61)$$

The external force is provided by the gravitational field:

$$\nabla^2 \Phi = 4\pi G \rho. \quad (3.62)$$

Jeans looked for a simple solution, an ansatz, called the background solution and then tried to perturb it: if it is stable than it was a good solution.

He started with $\rho = \text{const}$, $\vec{v} = 0$, $s = \text{const}$, $\Phi = \text{const}$.

It is obviously wrong! It cannot satisfy the Poisson equation.

However, we start from it and add some $\delta\rho$, $\delta\vec{v}$ (which we just call \vec{v}), δs and $\delta\Phi$; then we only keep the linear terms in these perturbations.

We find:

$$\partial_t \delta\rho + \rho_0 \vec{\nabla} \cdot \vec{v} = 0, \quad (3.63)$$

$$\partial_t \vec{v} = -\frac{1}{\rho_0} \nabla \delta P - \nabla \delta\Phi, \quad (3.64)$$

and

$$\nabla^2 \delta\Phi = 4\pi G \delta\rho, \quad (3.65)$$

and finally

$$\partial_t \delta s = 0. \quad (3.66)$$

We can expand

$$\delta P = \frac{\partial P}{\partial \rho} \delta\rho + \frac{\partial P}{\partial s} \delta s, \quad (3.67)$$

and here we define

$$c_s^2 = \frac{\partial P}{\partial \rho}. \quad (3.68)$$

We will then consider an exponential solution:

$$\delta\rho = \delta\rho_0 \exp\left(i\left(\vec{k} \cdot \vec{x} - \omega t\right)\right), \quad (3.69)$$

and similarly for \vec{v} , s , Φ .

We will see that we will need to stick to $\delta s = 0$, and find a dispersion relation with ω and \vec{k} : it will be

$$\omega^2 = c_s^2 k^2 - 4\pi G \rho, \quad (3.70)$$

so if the wavenumber is small enough we will have an imaginary ω .

Fri Nov 29 2019

We want to derive the Jeans instability criterion, starting from the structure equations:

1. continuity;
2. momentum conservation;
3. Poisson for the gravitational field;
4. entropy conservation.

Yesterday there was a mistake: the entropy conservation is actually

$$\frac{Ds}{Dt} + s \nabla \cdot \vec{v} = 0. \quad (3.71)$$

There is no need to make an ansatz for the pressure, since we can compute $P = P(\rho, s)$: we have the relation seen last time, involving the speed of sound.

Our *ansatz* is

$$\delta x_i = x_{i0} \exp(i\vec{k} \cdot \vec{x} - i\omega t), \quad (3.72)$$

with $x_i = \rho, \vec{v}, \Phi, s$.

Our equations become:

$$i\omega\delta\rho_0 + \rho_0 i\vec{k} \cdot \vec{v}_0 = 0 \quad (3.73a)$$

$$i\omega\vec{v}_0 = \frac{1}{\rho_0} i\vec{k} \left(c_s^2 \delta\rho_0 + \frac{\partial P}{\partial s} \delta s_0 \right) - i\vec{k} \delta\Phi_0 \quad (3.73b)$$

$$k^2 \delta\Phi = 4\pi G \delta\rho_0 \quad (3.73c)$$

$$\omega \delta s_0 = 0, \quad (3.73d)$$

so one class of solutions will have $\omega = 0$, that is, we consider time-independent ones:

$$\rho_0 i\vec{k} \cdot \vec{v}_0 = 0 \quad (3.74a)$$

$$0 = \frac{1}{\rho_0} \vec{k} \left(c_s^2 \delta\rho_0 + \frac{\partial P}{\partial s} \delta s_0 \right) - \vec{k} \delta\Phi_0 \quad (3.74b)$$

$$k^2 \delta\Phi = 4\pi G \delta\rho_0, \quad (3.74c)$$

and we have a result from Helmholtz: the fact that every velocity field can be decomposed into $\vec{v} = \nabla\Psi + \vec{T}$ with $\nabla \cdot \vec{T} = 0$.

As soon as we ask the perturbation to be time independent we only find vortical motions, with $\vec{v} \perp \vec{k}$.

Now we consider the velocity to be irrotational we consider the term $\delta s_0 = 0$, since $\omega \neq 0$ in general.

$$\omega\delta\rho_0 + \rho_0 \vec{k} \cdot \vec{v}_0 = 0 \quad (3.75a)$$

$$\omega \vec{v}_0 = \frac{1}{\rho_0} \vec{k} c_s^2 \delta \rho_0 - \vec{k} \delta \Phi_0 \quad (3.75b)$$

$$k^2 \delta \Phi = 4\pi G \delta \rho_0, \quad (3.75c)$$

which we can write as a linear system for the vector $\delta \rho_0, v_0, \delta \Phi_0$.

The 5x5 coefficient matrix is:

$$\begin{bmatrix} \omega & \rho_0 \vec{k} & 0 \\ \frac{1}{\rho_0} \vec{k} c_s^2 & \omega & \vec{k} \\ 4\pi G & 0 & k^2 \end{bmatrix}, \quad (3.76a)$$

which has determinant

$$\omega k^2 - \rho_0 \vec{k} \cdot \left(\frac{1}{\rho_0} \vec{k} c_s^2 - 4\pi G \vec{k} \right) = 0, \quad (3.77)$$

which gives the dispersion relation $\omega^2 = c_s^2 k^2 - \rho_0 4\pi G$.

A solution will be a combination of these. If $\omega^2 < 0$, which can happen because of the minus sign. If that happens, instead of a propagating wave we have a standing wave.

We have $\omega = (4\pi G \rho_0)^{1/2}$, so the characteristic time is

$$\tau = \frac{1}{\sqrt{4\pi G \rho_0}}, \quad (3.78)$$

The free fall time can be computed to be:

$$\tau_{\text{free fall}} = \left(\frac{3\pi}{32G\rho_0} \right)^{1/2}. \quad (3.79)$$

We have a typical Jeans wavenumber:

$$k_J^2 = \frac{4\pi G \rho_0}{c_s^2}, \quad (3.80)$$

corresponding to when the frequency becomes imaginary.

We can compare this to a plasma of charged particles, with the electrostatic potential instead of the gravitational field. In that case we find

$$\omega^2 = c_s^2 k^2 + \frac{4\pi n_e e^2}{m_e}, \quad (3.81)$$

where n_e is the number density of electrons, m_e is the electron mass.

We have the following analogies:

$$n_e \rightarrow \rho_0 / m \quad (3.82a)$$

$$m_e \rightarrow m \quad (3.82b)$$

$$e^2 \rightarrow G m^2 \quad (3.82c)$$

$$+ \rightarrow - ; \quad (3.82d)$$

the first of these are due to the inertial mass being equal to the gravitational mass; the plus becoming a minus is due to the fact that we do not have negative charge in the gravitational setting, so there cannot be a screening effect.

We now come back to the RW line element:

$$ds^2 = c^2 dt^2 - a^2(t) \left(\frac{dr^2}{1 - kr^2} + r^2 d\Omega^2 \right), \quad (3.83)$$

in which the most important thing is that a is a function of time.

We will now use a Newtonian approach: we ignore spatial curvature, that is, set $k = 0$: this is good for small enough scales. The thing is: the time scales of gravitational collapse are long, and we cannot ignore the effects of spacetime expansion.

We use a coordinate defined by $\vec{r} = a(t)\vec{x}$, so that the line element becomes

$$ds^2 = c^2 dt^2 - d|x|^2 - |x|^2 d\Omega^2. \quad (3.84)$$

right?

What will be done now could have been done by Newton: he did not just because he thought the universe was static. We have (dropping the vector sign, but still implying it):

$$u = \dot{r} = \dot{a}x + a\dot{x} = \frac{\dot{a}}{a}r + v, \quad (3.85)$$

where $v = a\dot{x}$ is called the peculiar velocity which galaxies and such can have.

Let us forget about the pressure gradient and just look at how the instability evolves: this will also apply to dark matter which has no pressure. We look at

$$\left[\frac{\partial \rho}{\partial t} \right]_{\vec{r}} + \nabla_{\vec{r}}(\rho \vec{u}) = 0, \quad (3.86)$$

the velocity equation is

$$\left[\frac{\partial \vec{u}}{\partial t} \right]_{\vec{r}} + (\vec{u} \cdot \nabla_{\vec{r}}) \vec{u} = -\nabla_{\vec{r}} \Phi, \quad (3.87)$$

and finally

$$\nabla_{\vec{r}}^2 = 4\pi G\rho. \quad (3.88)$$

Our expression for \vec{u} is $H\vec{r} + \vec{v}$, also we have

$$\rho(\vec{r}, t) = \rho_b(t) + \delta\rho(\vec{r}, t), \quad (3.89)$$

a background plus a perturbation.

To find the background, we assume that the density of matter is space independent but time dependent: then we find

$$\Phi_b(\vec{r}, t) = \frac{2\pi G}{3} \rho_b(t) r^2, \quad (3.90)$$

which has gradient

$$\nabla_{\vec{r}} \Phi_b = \frac{4\pi G}{3} \rho_b(t) \vec{r}, \quad (3.91)$$

and laplacian

$$\nabla_{\vec{r}}^2 \Phi_b = 4\pi G \rho_b, \quad (3.92)$$

but this diverges at $r \rightarrow \infty$! this is not an issue with the solution, but with Newtonian mechanics.

This is an elliptic differential equation: hyperbolic ones look like $\square \Phi = 0$, elliptic ones like $\nabla^2 \Phi = 0$, while parabolic ones (???)

We want to get equations in the comoving coordinates, not the local inertial ones. Take a generic function $f(\vec{r}, t)$, which can also be expressed with respect to (\vec{x}, t) .

Then, the difference between the derivatives is:

$$\left[\frac{\partial f}{\partial t} \right]_{\vec{x}} = \left[\frac{\partial f}{\partial t} \right]_{\vec{r}} + H(\vec{r} \cdot \nabla_{\vec{r}}) f. \quad (3.93)$$

Take the convective derivative

$$\frac{Df}{Dt} = \frac{\partial f}{\partial t}_{\vec{r}} + \frac{\partial f}{\partial \vec{r}} \cdot \vec{r}, \quad (3.94)$$

where $\vec{r} = \vec{u} = H\vec{r} + \vec{v}$. From what we saw, this can be expressed as

$$\frac{\partial f}{\partial t}_{\vec{r}} + H(\vec{r} \cdot \nabla_{\vec{r}}) + (\vec{v} \cdot \nabla_{\vec{r}}) f, \quad (3.95)$$

but we can also do it the other way round, keeping \vec{x} constant: that way, we find

$$\frac{Df}{Dt} = \frac{\partial f}{\partial t}_{\vec{x}} + \frac{\partial f}{\partial \vec{x}}, \quad (3.96)$$

and identifying the terms we get the desired relation, equation (3.93).

Skipping some passages, from the continuity equation we get:

$$\frac{\partial \rho}{\partial t}_{\vec{x}} + 3H\rho + \frac{1}{a} \nabla_{\vec{x}}(\rho \vec{v}) = 0, \quad (3.97)$$

while for the Euler equation we have:

$$\frac{\partial H\vec{r}}{\partial t}_{\vec{r}} + H^2(\vec{r} \cdot \nabla_{\vec{r}})\vec{r} = -\nabla_{\vec{r}} \Phi_b. \quad (3.98)$$

The divergence term is just $r^j \partial_j r^i = r^i$. Inserting the expression we know for the gravitational potential, whose gradient is proportional to \vec{r} , we get:

$$\vec{r} \left(\dot{H} + H^2 = -\frac{4\pi G}{3} \rho_b \right), \quad (3.99)$$

which must hold for any \vec{r} , so we drop it and recover one of the Friedmann equations! Note that we did not use any GR, everything was Newtonian.

We will get another result:

$$\frac{\partial \vec{v}}{\partial t} + H \vec{v} + \frac{1}{a} (\vec{v} \cdot \nabla_{\vec{x}}) \vec{v} = -\frac{1}{a} \nabla_{\vec{x}} \delta \Phi, \quad (3.100)$$

our new perturbed Euler equation; while the new Poisson equation is

$$\nabla^2 \delta \Phi = a^2 4\pi G \delta \rho. \quad (3.101)$$

Some cosmologists like to define

$$\delta(\vec{x}, t) = \frac{\delta \rho}{\rho_b} = \frac{\rho(\vec{x}, t) - \rho_b(t)}{\rho_b(t)}, \quad (3.102)$$

which can be anywhere from -1 to $+\infty$.

Then, we can interpret a negative $\delta \rho$ as a negative charge.

Why the Newtonian approach? Nobody knows how to write down the equations for a general relativistic self-gravitating fluid.

Is the Newtonian approximation a good one? In it, we have

$$\nabla^2 \delta \Phi = 4\pi G \rho_b a^2 \delta, \quad (3.103)$$

can we make the weak field approximation?

Typically

$$\delta \Phi \sim 4\pi G \rho_b a^2 \delta \lambda^2, \quad (3.104)$$

where λ is the variation scale, from the Laplacian.

From the Friedmann equation, $H^2 = \frac{8\pi G}{3} \rho_b$, then we have

$$\delta \Phi \sim \frac{3}{2} H^2 \delta a^2 \lambda^2 \sim \left(\frac{\lambda^2}{\lambda_{\text{hor}}^2} \right) \delta, \quad (3.105)$$

and $\lambda \sim \text{Mpc}$, the galactic perturbation scale, while $\lambda_{\text{hor}} \sim \text{Gpc}$.

As long as the perturbations are only galactic, the Newtonian approximation is good.

Thu Dec 05 2019

There will be a section on gravitational waves in cosmology in January.

We found the relation:

$$\vec{r} = a(\lambda)\vec{x}, \quad (3.106)$$

and

$$\vec{r} = \vec{u} = \dot{a}\vec{x} + a\dot{\vec{x}} = H\vec{r} + \vec{v}, \quad (3.107)$$

since $H = \dot{a}/a$.

An issue we have is the *redshift space distortion*, caused by *peculiar velocities* which cause the light we see to have redshift beyond the one caused by cosmology alone.

In comoving coordinates, we can use the regular Newtonian fluid dynamics equations.

Since we know that the gravitational instability is relevant when the gravitational force overcomes the pressure forces, we set the pressure gradient to be equal to zero. We have

$$\frac{\partial \rho}{\partial t}|_{\vec{r}} + \nabla_{\vec{r}}(\rho \vec{u}) = 0 \quad (3.108a)$$

$$\frac{d\vec{u}}{dt} + (\vec{u} \cdot \nabla_{\vec{r}})\vec{u} = 0 \quad (3.108b)$$

$$\nabla_{\vec{r}}^2 \Phi = 4\pi G \rho, \quad (3.108c)$$

and we consider a background $\Phi_b(\vec{r}, t)$, and $\rho = \bar{\rho}(t)$: we have a peculiar gravitational field beyond the Robertson-Walker background: we define

$$\phi = \Phi - \Phi_b, \quad (3.109)$$

and analogously we define by

$$\delta\rho = \rho - \bar{\rho} = \bar{\rho}(1 + \delta), \quad (3.110)$$

with $-1 \leq \delta \leq \infty$.

The Laplace equation for ϕ is given by

$$\nabla_{\vec{x}}^2 \phi(\vec{x}, t) = 4\pi G a^2(t) \bar{\rho}(t) \delta(\vec{x}, t), \quad (3.111)$$

and we actually *can* have negative quantities on the RHS. Not repulsive gravity, but kind of.

The relation between the derivatives is

$$\frac{\partial}{\partial r} = \frac{1}{a} \frac{\partial}{\partial x}, \quad (3.112)$$

therefore the continuity equation becomes:

$$\frac{\partial \rho}{\partial t} + 3H\rho + \frac{1}{a} \nabla_x(\rho \vec{v}) = 0, \quad (3.113)$$

we get that the density scales like a^{-3} if there is no velocity: then it becomes

$$\frac{\dot{\rho}}{\rho} + 3\frac{\dot{a}}{a} = 0. \quad (3.114)$$

For the momentum equation the calculation is longer, but in the end we find:

$$\frac{\partial \vec{v}}{\partial t} + H\vec{v} + \frac{1}{a}(\vec{v} \cdot \nabla_x)\vec{v} = -\frac{1}{a}\nabla_x\phi, \quad (3.115)$$

and once again we remark that we are not using GR because, in the absence of symmetry, that is definitely too difficult.

If we perturb, we will find that plane waves are not solutions anymore, because we will have time-dependent coefficients.

The perturbation of the density is:

$$\rho(\vec{x}, t) = \bar{\rho}(t)(1 + \delta(\vec{x}, t)), \quad (3.116)$$

and for the field:

$$\phi(\vec{x}, t) = \Phi - \Phi_b. \quad (3.117)$$

The derivative is:

$$\frac{\partial \rho}{\partial t} = \frac{\partial \bar{\rho}}{\partial t}(1 + \delta) + \bar{\rho} \frac{\partial \delta}{\partial t}, \quad (3.118)$$

and we can discard higher-order terms: simplifying the zeroth-order equation we find for the continuity equation:

$$\frac{\partial \delta}{\partial t} + \frac{1}{a}\vec{\nabla} \cdot \vec{v} = 0. \quad (3.119)$$

On the other hand, for the momentum equation we find:

$$\frac{\partial \vec{v}}{\partial t} + H\vec{v} = -\frac{1}{a}\nabla\phi, \quad (3.120)$$

and in order to solve these we can expand in Fourier space:

$$\delta(\vec{x}, t) = \frac{1}{(2\pi)^3} \int d^3\vec{k} \tilde{\delta}(t) \exp(i\vec{k} \cdot \vec{x}), \quad (3.121)$$

and similarly for \vec{v} and ϕ .

The actual quantities must be real: therefore we know that $\tilde{\delta}_k^*(t) = \tilde{\delta}_{-\vec{k}}(t)$.

Any vector field \vec{V} can be decomposed by the Helmholtz theorem into a gradient and a divergence:

$$\vec{V} = \nabla\Psi + \vec{T}, \quad (3.122)$$

where $\nabla \cdot \vec{T} = 0$.

Using the convective time derivative, we have

$$\frac{D\vec{\sigma}}{dt} + H\vec{\sigma} = -\frac{1}{a}\nabla\phi. \quad (3.123)$$

Inserting here the Helmholtz decomposition applied to the velocity, we find

$$\frac{D\vec{T}}{Dt} + H\vec{T} = 0, \quad (3.124)$$

which means that vorticity is conserved in fluid motion.

This means that any vorticity which might have been there at the beginning would have been diluted out over time.

We project the equations along the versor $\vec{u}_{\vec{k}} = \vec{k}/|\vec{k}|$.

Then the equation becomes

$$\dot{\delta}_{\vec{k}} + \frac{ik}{a}v_{\vec{k}} = 0, \quad (3.125)$$

and

$$\dot{v}_{\vec{k}} + Hv_{\vec{k}} = -\frac{i\vec{k}}{a}\phi_{\vec{k}}, \quad (3.126)$$

and finally

$$-k^2\phi_{\vec{k}} = 4\pi Ga^2\bar{\rho}\delta_{\vec{k}}. \quad (3.127)$$

Now, by differentiating in one single step we will get a linear second order differential equation for δ which is separated from the others.

Differentiating one more time and making the \vec{k} implicit we get:

$$\ddot{\delta} + \frac{ik}{a}\dot{v} - \frac{ik}{a^2}\dot{a} = 0 \quad (3.128a)$$

$$\ddot{\delta} + \frac{ik}{a}\left(-Hv - \frac{ik}{a}\phi\right) - \frac{ik}{a}Hv = 0 \quad (3.128b)$$

$$\ddot{\delta} - \frac{2ik}{a}Hv + \frac{k^2\phi}{a^2} = 0, \quad (3.128c)$$

but the last term is $-4\pi Ga^2\bar{\rho}\delta$: then in the end we find

$$\ddot{\delta} + 2H\dot{\delta} - 4\pi G\bar{\rho}\delta = 0, \quad (3.129)$$

and we use the following solutions of the background equations:

$$a(t) \propto t^{2/3} \quad (3.130a)$$

$$H = \frac{2}{3t} \quad (3.130b)$$

$$\bar{\rho} = \left(6\pi Gt^2\right)^{-1}, \quad (3.130c)$$

and plugging these in we find

$$\ddot{\delta} + \frac{4}{3t}\dot{\delta} - \frac{2}{3t^2}\delta = 0, \quad (3.131)$$

which will be a power of t : plugging t^α we find the equation

$$\alpha(\alpha - 1) + \frac{4}{3}\alpha - \frac{2}{3} = 0, \quad (3.132)$$

which gives $\alpha = -1, 2/3$.

The solution with $\alpha = 2/3$ is called the growing mode, while the one with $\alpha = -1$ is called the decaying mode. When the inflaton field becomes classical we lose a degree of freedom: this removes the decaying solutions. Another way to see it is to see that the decaying mode explodes as $t \rightarrow 0$.

Thus, we usually remove the decaying solution.

An interesting observation is the fact that the growing mode grows just as fast as the scale factor.

Inserting this into the other equations we find $v \propto t^\beta$ with $\beta = 1/3$ or $-4/3$, while $\phi \propto t^\gamma$ with $\gamma = 0, -5/2$.

understand what the equation

$$\chi_J = c_s \sqrt{\frac{\pi}{4\bar{\rho}}}, \quad (3.133)$$

The Jeans density is given by

$$\rho_J = \frac{3}{4\pi M^2} \left(\frac{3k_B T}{2G\bar{m}} \right)^3. \quad (3.134)$$

$$\chi_J = c_s \left(\frac{\pi}{4\bar{\rho}} \right)^{1/2}, \quad (3.135)$$

$$M_J = \frac{4\pi}{3} \bar{\rho} \left(\frac{\lambda}{2} \right)^3, \quad (3.136)$$

$$c_s^2 \sim \frac{k_B T}{\bar{m}}, \quad (3.137)$$

and now we are assuming that $\rho > \rho_J$.

Stars are made of baryons.

Molecular hydrogen:



$$H^- + H \leftrightarrow H_2 + e \quad (3.138b)$$

$$H + p \leftrightarrow H_2^+ + \gamma \quad (3.138c)$$

$$H_2^+ + H \leftrightarrow H_2 + p, \quad (3.138d)$$

and these processes are generally not very efficient: they start to be efficient at redshifts of about $z \sim 200$.

It is easier to get the Jeans density if the mass is high and the temperature is low.

The formation of stars is somewhat top-down: larger structures form first.

The critical energy density today is $\rho_{0c} \sim 10^{-29} \text{ gcm}^{-3}$, while for baryons we have $\rho_{0b} \sim 1 \times 10^{-31} \text{ gcm}^{-3}$.

On the other hand, at $z \sim 200$ we must multiply the density by a factor $(1+z)^3$ we get a density of $\rho_b(z \sim 200) \sim 10^{-22} \text{ kgm}^{-3}$.

If we take $M = 1000M_\odot$ and $T \approx 20 \text{ K}$, we have a Jeans critical density of around $10^{-20} \text{ kgm}^{-3}$.

Is 20 K a typical temperature then?

The matter starts free-falling: then the kinetic energy increases. However, it is used in chemical processes. That way matter can become opaque.

We have the equation

$$\frac{1}{2} \left(\frac{dr}{dt} \right)^2 = \frac{Gm_0}{r} - \frac{Gm_0}{r_0}, \quad (3.139)$$

which gives us a typical free-fall time of $2 \times 10^4 \text{ yr}$.

The energy needed to ionize all the hydrogen is

$$E = \frac{M}{2m_H} \epsilon_D + \frac{M}{m_H} \epsilon_I, \quad (3.140)$$

where M is the mass of the star, m_H is the mass of a hydrogen atom, while $\epsilon_D \approx 4.5 \text{ eV}$ and $\epsilon_U \approx 13.6 \text{ eV}$.

We are going from a radius R_1 to a radius R_2 : we get $R_1 \sim \sqrt[3]{3M/4\pi\rho_J}$, since that is the radius at which the collapse starts.

However, the R_1 term hardly contributes: the term R_2 is typically 4 orders of magnitude smaller.

Our objects must still become hotter and smaller in order to reach the 10^7 K needed for ignition.

The final equation we get comes from the virial theorem:

$$2E_k + E_{\text{gr}} = 0, \quad (3.141)$$

and

$$E_k = \frac{3}{2} N k_B T = \frac{3}{2} \frac{M}{m} k_B T \approx \frac{3M}{m_H} k_B T, \quad (3.142)$$

since $\bar{m} = 0.5m_H$: so in the end, equating the energies, we have

$$2 \times 3k_B T \frac{M}{m_H} = \frac{M}{m_H} \left(\frac{\epsilon_D}{2} + \epsilon_I \right), \quad (3.143)$$

which means that $k_B T = \frac{1}{12}(\epsilon_D + 2\epsilon_I) \sim 2.6 \text{ eV}$, which is still very low compared to the temperature we need.

Fri Dec 06 2019

Today we will discuss another possible *caveat* for star formation: the case where the star collapses but does not ignite.

After recombination the baryons' temperature decays more rapidly than the radiation's, so we get a temperature low enough to satisfy the instability criterion.

The temperature stays low during the collapse: the thermal energy is used in order to break the bonds in H_2 and ionize the hydrogen.

Then we have a gas which is opaque to radiation: there is scattering, which means we lose energy through radiation very slowly. Then the Virial theorem is very close to being true.

At the end of the process, even the mass does not matter anymore: we got 2.6 eV as the temperature regardless of the mass. This is equivalent to around $30 \times 10^3 \text{ K}$. We must compare this to the ignition temperature: that is in the order of keV ($15 \times 10^6 \text{ K}$ is equivalent to around 1.3 keV).

After free-fall the radius is on the order of 10^{10} m for a solar mass star, while the Sun's radius is smaller by two orders of magnitude.

Now we discuss the *conditions for stardom*: we need to account for the fermionic nature of protons and electrons, which will give us a maximum density due to the Pauli exclusion principle.

The way we will treat this today will be quite rough.

We know that the De Broglie wavelength is given by

$$\lambda = \frac{h}{p} = \frac{2\pi\hbar}{p}. \quad (3.144)$$

How do we calculate p ? we assume that the particles are nonrelativistic and apply $E_k = p^2/2m_e$.

Objects which will not satisfy the conditions we talk about today become brown dwarves.

The kinetic energy is of the order $k_B T$, therefore

$$p \sim \sqrt{2m_e k_B T}, \quad (3.145)$$

and the critical density is defined by

$$\rho_c = \frac{\bar{m}}{\lambda^3}, \quad (3.146)$$

where from the formula we found

$$\lambda = \frac{2\pi\hbar}{\sqrt{2m_e k_B T}}, \quad (3.147)$$

which gives approximately

$$\rho_c \sim \bar{m} \frac{(m_e k_B T)^{3/2}}{(2\pi\hbar)^3}, \quad (3.148)$$

and from the virial theorem $2E_k + E_{\text{gr}} = 0$, with

$$E_k = \frac{3}{2} N k_B T = \frac{3}{2} \frac{M}{\bar{m}} k_B T, \quad (3.149)$$

while

$$E_{\text{gr}} = -\frac{GM^2}{R}, \quad (3.150)$$

which means

$$3Nk_B T = \frac{GM^2}{R}, \quad (3.151)$$

and we can rewrite this as

$$\frac{3k_B T}{\bar{m}} = \frac{GM}{R}, \quad (3.152)$$

and we can express the mass as

$$M = \frac{4}{3} \pi \bar{\rho} R^3, \quad (3.153)$$

so

$$\frac{1}{R} = \left(\frac{4\pi}{3} \frac{\bar{\rho}}{M} \right)^{1/3}, \quad (3.154)$$

which gives us the result

$$k_B T = \frac{GM\bar{m}}{3} \left(\frac{4\pi}{3} \frac{\bar{\rho}}{M} \right)^{1/3}, \quad (3.155)$$

and if we substitute the critical density in for $\bar{\rho}$ we will get the maximum possible temperature allowed at a given mass.

This yields

$$k_B T = \frac{GM\bar{m}}{3} \left(\frac{4\pi}{3M} \right)^{1/3} \frac{\bar{m}^{1/3}}{(2\pi\hbar)} (m_e k_B T)^{1/2}, \quad (3.156)$$

which it is convenient to square:

$$(k_B T)^2 = \frac{G^2 M^2 \bar{m}^2}{9} \left(\frac{4\pi}{3M} \right)^{2/3} \frac{\bar{m}^{2/3}}{(2\pi\hbar)^2} m_e k_B T, \quad (3.157)$$

so we can simplify, and up to an order-1 constant

$$k_B T = \frac{G^2 \bar{m}^{8/3} M^{4/3}}{(2\pi\hbar)^2}, \quad (3.158)$$

and inserting the ignition temperature of around 1 keV we get $M_{\min} \sim 0.08 M_\odot$.

This is confirmed experimentally.

Let us consider the Sun. Its mass is $M_\odot \approx 1.99 \times 10^{30}$ kg, the radius is $R_\odot \approx 6.96 \times 10^8$ m, the electromagnetic luminosity is $L_\odot = 3.86 \times 10^{26}$ W.

The age of the Sun is around $t_\odot \approx 4.55 \times 10^9$ yr, which is comparable to the age of the Universe.

The central density is $\rho_c \approx 1.48 \times 10^5$ kgm⁻³, while the central temperature is $T_c = 1.56 \times 10^7$ K, and the central pressure is around $P_c = 2.29 \times 10^{16}$ Pa.

The effective temperature is around $T_E \approx 5780$ K.

Definition?

What is the corresponding free fall time? It is much shorter than the age of the Sun: the Sun is not in free fall.

What is the number?

We know that

$$\langle P \rangle = -\frac{1}{3} \frac{E_{\text{gr}}}{V}, \quad (3.159)$$

where $E_{\text{gr}} = -GM^2/R$ while $V = 4\pi R^3/3$: plugging the Sun's numbers we get

$$\langle P \rangle = 10^{14} \text{ Pa}, \quad (3.160)$$

100 times less than the central density.

The density of the Sun is actually very similar to the density of water.

Was it correct to use nonrelativistic equations? (???)

$$\langle P \rangle = \frac{\bar{\rho}}{\bar{m}} k_B T_I, \quad (3.161)$$

where T_I is the mean internal temperature of the Sun. The value of $\bar{m} \approx 0.61$ instead of 0.5 when considering the proper composition of the Sun.

We get

$$k_B T_I \approx \frac{GM_\odot \bar{m}}{3R_\odot} \approx 1.5 \text{ keV} \approx 6 \times 10^6 \text{ K}. \quad (3.162)$$

We have evidence that the Sun is a blackbody, we can write

$$L_\odot = 4\pi R_\odot \sigma T_E^4, \quad (3.163)$$

where σ is Stefan's constant: $\sigma \approx 5.67 \times 10^{-8} \text{ W m}^{-2} \text{ K}^{-4}$.

In principle we could define another quantity:

$$L'_{\odot} = 4\pi R_{\odot}^2 \sigma T_I^4, \quad (3.164)$$

which does not fit the data. Why is this?

We must consider a photon which comes from the interior of the Sun and goes towards the outside: it will follow a random walk scattering many times. However, depending on the density of electron in the outer regions (electron scattering dominates in the outside), the last scattering is in the outermost regions of the star.

We need to deal with the Fourier equation. First of all, we need to write the Langevin equation: this is connected to the work by Einstein on Brownian motion.

These are processes in which there is a "random" force (due to the fact that our description of the microscopic state is probabilistic), and possibly deterministic forces.

What is the stuff about viscosity?

The equation in the end is like:

$$\dot{\vec{x}} = \vec{\eta}, \quad (3.165)$$

What?

with $\langle \eta_i(t) \eta_j(t') \rangle = 2D \delta_{ij} \delta(t - t')$: there is complete uncorrelation. This is the Markov property: the process has no memory.

This gives us a Gaussian distribution of the positions of the particles, and we get the equation:

$$\frac{\partial P}{\partial t} = D \nabla^2 P, \quad (3.166)$$

where $P = P(\vec{x}, t)$. This is a *parabolic equation*, the *Fokker-Planck* formula. We need to give it both initial conditions and boundary conditions.

What is P ?

There are three kinds of boundary conditions:

1. nothing: free boundary, the solution can diverge;
2. a reflective boundary: particles "bounce back", in order to deal with this we use the images method, as in electromagnetism;
3. an absorbing boundary: particles disappear if they reach the boundary.

This is equivalent to the Fourier transport equation.

The solution without boundary is a Gaussian with variance $\langle x^2 \rangle = 2Dt$ and centered around zero.

If $\sigma = \sqrt{\langle x^2 \rangle}$ becomes greater than the boundary then most of the particles have escaped.

\vec{D} is the displacement vector of the Sun: it is

$$\vec{D} = \sum_i \vec{l}_i, \quad (3.167)$$

where the \vec{l}_i are the displacement vectors of its various steps.

$$\langle \vec{D}^2 \rangle = \sum_i \langle \vec{l}_i^2 \rangle + \sum_{i < j} \langle \vec{l}_i \cdot \vec{l}_j \rangle, \quad (3.168)$$

but if we have isotropy then the scalar products have mean zero. This might be unphysical since the steps are larger at the boundary than at the center...

Then we find:

$$\langle \vec{D}^2 \rangle = Nl^2 = R_\odot^2, \quad (3.169)$$

where $N = R_\odot^2/l^2$.

The time it takes for a photon to cover a distance l is $t = l/c$. Then we have $t_{RW} = Nt = R_\odot^2 l / (l^2 c)$, which means

$$t_{RW} = \frac{R_\odot^2}{cl}, \quad (3.170)$$

while in direct flight the photon would take $t_0 = R_\odot/c$: their ratio is

$$\frac{t_{RW}}{t_0} = \frac{R_\odot}{l}, \quad (3.171)$$

and then

$$L_\odot = L'_\odot \frac{l}{R_\odot}, \quad (3.172)$$

which means

$$T_E = \left(\frac{l}{R_\odot} \right)^{1/4} T_I, \quad (3.173)$$

so we can gather l by knowing the other three parameters: we get $l = 1$ mm. This is actually an average of the mean free paths.

$$L = L' \frac{l}{R} = 4\pi R_\odot^2 \sigma T_I^4 \frac{l}{R}, \quad (3.174)$$

and we know that $k_B T_I = \frac{GM\bar{m}}{3\hbar}$: replacin this we find

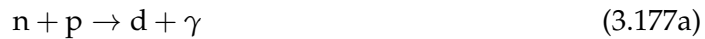
$$L = 4\pi R_\odot^2 \sigma \left(\frac{GM\bar{m}}{3Rk_B} \right)^4 \frac{l}{R} = \frac{(4\pi)^2}{3^5} \frac{\sigma}{k_B^4} G^4 \bar{m}^4 \rho l M^3. \quad (3.175)$$

What is the typical lifespan of a star? This gives us $L \sim M^3$, which means $\tau \sim M/L \sim M^{-2}$. Observationally, this matches the data pretty well.

Now we discuss thermonuclear fusion. The reactions we need are



which involves the weak interaction (for the first process: $\tau \sim 5 \times 10^9$ yr), EM interaction (for the second process: $\tau \sim 1$ s) and strong interaction (for the third process: $\tau \sim 3 \times 10^5$ yr). On the other hand, the process



which does not use the weak interactions. However, later there are no more free neutrons: this means that even if it is slower the first process is the only one which can happen.

The net balance is 4 protons in, 1 ${}^4\text{He}$ out.

Thu Dec 12 2019

Check room availability for the first week of January.

Tomorrow / next week we will discuss some black holes and neutron stars.

Last week we discussed the nuclear processes which occur in the center of the Sun when $T \sim 10^7$ K is reached.

The minimum mass for a star in order to ignite fusion, as we found, is around $0.08M_\odot$.

The most important difference between the processes outlined last week is in the first equation of each: there are very few free neutrons.

So, the weak interaction process dominates: since it is so slow, it has a very low power density: $P = 4 \times 10^{26}$ W, but we need to take its ratio to the volume of the Sun. This gives a density lower than that of a human.

For each ${}^4\text{He}$ nucleus we get 26 MeV, and we need 4 protons to make it. Then, this gives us the number of protons per second the Sun uses in order to produce the power it does. We will use the following relation:

$$1 \text{ MeV} \approx 1.78 \times 10^{-30} \text{ kg} \approx 1.6 \times 10^{-13} \text{ J}, \quad (3.178)$$

and then, in SI units: we get

$$\frac{4 \times 10^{26}}{2.6 \times 10^{-13} / 4} \approx 4 \times 10^{38} \frac{\text{protons}}{\text{s}}. \quad (3.179)$$

For each process we also emit one electron neutrino, and we need to do the first two steps of the process twice for each ${}^4\text{He}$ nucleus, so we are producing 2×10^{38} neutrinos per second.

A proton's mass is around $m_p \approx 1 \text{ GeV} \approx 1.78 \times 10^{-27} \text{ kg}$, so in the Sun there are around 10^{56} of them: this means that the typical lifetime of the Sun is around 10^{10} years.

What is the final state of the Sun? After hydrogen runs out, the Sun should start the next process: helium burning. The core contracts and heats. However we also have the degeneracy pressure.

There is a boundary, the Chandrasekar mass $M_C \approx 1.4M_\odot$, between the final fate of the star being a white dwarf or a neutron star.

From the point of view of the state of matter, brown dwarves and white dwarves are very similar.

So the core becomes denser and hotter and starts burning helium. The external parts, by the virial theorem, must then expand.

So the possibilities, in order of mass, are white dwarf, neutron star, black hole.

The matter which is expelled can form a *planetary nebula*.

Process	Fuel	Products	T_{\min}	M_{\min}
Hydrogen burning	Hydrogen	Helium	10^7 K	$0.08M_\odot$
Helium burning	Helium	Carbon, Oxygen	10^8 K	$0.5M_\odot$
Carbon burning	Carbon	Oxygen, Neon, Sodium	$5 \times 10^8 \text{ K}$	$8M_\odot$
Neon burning	Neon	Magnesium, Oxygen	10^9 K	$9M_\odot$
Oxygen burning	Oxygen	Magnesium to Sulphur	$2 \times 10^9 \text{ K}$	$10M_\odot$
Silicon burning	Silicon	Iron and nearby elements	$3 \times 10^9 \text{ K}$	$11M_\odot$

Figure 3.1: Solar processes.

A plot of the binding energy per nucleon $B = E - M_{\text{nucleons}}$ shows that it is maximum for iron.

Something about the possibility to have ${}^8\text{Be}$ be stable in order for potassium to be formed.

When the nucleus cannot reach the temperature needed for the next process, the collapse is stopped by electron degeneracy.

The final radius of the red giant phase of the Sun is around 70 times the radius of the Sun, while the white dwarf phase is 70 times smaller.

Is this correct? I couldn't quite hear.

Our equation of hydrostatic equilibrium is:

$$\frac{dP}{dr} = -\frac{Gm(r)\rho(r)}{r^2}, \quad (3.180)$$

and the equation giving the variation of the mass is $\frac{dm}{dr} = 4\pi r^2 \rho(r)$. So we get

$$\frac{r^2}{\rho(r)} \frac{dP}{dr} = -Gm(r), \quad (3.181)$$

which can be restated as

$$\frac{d}{dr} \left(\frac{r^2}{\rho(r)} \frac{dP}{dr} \right) = -G \frac{dm}{dr} = -4\pi G \rho(r) r^2, \quad (3.182)$$

more commonly stated as

$$\frac{1}{r^2} \frac{d}{dr} \left(\frac{r^2}{\rho(r)} \frac{dP}{dr} \right) = -4\pi G \rho(r), \quad (3.183)$$

which holds if we have hydrostatic equilibrium.

Commonly the equation of state used for this is called *polytropic*:

$$P = k \rho^{\frac{n+1}{n}}, \quad (3.184)$$

where $k = \text{const}$ and $n = 1/(\gamma - 1)$: so if $\gamma = 5/3$, which holds for a monoatomic gas, then $n = 3/2$ while if $\gamma = 4/3$, which holds for an ultrarelativistic gas, then $n = 3$.

Using this law, we can write this equation in terms of either only the density or only the pressure.

Let us have only the density:

$$\frac{1}{r^2} \frac{d}{dr} \left(\frac{r^2}{\rho(r)} \frac{d}{dr} \left(\rho^{\frac{n+1}{n}} \right) \right) = -4\pi G \rho(r), \quad (3.185)$$

and we need two boundary conditions: we can fix the central density $\rho(r=0) = \rho_c$ and the derivative of the density at the center:

$$\frac{d\rho_c}{dr}(r=0) = 0. \quad (3.186)$$

This is because:

$$\rho^{1/n} \frac{d\rho}{dr} \propto -\frac{Gm(r)}{r^2} \rho(r), \quad (3.187)$$

and the mass is given by $m(r) \sim \rho_c r^3$: therefore $\rho^{1/n} d\rho/dr \propto r$, so it makes sense to set the derivative to zero.

The radius of the star is defined as the one at which the density goes to 0: $\rho(R) = 0$, and the mass of the star is given by $m(R) = M$.

These equations can be solved numerically. Also, we can use an ansatz.

A model by Clayton 1986: $P_c = 2 \times 10^{16}$ Pa, and since

$$\langle P \rangle = -\frac{1}{3} \frac{E_{\text{gr}}}{V_{\odot}}, \quad (3.188)$$

where $E_{\text{gr}} = GM_{\odot}^2/R_{\odot}$, and $V_{\odot} = \frac{4}{3}\pi R_{\odot}^3$: therefore the mean pressure is approximately 1/200 times the pressure at the center.

We can write the equation of hydrostatic equilibrium as

$$\frac{dP}{dr} = -\frac{Gm\rho}{r^2} \approx -\frac{4\pi G}{3}\rho_c^2 r, \quad (3.189)$$

so the pressure gradient goes to zero linearly in r . So, we know that

$$\frac{dP}{dr}(r=0) = 0, \quad (3.190)$$

also, as $r \rightarrow R$ we get $dP/dr \rightarrow 0$ as well, since it is proportional to $\rho(r)$ in that region.

The pressure gradient around the center follows the density. Its variation is much larger near the center than on the outside. The ansatz by Clayton is

$$\frac{dP}{dr} = -\frac{4\pi}{3}G\rho_c^2 r \exp\left(-\frac{r^2}{a^2}\right), \quad (3.191)$$

where the parameter a has the dimensions of a length, and we take it to be $a \ll R$. This model is quite accurate near the center, not so much near the surface!

Integrating we find:

$$P(r) = \frac{2\pi}{3}G\rho_c^2 a^2 \left(\exp\left(-\frac{r^2}{a^2}\right) - \exp\left(-\frac{R^2}{a^2}\right) \right), \quad (3.192)$$

so that the pressure is exactly zero at the surface: $P(R) = 0$. We will check *a posteriori* that this model makes sense for a typical star.

Do we set the pressure to stay at zero for $r > R$? Maybe does not matter...

We have the relation

$$Gm(r) dm = -4\pi r^4 dP, \quad (3.193)$$

which can be integrated in order to give the following result:

$$\frac{1}{2}Gm^2(r) = -4\pi \int_0^r d\tilde{r} \tilde{r}^4 \frac{dP}{d\tilde{r}}. \quad (3.194)$$

We recover $m(r)$ by taking the square root:

$$m(r) = \frac{4\pi a^3}{3}\rho_c \Phi(x), \quad (3.195)$$

where $x = r/a$ and $\Phi(x)$ is defined from the integral from before:

$$\Phi^2(x) = 6 \int_0^x dy y^5 e^{-y^2} = 6 - 3(x^4 + 2x^2 + 2)e^{-x^2}, \quad (3.196)$$

and the second term is almost zero near the surface if $a \ll R$: the power series starts from the sixth power (? what?).

The density profile then is given by:

$$\rho(r) = \frac{1}{4\pi r^2} \frac{dm}{dr} = \rho_c \left(\frac{x^3 e^{-x}}{\Phi(x)} \right), \quad (3.197)$$

and we can also recover the temperature profile from the ideal gas law:

$$T(r) = \frac{\bar{m}}{k_B} \frac{P(r)}{\rho(r)}, \quad (3.198)$$

and, when $x \ll 1$ we find

$$\Phi(x) \sim \left(x^6 - \frac{3}{4}x^8 + \frac{3}{10}x^{10} - \frac{1}{12}x^{12} + \dots \right)^{1/2}, \quad (3.199)$$

and inserting this we find

$$\rho(r) \approx \rho_c \left(1 - \frac{5}{8} \frac{r^2}{a^2} + \dots \right), \quad (3.200)$$

and

$$T(r) \approx T_c \left(1 - \frac{3}{8} \frac{r^2}{a^2} + \dots \right), \quad (3.201)$$

which allows us to get the total mass

$$M = m(R) = \frac{4\pi\rho_c a^3}{3} \Phi(R/a) \approx \frac{4\pi\rho_c a^3 \sqrt{6}}{3}, \quad (3.202)$$

and lots of it are within a radius a :

$$\rho(a) = 0.53\rho_c \quad \text{and} \quad m(a) = 0.28M_\odot. \quad (3.203)$$

We also have the relation $a = R_\odot/5.4$, and

$$\frac{\langle \rho \rangle}{\rho_c} \sim \left(\frac{a}{R} \right)^3 \sim 1. \quad (3.204)$$

Where does this come from?

We have the relations

$$m(r) = \frac{4\pi a^3}{3} \rho_c \Phi(x), \quad (3.205)$$

and

$$P_c = \frac{2\pi}{3} G \rho_c^2 a^2. \quad (3.206)$$

Then we can combine these:

$$P_c = \left(\frac{\pi}{36}\right)^{1/3} GM^{2/3} \rho_c^{4/3}. \quad (3.207)$$

The factor $(\pi/36)^{1/3} \sim 0.44$. Changing γ we get: for $\gamma = 5/4$ the factor is 0.48. We use the ideal gas relation: then we find

$$k_B T_c = \left(\frac{\pi}{36}\right)^{1/3} G \bar{m} M^{1/3} \rho_c^{1/3}. \quad (3.208)$$

Then, we can figure out which processes actually can happen.

Fri Dec 13 2019

We discuss the maximum mass of stars.

We were able to give meaning to the parameter a : the pressure is given by

$$P(r) = \frac{2\pi}{3} G \rho_c^2 a^2 \left(\exp\left(-\frac{r^2}{a^2}\right) - \exp\left(-\frac{R^2}{a^2}\right) \right), \quad (3.209)$$

so we can see that

$$a = \left(\frac{3M}{4\pi \rho_c \sqrt{6}} \right)^{1/3}, \quad (3.210)$$

and then we got an expression for the central pressure P_c : the parameter multiplying it is approximately 0.44, while more accurate models give: if $\gamma = 5/3$ (ideal gas) we get 0.48 while if $\gamma = 4/3$ (ultrarelativistic) we get 0.36.

We have the relation

$$P_c = \frac{\rho_c}{\bar{m}} k_B T_c, \quad (3.211)$$

which we use to get the last relation from last time.

This allows us to get some figures for main sequence (hydrogen burning) stars. In the Hertzsprung Russel diagram we plot L/L_\odot versus T_{eff} , the latter increasing right to left.

The Main Sequence runs from the upper left to the lower right, we have Red Giants on the upper right and White Dwarves on the lower left. Most of the stars are on the Main Sequence: the hydrogen burning phase lasts a long time.

What is the maximum mass for Main Sequence stars? A star becomes unstable when most of its material becomes ultrarelativistic: then, its total energy goes from a negative value to 0 and the adiabatic index approaches $4/3$.

Suppose that the central energy is partly given by radiation and partly by matter. We write

$$P_c = P_\rho + P_r = \beta P_c + (1 - \beta) P_c, \quad (3.212)$$

where the terms of the two sums exactly correspond to each other, and

$$\beta P_c = P_\rho = \frac{\rho_c k_B T_c}{\bar{m}} \quad (3.213)$$

and

$$(1 - \beta) P_c = P_r = \frac{1}{3} a, \quad (3.214)$$

where

$$a = \frac{\pi^2 k_B^2}{15 \hbar^3 c^3}. \quad (3.215)$$

So we have

$$(\beta P_c)^4 = \frac{\rho_c^4}{\bar{m}^4} (k_B T_c)^4, \quad (3.216)$$

and

$$\frac{1 - \beta}{\beta^4} P_c^{-3} = \frac{a}{3} \left(\frac{k_B \rho_c}{\bar{m}} \right)^{-4}. \quad (3.217)$$

Inverting this we can eliminate the temperature dependence:

$$P_c = \left(\frac{3}{a} \frac{1 - \beta}{\beta^4} \right)^{1/3} \left(\frac{k_B \beta}{\bar{m}} \right)^{4/3} = \left(\frac{\pi}{36} \right)^{1/3} G M^{1/3} \rho_c^{4/3}, \quad (3.218)$$

so

$$\left(\frac{\pi}{36} \right)^{1/3} G M^{2/3} = \left(\frac{3}{a} \frac{(1 - \beta)}{\beta^4} \right)^{1/3} \left(\frac{k_B}{\bar{m}} \right)^{4/3}, \quad (3.219)$$

so as β decreases, M increases.

[Plot of $1 - \beta$ versus M/M_\odot , showing this.]

Now we deal with the degenerate electron gas in stars.

The distribution function is approximately given by

$$f(p) \propto \frac{1}{\exp\left(\frac{\epsilon_p - \mu}{k_B T}\right) + 1}, \quad (3.220)$$

where $\epsilon_p = c(p^2 + m^2 c^2)^{1/2}$. Then, as $T \rightarrow 0$ we get: $f(\epsilon_p) = 1$ if $\epsilon_p \leq \epsilon_F$ and $f(\epsilon_p) = 0$ if $\epsilon_p > \epsilon_F$; we can express this energy in terms of the momentum:

$$\epsilon_F^2 = c^2 p_F^2 + m^2 c^4. \quad (3.221)$$

The number of electrons is given by

$$n_e = 2 \int_0^{p_F} dp p^2 4\pi \frac{1}{h^3} = \frac{8\pi}{3} \left(\frac{p_F}{h} \right)^3, \quad (3.222)$$

where we have a factor of 2 to account for the spin-1/2 nature of the electrons. This means that

$$p_F = \left(\frac{3n}{8\pi} \right)^{1/3} h. \quad (3.223)$$

The density is given by

$$\rho = \frac{2}{h^3} \int_0^{p_F} dp p^2 4\pi \epsilon_p, \quad (3.224)$$

and we can consider either the nonrelativistic or the ultrarelativistic limit. In the nonrelativistic limit we find

$$\epsilon_p = mc^2 + \frac{p^2}{2m}, \quad (3.225)$$

so the result becomes

$$\rho = n \left(mc^2 + \frac{3}{10} \frac{p_F^2}{m} \right), \quad (3.226)$$

and we know that for a nonrelativistic gas the pressure is given by

$$P = \frac{2}{3} \frac{E_k}{V}, \quad (3.227)$$

where E/V is the kinetic energy density.

$$P = n \frac{p_F^2}{5m}, \quad (3.228)$$

where does this come from?

$$P = k_{NR} n^{5/3}, \quad (3.229)$$

where

$$k_{NR} = \frac{h^2}{5m} \left(\frac{3}{8\pi} \right)^{2/3}. \quad (3.230)$$

In the relativistic case, on the other hand, we get $\epsilon_p \approx cp$, and $\rho = \frac{3}{4} n \rho_{Fc}$. In this case, we also know that the pressure becomes

$$P = \frac{1}{3} \frac{E_k}{V}. \quad (3.231)$$

So we find

$$P = K_{VR} n^{4/3}, \quad (3.232)$$

where

$$K_{VR} = \frac{hc}{4} \left(\frac{3}{8\pi} \right)^{1/3}. \quad (3.233)$$

We make a plot: on the x axis we have the number density in m^{-3} , on the y axis we have the temperature in K.

We divide the plot into:

1. Classical UR: $P \propto nk_B T$;
2. classical NR (like the Sun);
3. degenerate NR: $P = K_{NR} n^{4/3}$
4. degenerate UR.

Classical vs degenerate is marked by a line similar to $T \sim n$, while we have NR for both T and n lower than certain critical values (since a degenerate gas can become ultrarelativistic even at low temperatures! this is the point).

We have

$$P_c = \frac{\rho_c}{\bar{m}} k_B T_c, \quad (3.234)$$

and

$$k_B T_c = \left(\frac{\pi}{36} \right)^{1/3} G \bar{m} M^{2/3} \rho_c^{1/3}, \quad (3.235)$$

and it can be (easily?) shown that

$$\bar{m} = 2m_H \times \frac{1}{1 + 3x_1 + 0.5x_4}, \quad (3.236)$$

where $x_{1,4}$ are the concentrations of hydrogen and helium respectively.

To estimate the maximum achievable central temperature we do:

$$P_c = k_{NR} n_e^{5/3} + n_i k_B T_c, \quad (3.237)$$

where $n_e = n_i = \rho_c / \bar{m}_H$.

$$\left(\frac{\pi}{36} \right)^{1/3} G M^{2/3} \rho_c^{4/3} = k_{NR} \left(\frac{\rho_c}{m_H} \right)^{5/3} + \frac{\rho_c}{m_H} k_B T_c, \quad (3.238)$$

which implies

$$k_B T_c = \left(\frac{\pi}{36} \right)^{1/3} G m_H M^{2/3} \rho_c^{1/3} - k_{NR} \left(\frac{\rho_c}{m_H} \right)^{2/3}, \quad (3.239)$$

which is in the shape $k_B T_c = A\rho_c^{1/3} - B\rho_c^{2/3}$: we can find a maximum of this function, which comes out to be at $\rho_c = (A/2B)^3$, where we have

$$k_B T_c = \frac{A^2}{2B} = \left(\frac{\pi}{36}\right)^{2/3} \frac{G^2 m_H^{8/3}}{4k_{NR}} M^{4/3}. \quad (3.240)$$

Now, we can set this temperature to be larger than the ignition temperature for any process we want, to see whether it will happen.

The minimum mass is

$$M_{\min} = \left(\frac{36}{\pi}\right)^{1/2} \left(\frac{4k_{NR}}{G^2 m_H^{8/3}}\right)^{3/4} (k_B T_{\text{ign}})^{3/4}. \quad (3.241)$$

The potential energy between two hydrogen nuclei separated by a distance equal to their quantum wavelength is

$$E_g = -\frac{Gm_H^2}{r} = -\frac{Gm_H^3 c}{\hbar}, \quad (3.242)$$

where we inserted $r = \hbar/m_H c$. This corresponds to an energy $E = m_H c^2$, and we have a value

$$\alpha_G = \frac{E_g}{E} = \frac{Gm_H^2}{\hbar c} \sim 5.9 \times 10^{-39}. \quad (3.243)$$

For electromagnetic interaction, we get

$$\alpha_{EM} = \frac{e^2}{4\pi\epsilon_0 \hbar c} \approx \frac{1}{137}, \quad (3.244)$$

which is *much greater*.

Then we find

$$M_{\min} \approx 16 \left(\frac{k_B T_{\text{ign}}}{m_e c^2}\right)^{3/4} \alpha_G^{-3/2} m_H. \quad (3.245)$$

If $T_{\text{ign}} \sim 1.5 \times 10^6$ K, one tenth of the temperature of the Sun, we find

$$M_{\min} \sim 0.03 \alpha_G^{-3/2} m_H, \quad (3.246)$$

while for the maximum mass, from before with $\beta = 0.5$ and $\bar{m} = 0.61 m_H$ we get

$$M_{\max} \approx 56 \alpha_G^{-3/2} m_H, \quad (3.247)$$

so this hints to the fact that $m_* = \alpha_G^{-3/2} m_H$ is an important characteristic mass. This is around $1.85 M_\odot$.

This corresponds to a number of particles:

$$N_* = \frac{m_*}{m_H} \approx 2 \times 10^{52}. \quad (3.248)$$

Suppose the core of a star is held together by the pressure of degenerate electrons: we discuss *white dwarves*. We define

$$n_e = Y_e \frac{\rho_c}{m_H}, \quad (3.249)$$

where $Y_e = (1 + x_1)/2$. The pressure is

$$P = k_{NR} n_e^{5/3} = k_{NR} \left(\frac{Y_e \rho_c}{m_H} \right)^{5/3}, \quad (3.250)$$

which must be compared to

$$P_c = \left(\frac{\pi}{36} \right)^{1/3} G M^{2/3} \rho_c^{4/3}, \quad (3.251)$$

and we assume that $P = P_c$: then we get

$$\rho_c \approx \frac{3.1}{Y_e^5} \frac{M}{m_*} \frac{m_H}{(h/m_e c^2)^3}. \quad (3.252)$$

Missing square on the M/m_* ?

The pressure is

$$P = k_{UR} n_e^{4/3} = k_{UR} \left(\frac{Y_e \rho_c}{m_H} \right)^{4/3}, \quad (3.253)$$

so we get

$$k_{UR} \left(\frac{Y_e \rho_c}{m_H} \right)^{4/3} \approx \left(\frac{\pi}{36} \right)^{1/3} G M^{2/3} \rho_c^{4/3}, \quad (3.254)$$

so we can see that in this limit the expression becomes independent of ρ_c . This will give us a limit mass: the Chandrasekhar mass.

$$M_{CH} = \left(\frac{36}{\pi} \right)^{1/2} \left(\frac{Y_e}{m_H} \right)^2 \left(\frac{k_{UR}}{G} \right)^{3/2} \approx 2.3 Y_e^2 m_* \approx 4.3 Y_e^2 M_\odot \approx 1.4 M_\odot. \quad (3.255)$$

This is the maximum mass of a white dwarf to remain stable, held together by the degeneracy pressure of electrons. Above this, it becomes a neutron star.

Thu Dec 19 2019

The lessons on gravitational waves will be on the 8th and 9th of January, from 14:30 to 16:30, in rooms LUF2 and P2B respectively.

We want to find the Chandrasekhar limit in a more precise manner.

The number density of electrons is given by

$$n_e = Y_e \frac{\rho_c}{m_H}, \quad (3.256)$$

in the nonrelativistic case the pressure is given by

$$P_c = k_{\text{NR}} n_e^{5/3} = k_{\text{NR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{5/3}, \quad (3.257)$$

but we can also derive it by

$$P_c = k_{\text{NR}} n_e^{5/3} = k_{\text{NR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{5/3}, \quad (3.258)$$

and equating these we find:

$$k_{\text{NR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{5/3} = k_{\text{NR}} n_e^{5/3} = k_{\text{NR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{5/3}, \quad (3.259)$$

which implies

$$\rho_c = \frac{3.1}{Y_e^5} \left(\frac{M}{M_*} \right)^2 \frac{m_H}{(h/m_e c^2)^3}, \quad (3.260)$$

where $\alpha_G = G m_H^2 / (\hbar c) \approx 5.9 \times 10^{-30}$, while

$$m_* = \alpha_G^{-3/2} m_H = 1.85 M_\odot, \quad (3.261)$$

while for the ultrarelativistic case we get

$$P_C = k_{\text{UR}} n_e^{4/3} = k_{\text{UR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{4/3}, \quad (3.262)$$

so in this particular case the density ρ_c simplifies from the equations: we get a critical mass

$$M_{\text{CHANDRA}} = \left(\frac{36}{\pi} \right)^{1/2} \left(\frac{Y_e}{m_H} \right)^2 \left(\frac{k_{\text{UR}}}{G} \right)^{3/2} \approx 2.3 Y_e^2 m_* \approx 4.3 Y_e^2 M_\odot, \quad (3.263)$$

and we assume that we are in the fully degenerate case: in the integration over momenta of the phase space distribution we insert a cutoff at the Fermi energy.

We find:

$$P = \frac{4\pi}{3h^3} g_* \int_0^{p_F} dp p^2 \frac{p^2 c^2}{\epsilon_p}, \quad (3.264)$$

where $\epsilon_p = (p^2 c^2 + m^2 c^4)^{1/2}$. We integrate with the adimensional variable $x = p/(m_e c)$: we get

$$P = \frac{8\pi}{3h^3} m_e^4 c^5 \int_0^{x_F} \frac{x^4}{(1+x^2)^{1/2}} dx, \quad (3.265)$$

so we get $P = k_{\text{UR}} n_e^{4/3} I(x_F)$, where we incorporated the integral in the term $I(x_F)$: this is given by

$$I(x) = \frac{3}{2x^4} \left(x(1+x^2)^{1/2} \left(\frac{2x^2}{3} - 1 \right) + \log \left(x + (1+x^2)^{1/2} \right) \right), \quad (3.266)$$

and

$$x_F = \frac{p_F}{m_e c} = \left(\frac{3n_e}{8\pi} \right)^{1/3} \frac{h}{m_e c} = \left(\frac{3Y_e \rho_c}{8\pi m_H} \right)^{1/3} \frac{h}{m_e c}, \quad (3.267)$$

so if $x_F \gg 1$ we have $I(x_F) \sim 1$, the ultrarelativistic case, while if $x_F \ll 1$ we have $I(x_F) \sim 4x_F/5$. This then interpolates between our different cases.

$$k_{\text{UR}} \left(\frac{Y_e \rho_c}{m_H} \right)^{4/3} I(x_F) \approx \left(\frac{\pi}{36} \right)^{1/3} G M^{4/3} \rho_c^{4/3}, \quad (3.268)$$

so we can extract the mass:

$$M = I(x_F)^{3/2} M_{\text{CH}}, \quad (3.269)$$

and the important thing is that $x_F \propto n_e^{1/3} \propto \rho_c^{1/3}$.

[Graph: on the x axis M/M_{CHANDRA} , on the y axis ρ_c .]

If we increase the mass, the star is not able to support itself by the pressure due to being a gas of degenerate electrons.

This gives us

$$M_{\text{CHANDRA}} = 3.1 Y_e^2 m_* = 5.8 Y_e^2 M_{\odot} = 1.4 M_{\odot}. \quad (3.270)$$

It can be shown that the mean density is around

$$\langle \rho \rangle = \frac{1}{6} \rho_c = \frac{0.51}{Y_e^2} \left(\frac{M}{m_*} \right)^2 \frac{m_H}{(h/m_e c)^3}, \quad (3.271)$$

so we can estimate the radius as

$$R = \left(\frac{3M}{4\pi \langle \rho \rangle} \right)^{1/3} \approx 0.77 Y_e^{5/3} \left(\frac{M}{m_*} \right)^{1/3} \alpha_G^{-1/2} \frac{h}{m_e c}, \quad (3.272)$$

and the object at the end is

$$\alpha_G^{-1/2} \frac{h}{m_e c} \approx 3 \times 10^7 \text{ m}, \quad (3.273)$$

so if we take $Y_e = 0.5$ we find:

$$R = \frac{R_{\odot}}{74} \left(\frac{M_{\odot}}{M} \right)^{1/3}. \quad (3.274)$$

The luminosity is given by

$$L = 4\pi R^2 \sigma T_E^4 = \frac{1}{74^2} \left(\frac{M_\odot}{M} \right)^{4/3} \left(\frac{T_E}{6000 \text{ K}} \right) L_\odot, \quad (3.275)$$

so if we take a typical effective temperature of around 10^4 K (recall that neutron stars are in the blue part of the HR diagram), $M = 0.4M_\odot$ we get $L \approx 3 \times 10^{-3} L_\odot$.

We deal with degenerate stars: the Pauli exclusion principle plays a critical role, and the process

$$n \rightarrow p + e^- + \bar{\nu}_e \quad (3.276)$$

is inhibited, since there is no more room for electrons. On the other hand, the process

$$e^- + p \rightarrow n + \nu_e \quad (3.277)$$

is favoured. We can look at the Saha formula to get numerical estimates for this. The chemical potential of neutrinos can be neglected, therefore we find

$$\mu_n = \mu_p + \mu_e, \quad (3.278)$$

and the same equation holds for their Fermi energies.

This means that at a certain point we will have only neutrons.

We need to exploit the Fermi exclusion principle. Why does the equation

$$\epsilon_{F,n} = \epsilon_{F,p} + \epsilon_{F,e}, \quad (3.279)$$

favour neutrons? we have the constraint that the number of protons must equal the number of electrons, but the Fermi energy depends on the number density of these. typical numbers then become

$$n_p = n_e = \frac{n_n}{200}. \quad (3.280)$$

We will have

$$n_n = Y_n \frac{\rho_c}{m_n} \approx \frac{\rho_c}{m_n} \quad \text{since} \quad Y_n \approx 1. \quad (3.281)$$

For a neutron star we have

$$\rho_c \approx 3.1 \left(\frac{M}{M_*} \right)^2 \frac{m_n}{(h/m_n c)^3}, \quad (3.282)$$

while for a white dwarf we have

$$\rho_c \approx \frac{3.1}{Y_e^5} \left(\frac{M}{M_*} \right)^2 \frac{m_H}{(h/m_e c^2)^3}, \quad (3.283)$$

the mass is given by $M_* = \alpha_G^{-3/5} m_n \approx 1.85 M_\odot$ and for the radius we get

$$R = 0.77 \left(\frac{M_*}{M} \right)^{1/3} \alpha_G^{-1/2} \frac{h}{m_n c}, \quad (3.284)$$

where the characteristic length is given by

$$L_n = \alpha_G^{-1/2} \frac{h}{m_n c} \approx 17 \text{ km} \approx \frac{1}{1200} L_e, \quad (3.285)$$

1200 times smaller than the corresponding length scale for electrons: L_n is the characteristic length scale for a NS, L_e is the characteristic length for a white dwarf. The maximum mass is $M_{\text{max}}^{\text{NS}} = 3.1 M_* = 5.8 M_\odot$. This is the mass of a star *remnant*, which encompasses mass from the core only: the initial star will be much larger.

Can this NS become a BH? we need to compute

$$\frac{GM}{Rc^2} \approx 0.2 \left(\frac{M}{M_*} \right)^{4/3}, \quad (3.286)$$

so if the mass is large enough we can reach the critical value of $GM/Rc^2 = 2$.

Neutron stars were first detected as rotating objects: *pulsars*. We have

$$\frac{GM}{R^2} \approx R \omega_{\text{max}}^2, \quad (3.287)$$

the maximum angular velocity which can be supported gravitationally: it comes out to be

$$\tau_{\text{min}} = \frac{2\pi}{\omega_{\text{max}}} \approx 2\pi \left(\frac{R^3}{GM} \right)^{1/2}, \quad (3.288)$$

which is of the order

$$\tau_{\text{min}} \approx 11 \left(\frac{M_*}{M} \right) \alpha_G^{-1/2} \frac{h}{m_n c^2} \approx 0.6 \frac{M_*}{M} \text{ ms}. \quad (3.289)$$

This gives us a bound for gravitational waves of astrophysical origin.

no comment on the GR corrections to this formula: however I'd expect them to be significant

We move on to *the GR issue*.

In a star, the classical equation of hydrostatic balance is:

$$\frac{dP}{dr} = - \frac{Gm\rho}{r^2}, \quad (3.290)$$

while the GR equations for this are the TOV equation: Tolman-Oppenheimer-Volkov:

$$\frac{dP}{dr} = - \frac{Gm\rho}{r^2} \left(1 + \frac{P}{\rho c^2} \right) \left(1 + \frac{4\pi r^3 P}{mc^2} \right) \left(1 - \frac{2Gm}{rc^2} \right)^{-1}, \quad (3.291)$$

and the first correction is reminiscent of cosmology: *the pressure itself contributes to the inertia of the system.*

If we have constant density, in the Newtonian case we get

$$m(r) = \frac{4\pi}{3}\rho_0 r^3, \quad (3.292)$$

so than we can integrate and get

$$P(r) = \frac{2\pi G}{3}\rho_0^2(R^2 - r_0^2), \quad (3.293)$$

where we inserted the boundary condition $P(R) = 0$. In the first-order GR case, this can still be solved analytically! We get

$$P(r) = \rho_0 c^2 \left(\frac{(1 - 2GMr^2/R^3c^2)^{1/2} - (1 - 2GM/Rc^2)^{1/2}}{3(1 - 2GM/Rc^2)^{1/2} - (1 - 2GMr^2/R^3c^2)^{1/2}} \right), \quad (3.294)$$

Check exponents

so we get

$$P_c = \frac{2\pi}{3}G\rho_0^2 R^2 = \left(\frac{\pi}{6}\right)^{1/3} GM^{2/3} \rho_c^{4/3}, \quad (3.295)$$

so we can look at what happens when we consider $r = 0$: we get

$$P_c = \rho_0 c^2 \left(\frac{1 - \sqrt{1 - 2GM/Rc^2}}{3\sqrt{1 - 2GM/Rc^2} - 1} \right), \quad (3.296)$$

so we can see that the central pressure is finite as long as

$$\frac{GM}{Rc^2} < \frac{4}{9}, \quad (3.297)$$

which is not 1/2 since we have made some approximations.

Then we have a bound

$$M_{\max} \approx \left(\frac{8\pi f}{9}\right)^{3/2} M_*, \quad (3.298)$$

with $f \sim 1$. This means than the objects becomes contained inside its Schwarzschild radius, $R = 2GM/c^2$.

Tomorrow we will speak of galaxy formation.

Fri Dec 20 2019

The lectures in January will be about *cosmological* gravitational waves mainly.

Write to him if you do not have access to the Dropbox.

On the exam: a traditional oral exam, with questions on the main topics dealt with in class. The days on the calendar do not mean anything. The exams should be agreed upon by email.

Now we talk about the formation of dark matter halos. This is in Sabino's notes in the Dropbox.

We say we have a spherical object in the universe, focus on it and apply Birkoff's theorem: we can study it independently of the surroundings.

We can consider halos of different densities to account for over and under densities. In the real world, we will not have spheres, but three-axial ellipsoids: it is known that if we have over-densities the non-sphericity will increase, if we have under-densities it will decrease.

However we treat spherical models because it is simple. Historically the Americans supported the spherical models, the Russians supported the "pancake model". Now we know that we have pancakes with spheres inside (?).

We consider a sphere. We define:

$$\delta(\vec{x}, t) = \frac{\rho(\vec{x}, t) - \bar{\rho}(t)}{\bar{\rho}(t)}, \quad (3.299)$$

and we assume $0 < \delta \ll 1$ (a *small over-density*), although in general we could have $-1 < \delta < +\infty$.

We found in previous lessons that we have a growing mode $\delta \propto t^{2/3}$ and a decaying mode $\delta \propto t^{-1}$, for which we had $v \propto t^{1/3}$ and $v \propto t^{-4/3}$ respectively.

So we choose a time t_i such that

$$\delta(t) = \delta_+(t_i) \left(\frac{t}{t_i} \right)^{2/3} + \delta_-(t_i) \left(\frac{t}{t_i} \right)^{-1}. \quad (3.300)$$

The linearized continuity equation was:

$$v = i \frac{\dot{\delta}}{k} a \propto \left(\frac{2}{3} \delta_+(t_i) \left(\frac{t}{t_i} \right)^{1/3} - \delta_-(t_i) \left(\frac{t}{t_i} \right)^{-4/3} \right), \quad (3.301)$$

since $a \propto t^{2/3}$. This is in order to write explicitly the fact that we need two initial conditions.

We suppose that at $t = t_i$ we have *unperturbed Hubble flow*: $v(t_i) = 0$, which means

$$\delta_-(t_i) = \frac{2}{3} \delta_+(t_i), \quad (3.302)$$

and we can have a generic initial density, which we can express as:

$$\delta(t_i) = \delta_i = \delta_+ + \delta_- = \frac{5}{3} \delta_+. \quad (3.303)$$

Our perturbation can be dealt with cosmologically as being a *local FRLW closed universe*. Say our sphere has a radius R : then if we have $\delta = \delta_i$ then

$$\Omega(t_i) = 1 + \delta_i. \quad (3.304)$$

We have then, from the Friedmann equations:

$$\dot{a}^2 = \frac{8\pi G}{3}\rho a^2 - k, \quad (3.305)$$

where $k = +1$ if $\delta_i > 0$, because of the fact that our sphere is locally a closed universe. This then means:

$$-k = (1 - \Omega)a^2 H^2, \quad (3.306)$$

so

$$\frac{\dot{a}^2}{a_i^2} = H_i^2 \left(\Omega_p(t_i) \frac{a_i}{a} + (1 - \Omega_p(t_i)) \right), \quad (3.307)$$

where the index p denotes the fact that we are talking about the perturbation. We have:

$$\rho_p(t) = \rho_p(t_i) \left(\frac{a_{pi}}{a_p} \right)^3 \quad (3.308a)$$

$$= \rho_c(t_i) \Omega_p(t_i) \left(\frac{a_{pi}}{a} \right)^3. \quad (3.308b)$$

We want to derive a time for the *turnaround time* t_m :

$$\rho_p(t_m) = \rho_c(t_i) \Omega_p(t_i) \left(\frac{\Omega_p(t_i) - 1}{\Omega_p(t_i)} \right)^3 \quad (3.309a)$$

$$= \rho_c(t_i) \frac{(\Omega_p(t_i) - 1)^3}{\Omega_p(t_i)^2}, \quad (3.309b)$$

since at $t = t_m$ we have $\frac{\dot{a}^2}{a_i^2} = 0$.

since $\dot{a}(t_m) = 0$?

Some time ago we had found:

$$t(\theta_m = \pi) = t_m = \frac{\pi}{2H_0} \frac{\Omega_0}{(\Omega_0 - 1)^{3/2}}, \quad (3.310)$$

but calculating this *now* is arbitrary: the same formula holds replacing 0 with i .

Then we have

$$t_m = \frac{\pi}{2H_i} \left(\frac{\rho_c(t_i)}{\rho_p(t_m)} \right)^{3/2}, \quad (3.311)$$

since we have found that in the relation from a few lessons ago we have exactly the inverse of the relation we just derived connecting Ω and ρ .

It is convenient to consider unperturbed Hubble flow so that the Hubble parameter inside and outside is the same.

Then we have

$$H^2(t_i) = \frac{8\pi G}{3} \rho_c(t_i), \quad (3.312)$$

so we have a cancellation: we find

$$t_m = \frac{\pi}{2H_i} \left(\frac{\rho_c(t_i)}{\rho_p(t_i)} \right)^{1/2} = \left(\frac{3\pi}{32G\rho_p(t_m)} \right)^{1/2}, \quad (3.313)$$

so we get

$$\rho_p(t_m) = \frac{3\pi}{32Gt_m^2}, \quad (3.314)$$

which holds inside the sphere, while the critical density *outside* is

$$\rho_c(t_m) = \frac{1}{6\pi Gt_m^2}. \quad (3.315)$$

We are implicitly using the *synchronous gauge*, and we are finding an exact solution to the EFE: The Lemaître-Tolman-Bondi solution.

We can ask: at a certain given time, how much is the interior density larger than the exterior one? this is given by

$$1 + \delta_p(t_m) = \chi(t_m) = \frac{\rho_p(t_m)}{\rho_c(t_m)} = \frac{3\pi}{32G} 6\pi G = \left(\frac{3\pi}{4} \right)^2 \approx 5.6. \quad (3.316)$$

Then we have $\delta_p(t_m) \approx 4.6$. How does this compare to linear theory?

$$\delta_p(t_m) = \delta_t(t_i) \left(\frac{t_m}{t_i} \right)^{2/3} + \delta_-(t_i) \left(\frac{t_m}{t_i} \right)^{-1} \approx \frac{3}{5} \delta_p(t_i) \left(\frac{t_m}{t_i} \right)^{2/3}, \quad (3.317)$$

since the second term vanishes.

We know that $H_i = 2/(3t_i)$, so

$$t_m = \frac{3\pi t_i}{4} \times \left(\frac{1 + \delta_i}{\delta_i^{3/2}} \right) \approx \frac{3\pi t_i}{4} \delta_i^{-3/2}, \quad (3.318)$$

which comes from the equation from the perturbation density at t_m , in which we substitute $\Omega_p - 1 = (\delta_p + 1) - 1 = \delta_p$, and we approximate $1 + \delta \approx 1$ since δ is small.

So we have

$$\frac{t_m}{t_i} = \frac{3\pi}{4} \delta_i^{-3/2}, \quad (3.319)$$

which gives us a correction

$$\delta_p(t_m) = \frac{3}{5} \delta_i \left(\frac{3\pi}{4} \delta_i^{-3/2} \right)^{2/3} = \frac{3}{5} \left(\frac{3\pi}{4} \right)^{2/3} \approx 1.87. \quad (3.320)$$

So linear theory would have given us a wrong answer, as we should expect: we are very far from the $\delta \ll 1$ regime in which we expect linearization to hold.

We speak of the virialization time (the time after which the VT starts applying?). It is hard to calculate, different books give different values.

$$t_{\text{vir}} \approx t_c = 2t_m. \quad (3.321)$$

Then we have

$$E_{\text{tot}} = T + E_{\text{gr}} = \frac{1}{2} E_{\text{gr}} = -T, \quad (3.322)$$

since $2T + E_{\text{gr}} = 0$.

The total energy at virialization is given by

$$E_{\text{eq}} = -\frac{1}{2} \frac{3}{5} \frac{GM^2}{R_{\text{eq}}}, \quad (3.323)$$

where the factor $3/5$ is given by the geometry of the system, we are assuming constant density for our sphere, while the $1/2$ comes from the VT. At the time of collapse instead the energy is given by

$$E_m = -\frac{3}{5} \frac{GM^2}{R_m}, \quad (3.324)$$

This then means that the radius at virialization, R_{eq} , is twice the radius at the start of the collapse, R_m :

$$2R_{\text{eq}} = R_m. \quad (3.325)$$

This also means, by mass conservation, that $\rho_{\text{vir}} = 8\rho_m$ since the volume shrinks by 2^3 .

We want to know the density at collapse::

$$\frac{\rho_p(t_c)}{\rho_c(t_c)} = \underbrace{\frac{\rho_p(t_c)}{\rho_p(t_m)}}_8 \underbrace{\frac{\rho_p(t_m)}{\rho_c(t_m)}}_{\chi \approx 5.6} \underbrace{\frac{\rho_c(t_m)}{\rho_c(t_c)}}_{2^2} \approx 180. \quad (3.326)$$

This object is called $1 + \delta(t_c)$, so $\delta(t_c) \approx 179$.

What would happen if we were to use linear theory at this time? Then we would find

$$\delta_+(t_c) = \delta_+(t_m) \left(\frac{t_c}{t_m} \right)^{2/3}, \quad (3.327)$$

which gives us

$$\delta_+(t_c) \approx \frac{3}{5} \left(\frac{3\pi}{4} \right)^{2/3} 2^{2/3} \approx 1.686, \quad (3.328)$$

which is a kind of “clock”: how long can we use linear theory for?

We deal with *** theory: we introduce

$$n(M) = \frac{dN}{dM} = \# \text{ of objects per unit volume with mass in } [M, M + dM]. \quad (3.329)$$

Let us consider linear perturbations $\delta(\vec{x}, t)$: perturbations in the matter density dealt with using linear theory only.

These tend to fluctuate a lot. We need a *filter*: we use a *low-pass* filter, ignoring the high-frequency modes. It is $W_R(\vec{x})$: R is a spatial radius, we ignore modes with spatial frequency larger than $1/R$.

We use a filter labelled by a mass M , which we find by assuming a certain density, and then using $M \propto R^3$.

It is generally a good idea to assume gaussianity. In the Planck data it was found by Sabino’s team that the bounds on non-gaussianity are very low, the data are almost gaussian. So, we have

$$p(\delta_M) d\delta_M = \frac{1}{\sqrt{2\pi\sigma_M^2}} \exp\left(-\frac{\delta_M^2}{2\sigma_M^2}\right) d\delta_M, \quad (3.330)$$

where the variance is typically diverging if we do not apply the filter: we have

$$\sigma_M^2 = \langle \delta_M^2 \rangle \propto M^{-2\alpha}, \quad (3.331)$$

and typically $\alpha \sim 1/2$.

We define a threshold for the value of δ_M , and we want to compute the probability of the value becoming larger than it. We use for the critical value $\rho_c = 1.686$ from before. We have

$$\mathbb{P}_{>\delta_c}(M) = \int_{\delta_c}^{\infty} d\delta_M p(\delta_M). \quad (3.332)$$

so we have

$$n(M)M dM = \rho_m (\mathbb{P}_{>\delta_c}(M) - \mathbb{P}_{>\delta_c}(M + dM)) \quad (3.333a)$$

$$= \rho_m \left| \frac{d\mathbb{P}_{>\delta_c}}{dM} \right| dM \quad (3.333b)$$

$$= \rho_m \left| \frac{d\mathbb{P}_{>\delta_c}}{d\sigma_M} \right| \left| \frac{d\sigma_M}{dM} \right| dM. \quad (3.333c)$$

Integrating $\frac{d\mathbb{P}_{>\delta_c}}{dM}$ we expect to find the matter density again, but we find 1/2 of it.

This comes from a miscount: as the mass we are considering shrinks, we might be already including smaller objects inside the gravitational influence of larger ones. Properly accounting for this one gets precisely a factor 2.

Integrating we get

$$n(M) = \frac{2}{\sqrt{\pi}} \frac{\rho_m}{M_*^2} \alpha \left(\frac{M}{M_*} \right)^\alpha \exp \left(- \left(\frac{M}{M_*} \right)^{2\alpha} \right), \quad (3.334)$$

where $M_* = (2/\delta_c)^{1/2\alpha} M_0$.

Accounting for non spherical collapse we get much better estimates.

Wed Jan 08 2020

3.1 Gravitational waves and interferometry

Guest lecture by Angelo Ricciardone.

An outline:

1. introduction about gravitational waves;
2. frequency bands \iff sources of GW \iff detectors;
3. GWs and observables;
4. stochastic background of GW: characterization, sources, detection (here we discuss *cosmological* sources).

References:

1. Book: “Gravitational Waves - theory and experiments” by Michele Maggiore (Oxford University Press, 2007);
2. Book: “Gravitational Waves - Astrophysical sources” (2018);
3. Paper: “The basics of gravitational wave theory”, F. Flanagan (<https://arxiv.org/abs/gr-qc/0501041>);
4. Review: “Gravitational waves from inflation”, (<https://arxiv.org/abs/1605.01615>);
5. Review: “Cosmological background of gravitational waves”, (<https://arxiv.org/abs/1801.04268>);
6. Kind of unrelated review: “Inflation and the Theory of Cosmological Perturbations” <https://arxiv.org/abs/hep-ph/0210162>.

We start with some general facts.

Gravitational waves appear naturally in GR: they are propagating oscillations of the gravitational field.

The gravitational interaction is *weak*: this implies that GWs travel freely, but they are also hard to detect.

The frequencies we see for the EM spectrum are in the range $f_{EM} \sim 10^4 \text{ Hz} \div 10^{20} \text{ Hz}$, radio waves to γ rays.

The typical frequencies of GWs are, instead, in the range $f_{GW} \sim 10^{-16} \text{ Hz} \div 10^4 \text{ Hz}$, the lower end of this range is the frequency of the CMB GWs while the higher end is given off by astrophysical sources. Ground-based detectors can detect the higher end of this spectrum.

But the CMB is electromagnetic! Is there gravitational radiation corresponding to it?

The wavelength and frequency are typically comparable to the size of the object which is emitting the GWs.

Let us consider astrophysical sources:

$$T = 2\pi\sqrt{\frac{R^3}{GM}} \implies f = \sqrt{\frac{G\rho}{4\pi^2}}, \quad (3.335)$$

and since the mass is related to the density by $M = \rho V = \frac{4}{3}\pi R^3 \rho$, then we have

$$f_{GW} \approx \frac{1}{2\pi} \sqrt{\frac{3GM}{4\pi R^3}}. \quad (3.336)$$

We know that the radius of the object must be greater than the Schwarzschild radius $R_s = 2GM/c^2$: substituting this in we get

$$f_{GW} \approx \frac{1}{4\pi} \frac{c^3}{GM}, \quad (3.337)$$

which means, substituting the numbers, that

$$f_{GW} \approx 10^4 \text{ Hz} \frac{M_\odot}{M}, \quad (3.338)$$

Is the difference between the formulas coming from different geometries of the problem? Like, a rotating objects versus two inspiralling ones?

We make some estimates for neutron stars. Typically they have $M \sim 1.4M_\odot$ and $R \sim 10^4 \text{ m}$. So, we have $f \sim 10^4 \text{ Hz}$.

For small BHs, we have $M \sim 30M_\odot$, so $f_{GW} \sim 300 \text{ Hz}$. This is the band in which LIGO/VIRGO works: these interferometers cannot measure GWs with frequency smaller than 1 Hz.

If we increase the mass, if we use $M \sim 10^7 M_\odot$, we get $f_{GW} \sim 10^{-3} \text{ Hz}$: this is the band in which LISA will work.

There are indirect evidences of GW emission.

Pulsars slow down: this is the Hulse-Taylor binary pulsar, two neutron stars which are rotating around each other emit gravitational waves. The GW emission back-reacts on the dynamics of the binary, on a timescale which is observable.

A reference for this is <https://arxiv.org/abs/astro-ph/0407149>.

We can make a plot: radius of the system vs mass of the system. Since $f \propto M^{1/2} R^{-3/2}$, constant frequency means a powerlaw in this plot, so a straight line in $\log R$ vs $\log M$.

The chirp of the inspiral is *also* a powerlaw. The GR predictions for the period decrease of the Hulse-Taylor binary NS match observations very precisely.

On the 14/09/2015, we had the first direct evidence of GWs with the first detection.

On 17/08/2017 we had the first detection of a NS/NS merger: this was the birth of multimessenger astronomy, since we saw the event with optical telescopes also. This meant that the signal got to us with a speed which is the same as the electromagnetic speed of light.

Now, there is an app: the “GW events app” <https://apps.apple.com/us/app/gravitational-wave-events/id1441897107>.

In GR, gravitational waves can be polarized in two different ways. These are called “plus” and “cross” polarizations.

Can one of these be rotated into the other? There might be some issue since one is a pseudotensor. . .

We distinguish: the High Frequency band goes from 10^4 Hz to 1 Hz; the Low Frequency band goes from 1 Hz to 10^{-4} Hz; the Very Low Frequency band goes from 10^{-7} Hz to 10^{-9} Hz; the Extremely Low frequency band goes from 10^{-15} Hz to 10^{-18} Hz.

There are gaps since in certain frequency regions the methods we have on either side fail for different reasons leaving a gap.

Let us start from HF: it is the domain of Earth-based interferometers: LIGO (Livingstone & Hanford), which has 4 km arms, the two detectors are separated by 3000 km. Also, there is VIRGO near Pisa: an arm is 3 km long.

There is also Geo600 in Hannover, Germany: it has 600 m arms. In Japan there is Kagra: it has 4 km arms.

There are several reasons to have more than one detector: we can identify the position of the source, we can verify signals.

There will be the new Einstein Telescope, in a triangular configuration, in Italy or the Netherlands. There will be the Cosmic Explorer, with 40 km arms.

Now, let us discuss the main sources in the HF band. We have:

1. coalescence of stellar-mass BH binaries and NSs, for these we have an upper bound of $M \lesssim 10^3 M_\odot$;
2. rotation of neutron stars (pulsar);
3. stellar collapse: supernova to BH or NS.

In the LF band, the domain of space-based interferometers, we will have LISA, in a triangular shape with 2.5×10^6 km, and the Japanese DECIGO, with $L \sim 1000$ km.

In the LF band, the sources are:

1. white dwarves merging;
2. NS merging;

3. inspiral and coalescence of SMBH (masses from 100 to $10^8 M_\odot$).

In the VLF band, we can use Pulsar-Time Array. The sources in this frequency range are:

1. GWs from SMBH with $M > 11M_\odot$, but it seems like there are no black holes this large;
2. GWs from cosmic strings & from phase transitions;

In the ELF band, we have cosmological sources:

1. primordial GWs: here we have $h \sim (E_{\text{infl}}/M_P)^2$.

Typically we have amplitudes increasing as the frequency decreases.

For scalar perturbations we have

$$\frac{\Delta T}{T} \propto \delta\phi. \quad (3.339)$$

Vector perturbations decay with the expansion of the universe. We do have tensor perturbations. B mode polarizations correspond to primordial gravitational waves.

What are B modes?

In a simple MM interferometry, $\Delta L \propto h$.

LISA will orbit the Sun at 20° from the Earth. The astrophysical targets for LISA are MBHBs, EMRIs and compact WDs.

Also, there are potential cosmological sources. We have first order phase transitions around the TeV, inflationary GWs, cosmic strings and using MBHBs as standard sirens.

We put powerlaw amplitude spectra in a graph. What is Ω ?

We saw powerlaw spectra with Ω increasing with f : do they not have a UV problem since the energy density increased with the frequency? No, since because of the transfer function between the early universe and now at a certain point the powerlaw becomes decreasing.

Cosmological sources are stochastic: they give a background.

If we will have a detector with high sensitivity, we will see many events: they will form a stochastic background.

Can we not correlate the signals in such a way that we only look at GWs from a specific angular region?

A stochastic background of GW comes from a large number of independent uncorrelated sources that are not individually resolvable.

CSGWB is a candidate source for LISA.

For tomorrow, we can either give an overview of cosmological sources for GWs or we can derive the relation

$$\frac{\Delta L}{L} \propto h, \quad (3.340)$$

Thu Jan 09 2020

Many of these topics will be discussed in more details in the course by Nicola Bartolo on cosmological perturbations and in the course by Giacomo Ciani on gravitational waves.

Today we will focus on the LISA detector, which works for frequencies $f \sim 10^{-5} \div 10^{-1}$ Hz, and on the cosmological stochastic background of GWs (CSGWB).

We can make an analogy with the CMB: we expect to see a background of gravitational radiation coming from all directions.

The stochastic background can have different origins:

1. astrophysical origin: a coherent superposition of a large number of astrophysical sources, which are too weak to be detected separately;
2. cosmological origin: it is generated in the early universe by a variety of mechanisms:
 - (a) amplification of primordial tensor fluctuations via inflation;
 - (b) GWs from phase transition around the TeV scale (this is not the only energy scale at which they can be emitted, but it is the energy we'd need in order to detect them with LISA);
 - (c) GWs from topological defects.

A stochastic background of cosmological origin is expected to be

1. isotropic;
2. stationary;
3. unpolarized, which means that the cross and plus polarizations will have the same amplitude.

Let us discuss the main properties of the frequency spectrum: it is characterized

1. in terms of (normalized) energy density per unit logarithmic interval of frequency: this is called $h_0^2 \Omega_{GW}(f)$;
2. in terms of the spectral density of the ensemble average of the Fourier component of the metric $S_h(f)$;
3. more on the experimental side: in terms of a characteristic amplitude of the stochastic background $h_c(f)$.

Let us define these three quantities. The energy density is given by

$$\Omega_{GW} = \frac{1}{\rho_c} \frac{d\rho_{GW}}{d \log f}, \quad (3.341)$$

where ρ_{GW} is the energy density of SGWB, f is the frequency, while ρ_c is the present value of the critical energy density, defined as

$$\rho_c = \frac{3H_0^2}{8\pi G}, \quad (3.342)$$

where we usually write $H_0 = h_0 \times 100 \text{ km}/(\text{sMpc})$.

So, usually we plot $h_0^2 \Omega_{\text{GW}}$ in order to ignore the uncertainties on the measurements of H_0 .

We write stochastic GW at a given point $\vec{x} = 0$ in the transverse traceless (TT) gauge: $h_{ii} = 0$ and $\partial_i h^{ij}$: we get

$$h_{ab}(t) = \sum_{A=+, \times} \int_{\mathbb{R}} df \int_{S^2} d\Omega \hat{h}_A(f, \Omega) \exp(-2\pi i f t) e_{ab}^A(\hat{\Omega}), \quad (3.343)$$

where we must have the condition $\hat{h}_A(f, \Omega) = \hat{h}_A^*(-f, \Omega)$. Here $\hat{\Omega}$ is a unit vector representing the direction of propagation of the wave, and $d\hat{\Omega} = d\cos(\theta) d\phi$, e_{ab}^A are the polarization tensors:

1. $e_{ab}^+(\hat{\Omega}) = 2\hat{m}_{[a}\hat{m}_{b]}$;
2. $e_{ab}^\times(\hat{\Omega}) = 2\hat{m}_{(a}\hat{m}_{b)}$;

where $\hat{m}_{a,b}$ are unit vectors orthogonal to each other and to the propagation direction.

We have the condition $e_{ab}^A e^{Bab} = 2\delta^{AB}$.

Then, assuming a SGWB which is isotropic, unpolarized and stationary we will have

$$\langle \hat{h}_A^*(f, \hat{\Omega}) \hat{h}_{A'}(f', \hat{\Omega}') \rangle = \delta(f - f') \frac{1}{4\pi} \delta^{(2)}(\hat{\Omega}, \hat{\Omega}') \delta_{AA'} \frac{1}{2} S_h(f), \quad (3.344)$$

where

$$\delta^{(2)}(\hat{\Omega}, \hat{\Omega}') = \delta(\phi - \phi') \delta(\cos(\theta) - \cos(\theta')). \quad (3.345)$$

The factors of $1/4\pi$ and $1/2$ are for convention, for normalization, so that $\int S_h(f) df = 1$. Then the remaining bit, $S_h(f)$, is called the spectral density.

So, using the equations we found so far, we get

$$\langle h_{ab}(t) h^{ab}(t) \rangle = 2 \int_{\mathbb{R}} df S_h(f) \quad (3.346a)$$

$$= 4 \int_{f=0}^{f=\infty} d \log f f S_h(f). \quad (3.346b)$$

We define the characteristic amplitude $h_c(f)$ as

$$\langle h_{ab}(t) h^{ab}(t) \rangle = 2 \int_{f=0}^{f=\infty} d \log f h_c^2(f), \quad (3.347)$$

which means that $h_c^2(f) = 2f S_h(f)$.

The last step is to relate $h_c(f)$ and $h_0^2 \Omega_{\text{GW}}(f)$.

The energy density is defined as:

$$\rho_{\text{GW}} = \frac{1}{32\pi G} \langle \dot{h}_{ab} \dot{h}^{ab} \rangle, \quad (3.348)$$

where the average is performed over a wavelength, but by the ergodic theorem it can also be performed over a period. So,

$$\rho_{GW} = \frac{4}{32\pi G} \int_{f=0}^{f=\infty} d(\log f) f (2\pi f)^2 S_h(f), \quad (3.349)$$

whicm means that

$$\frac{d\rho_{GW}}{d \log f} = \frac{\pi}{4G} f^2 h_c^2(f) \quad (3.350a)$$

$$= \frac{\pi}{2G} f^3 S_h(f), \quad (3.350b)$$

so in the end we have

$$\Omega_{GW}(f) = \frac{2\pi^2}{3H_0^2} f^2 h_c^2(f), \quad (3.351)$$

therefore

$$\Omega_{GW}(f) = \frac{4\pi^2}{3H_0^2} f^3 S_h(f). \quad (3.352)$$

The *strain* is the quantity S_h .

An advantage of GWs is the fact that they decouple right after emission: they maintain the spectral shape.

If we start from the Einstein-Hilbert action

$$S = \int d^4x Fg \frac{M_P^2}{2} R, \quad (3.353)$$

and plug in a tensor-perturbed FRLW metric:

$$ds^2 = -a^2 \left(d\eta^2 + (\delta_{ij} + h_{ij}) d\vec{x}^2 \right), \quad (3.354)$$

we get the equations of motion for gravitational waves. The scenario of slow-roll inflation gives rise to an energy spectrum which is unobservable with LISA as well as ground-based detectors.

If we add an axion to the inflaton Lagrangian:

$$\mathcal{L} = \frac{1}{2} \partial_\mu \phi \partial^\mu \phi + V(\phi) + \underbrace{\frac{\phi}{\Lambda} F_{\mu\nu} \tilde{F}^{\mu\nu}}_{\text{Additional term}}, \quad (3.355)$$

we can see enhanced gravitational waves, which we hope would be detectable with LISA! However the enhancement is very small, we cannot see it.

Chaotic preheating models predict higher amplitudes, however they are at very high frequencies.

We can get observable signals by fine-tuning the parameters, unlikely scenario.

What about phase transition, the collisions of primordial vacuum bubbles? We can get numerical estimates on the spectral shape of this signal, and they seem promising and detectable.

What about cosmic defects? We can have Domain Walls, Cosmic Strings and Cosmic Monopoles.

The prediction here is a flat spectrum at observable frequencies, at possibly observable amplitudes depending on the emission time.

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