**Zeitschrift:** Helvetica Physica Acta
**Band:** 60 (1987)
**Heft:** 1

**Artikel:** Cosmology and particle physics
**Autor:** Straumann, N.
**DOI:** https://doi.org/10.5169/seals-115842

Helvetica Physica Acta
Vol. 60, 1987 9-39

018-0238/87/010009-31$1.50+0.20/0


# COSMOLOGY AND PARTICLE PHYSICS

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## INTRODUCTION

The last ten years have witnessed an increasing interaction between particle physics and cosmology. Two main reasons for the development of this new interdisciplenary field are: (i) It is likely that the properties of the presently observable Universe - its size, age, structure and content - have been strongly influenced by high energy processes in hot early epoches. (ii) The successful electroweak theory has encouraged particle physicists to realize Einstein's dream of unifying more - possibly all - of the fundamental interactions. From what we know, it is clear that a possible basic unity of all forces can only become manifest at very high energies, which may never be reached with man made accelerators [1).].

The large amount of work done so far has at least widened our outlook on cosmology and has opened new avenues. It is, however, difficult to assess what will be of lasting significance. More than once, the original beauty of an idea has been lost. Most of the work remains highly speculative - but interesting, or at least amusing.

Since this talk addresses a general audience of physicists, I will first lead you on a brief tour through the field of cosmology and particle physics, emphasizing some new developments and summarizing the present status of major issues which have been discussed in recent years. In a second part, I will then concentrate on one specific speculative topic, namely cosmic strings and their role in promoting the formation of galaxies.

## I. TOUR D'HORIZON

In talking about cosmology, it is always healthy to keep Fig. 1 in mind, which shows those parts of the Universe that are in principle accessible to us: the backward light cones for the duration of the human race and a small region from our immediate environment. If there is an event horizon, events beyond this horizon can never be observed. To quote G. Ellis [2]: "We are unable to obtain a model of the universe without some specifically cosmological assumptions which are completely unverifiable".

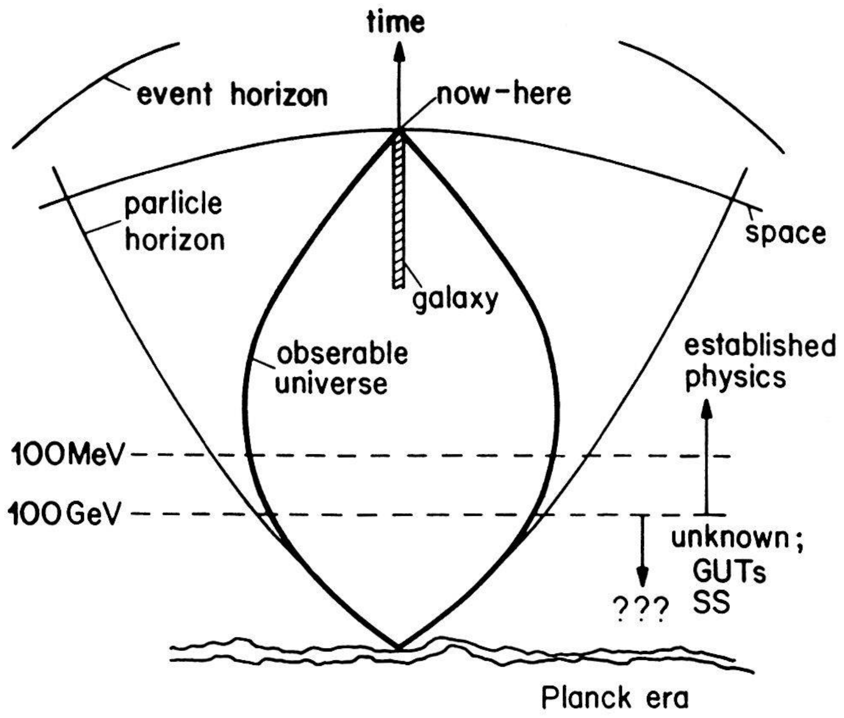


Fig. 1 Spacetime diagram showing the observable parts of the Universe.

## I.1 The Standard Model of Particle Physics

The standard model of particle physics, quantumchromodynamics and electroweak theory, based on the gauge group $SU(3)_c \times SU(2)_L \times U(1)_Y$, has turned out to be incredibly successful[3]. The overall satisfactory situation would have been unthinkable at the beginning of the gauge era about fifteen years ago. No "new physics" beyond the standard model has clearly been established (earlier indications for new physics have not stood the test of increasing statistics.) Further consolidation can be expected from future accelerators (TEVATRON, SLC, LEP, SSC,...). The Higgs boson - the missing link of the standard model - will perhaps then also be found and its nature studied.

It is well-known that many questions cannot be answered in the standard model. The large freedom in parameters calls for a deeper level of understanding, which might provide answers to questions such as:

(1) Why is there this remarkable replication of chiral fermion families and how many exist in Nature ?

(2) What are the properties of neutrinos (masses, lifetimes, mixings) ?

These questions are also, as we shall see, highly significant for astrophysics and cosmology. Conversely, cosmology gives interesting constraints to possible answers of these open problems.

On the particle physics side, an important result has recently been established by the UA1 and UA2 groups at CERN: The total number $N_\nu$ of neutrino flavors with conventional couplings to the $Z^0$ intermediate boson, and with masses small enough for the decay $Z^0 \rightarrow \nu\bar{\nu}$ to proceed without suppression is bounded by [4)]

$$N_\nu < 5.4 \pm 1 . \qquad (1)$$

The limit is obtained by using the theoretical estimate for the ratio of $W$ and $Z$ production in $\bar{p}p$ collisions, and taking the very firm theoretical values for their leptonic decay probabilities. The data for $\bar{p}p \rightarrow W \rightarrow \ell\nu$ and $\bar{p}p \rightarrow Z \rightarrow \ell\bar{\ell}$ then result in an upper limit for the $Z \rightarrow \ell\bar{\ell}$ branching ratio, which in turn yields a limit on $N_\nu$ . [An upper limit $N_\nu < 14$ from $e^+e^- \rightarrow \gamma$ + invisibles has recently been obtained [5)].] At SLC and LEP one additional neutrino beyond the three we know could be detected.

Another important very recent result is the new upper limit for the $\nu_e$-mass, established by the group of Kündig [6)] at the University of Zürich:

$$m_{\nu_e} < 18 \text{ eV} . \qquad (2)$$

This limit has been obtained from the endpoint region of the tritium β-spectrum. It includes instrumental and statistical uncertainties as well as uncertainties due to energy loss in the source and the final electron states. (See the report by R.E. Pixley at this meeting.) The upper limit (2) is in contradiction with the results of the ITEP group that have been published in a series of papers [7] and is compatible with all we know from neutrino-less 2β-decays.

Let me give also the best present limits for the $\nu$-masses of the other flavors [8]:

$$m_{\nu_\mu} < 270 \text{ keV} \quad \text{(SIN)},$$
$$m_{\nu_\tau} < 56 \text{ MeV} \quad \text{(PETRA)}\ . \tag{3}$$

I mention also that no oscillations have thus far be detected.[9] Mikheyev and Smirnov [10] have recently pointed out that the modifications of neutrino oscillations in matter [11] could explain the solar-neutrino puzzle in a very elegant way. If $m_2^2 - m_1^2 \simeq 6\text{x}10^{-5}\ \text{eV}^2$ for the masses of the mixing neutrino states, then the high energy neutrinos $\nu_e$ from $^8$B-decay would be converted into $\nu_\mu$ when traversing the sun, and therefore not be detected in the Davis experiment [12]. This mechanism works even for very small neutrino vacuum mixing angles.

Another issue not settled within the standard model, is the strong CP-problem. Quantumchromodynamics has a chiral anomaly which implies that CP is not a natural symmetry of the strong interactions. If the axion solution proposed by Peccei and Quinn is realized in Nature, then axions would be interesting candidates for cold dark matter (see section I.4).

## I.2 Cosmological Facts

There are only a few observational facts of cosmological significance which have been established with some degree of certainty. They all fit into the standard hot big bang model, whose underlying spacetime is a Friedman manifold with scale factor a(t) (t: cosmic time).

### A. Age determinations

Two important parameters characterizing the present Universe are the Hubble parameter $H_o = (\dot{a}/a)_{pres}$ and the density parameter $\rho_o$. (The subscript nought always refers to the present time.) It is customary to use dimensionless parameters, $h_o$ and $\Omega_o$, defined by

$$h_o = H_o/100 \text{ km s}^{-1}\text{Mpc}^{-1}, \quad \Omega_o = \rho_o/\rho_c, \tag{4}$$

where $\rho_c$ is the critical density, determined by $H_o$ and the gravitational constant $G$ *):

$$\rho_c = \frac{3H_o^2}{8\pi G} = 1.9\times10^{-29}\, h_o^2 \;(\text{g cm}^{-3}) . \tag{5}$$

If $\Omega_o < 1$ the Friedman universe expands forever; for $\Omega_o > 1$ it will recollapse in a big crunch. For the intermediate case $\Omega_o = 1$ the present age of the Friedman universe (with vanishing cosmological constant) is $t_o = 2/3H_o$.

$h_o$ is known within a factor of two [13]:

$$\tfrac{1}{2} \lesssim h_o \lesssim 1 . \tag{6}$$

It is remarkable that the Hubble age $H_o^{-1}$, the age of galactic globular clusters, and the age of the galaxy deduced from cosmochronometers all agree within existing uncertainties:

$$\begin{aligned} H_o^{-1} &\simeq \text{age of elements} \simeq \text{age of oldest stars} \\ &\simeq (1\text{-}2) \times 10^{10} \text{ yr}. \end{aligned} \tag{7}$$

It is probably too early to deduce a value for the cosmological constant from apparent discrepancies between the different ages.

### B. The density parameter

The density parameter $\Omega_o$ has been estimated in various ways.

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*) We use units where $\hbar = c = k_B = 1$. In these units the Planck mass, length and time are respectively:

$$m_{Pl} = G^{-\frac{1}{2}} = 1.22\times10^{19} \text{ GeV} = 2.2\times10^{-5} \text{ g},$$
$$\ell_{Pl} = G^{\frac{1}{2}} = 1.6\times10^{-33} \text{ cm},$$
$$t_{Pl} = G^{\frac{1}{2}} = 0.54\times10^{-43} \text{ s}.$$

(i) Infall to Virgo

Measurements of the peculiar local velocity field (deviation from the Hubble flow) induced by the overdensity of the Virgo complex (i.e. by the Virgo cluster and its extended halo) allow one to deduce a value for $\Omega_0$ under the assumption that the total mass is clustered like galaxies on scales $\sim 10\ h_0^{-1}$ Mpc. The result is [14]

$$\Omega_0 \text{ (Virgo-infall)} \simeq 0.1 - 0.2\ . \tag{8}$$

(ii) Mass-to-light ratios

The average luminosity density $\mathcal{L}$ has been determined to be [15]

$$\mathcal{L} = 1.8\text{x}10^8\ h_0\ L_\odot\ \text{Mpc}^{-3} \tag{9}$$

(with an uncertainty of perhaps a factor 2), implying

$$\Omega_0 = 0.6\text{x}10^{-3}\ h_0^{-1} \langle M/L \rangle_{solar}\ , \tag{10}$$

where $\langle M/L \rangle_{solar}$ is the average mass-to-light ratio.
Thus

$$\langle M/L \rangle_{solar} = 1500\ \Omega_0\ h_0\ (M_\odot / L_\odot)\ . \tag{11}$$

From the dynamics of galaxy pairs and clusters of galaxies average values for M/L have been determined. They lead to [15]

$$\Omega_0 \text{ (pairs and clusters)} \simeq 0.1 - 0.2\ . \tag{12}$$

(iii) $\Omega_B$ from big bang nucleosynthesis

The successful predictions of the abundance of light elements $^2D$, $^3He$, $^4He$ and $^7Li$ in the standard model (see subsection D) allows one to deduce [15] the following value of the baryonic part $\Omega_B$ of $\Omega_0$ :

$$\Omega_B \simeq (0.1 \pm 0.06)\ (h_0/0.5)^2\ . \tag{13}$$

(iv) Distribution of infrared galaxies

This is a new method. The surface distribution of infrared galaxies (with $z \lesssim 0.03$), determined by the satellite IRAS, shows a dipole axis, which agrees within the errors with the one inferred from the dipole anisotropy of the 3K-background. This dipole part determines the peculiar local acceleration and allows one - as for the Virgo infall - to deduce $\Omega_0$ on larger scales ($h_0^{-1}$ 100 Mpc). At the 2nd ESO/CERN Symposium a few weeks ago, R. Robinson gave the result [16]

$$\Omega_0 = 1 \pm 0.15 \, , \tag{14}$$

assuming that the infrared galaxies trace the matter distribution. This is a very important result, in particular in view of the prediction $\Omega_0 = 1.00....$ in inflationary models. If true, it would be a strong indication that not all matter is baryonic (see eq.(13)).

I recall that the luminous contribution to $\Omega_0$ is

$$\Omega_{lum} \sim 0.01$$

and thus at least 90 % (perhaps 99 %) of matter is dark.

The nature of this dark matter is a major problem of cosmology. In this connection, it may be useful to recall the contribution of one massive neutrino type to $\Omega_0$ . In the standard model, the present number density is

$$n_\nu + n_{\bar{\nu}} = \frac{3}{11} n_\gamma \simeq 109 \text{ cm}^{-3} \, . \tag{15}$$

Since $\rho_c = 10'540\ h_0^2$ (eV cm$^{-3}$), we conclude

$$\Omega_\nu \simeq \frac{m_\nu \text{ (eV)}}{100\ h_0^2} \, . \tag{16}$$

This gives the following restrictive cosmological bound

$$\sum_{\nu,\bar{\nu}} m_\nu \leq 100\ h_0^2\ \Omega_0 \text{ (eV)} \, . \tag{17}$$

A conservative limit for the right-hand side is 200 eV ($\Omega_0 \lesssim 2$ from limits for the deceleration parameter and the age of the Universe). The new limit (2) gives $\Omega_{\nu_e} < 1$ , but for the other $\nu$ -flavors the cosmological limit is much more stringent than the experimental bounds (3) .

C. The 3K-background radiation

Until recently there were some indications that the observed microwave background radiation might disagree with the blackbody spectrum in the peak region of the spectrum. These distortions have now disappeared [17]. Possible deviations from a thermal spectrum are less than 6 %. (The measurements in the peak region are exceedingly hard and experimenters have always warned theoreticians - often without success - to take the earlier deviations not too seriously.)

Apart from the dipole anisotropy, due to our motion through the background radiation, no intrinsic anisotropies have been observed; $\Delta T/T$ is constrained to be less than $10^{-4}$ on all angular scales above a few arcminutes. These cosmological precision experiments are severly constraining models of galaxy formation. A baryonic matter fluctuation of mass $M$ should induce a fluctuation in the background radiation at some level on an angular scale

$$\theta = 5' \left(h_0 \Omega_0{}^2\right)^{1/3} \left(\frac{M}{10^{14} M_\odot}\right)^{1/3} . \tag{18}$$

We shall return to this in section I.4.

### D. Big bang nucleosynthesis of $^4$He and other light elements and the number of neutrino flavors

Big bang nucleosynthesis occurred at a temperature $T \simeq 0.1$ MeV during the "lepton era". The physics during this era is very simple and so well-understood (apart from possible exotic particles) that the resulting abundance of the light elements $^2$D, $^3$He, $^4$He, and $^7$Li can be calculated reliably [18]. The really difficult and complicated task is to measure present abundances and infer from them the primordial ones, to be compared with the theoretical predictions. Since various complicated additions and depletions occurred at later times (stellar processes, spallation by cosmic rays, etc.), there remain considerable uncertainties for the values for the primordial abundances. Within them, one obtains good agreement with theory for all four elements [19], with one value $\Omega_B$ for the present baryon density parameter [*]. This gives us confidence that the Friedman-Lemaitre models can be further extrapolated back in time. For three $\nu$-species one finds the value for $\Omega_B$ given in eq.(13).

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[*] see next page.

From Fig. 2 it can be seen that the primordial abundance of $^4$He also depends on the number of neutrino flavors. Qualitatively, this is clear because additional species of light particles lead to an increase in the expansion rate and thus to an increase in the $n/p$ ratio (the freeze-out temperature for this ratio is increased). From the observations of D, $^3$He and $^7$Li one deduces $\eta = (1-10) \times 10^{-10}$ and thus the "observational" limit for the primordial abundance by mass of $^4$He, $Y_{prim} \leq 0.25$, leads to the remarkable restriction

$$N_\nu \leq 4 \; . \qquad (19')$$

This tight limit has recently been challenged by J. Ellis et al. [20] on the ground of the following arguments. New results for the beta decay of polarized neutrons imply a relatively low value of the neutron half-life $\tau_{\frac{1}{2}}$ and hence the authors impose the more conservative value $\tau_{\frac{1}{2}} > 10.2$ min. From the astronomical data, they conclude that the bound $Y_{prim} \leq 0.25$ is not really save and they adopt $Y_{prim} \leq 0.26$ as a reasonable bound. From Fig. 2 one sees that these relaxed conditions would allow five neutrino flavors if $\eta$ should be close to $10^{-10}$. Now Yang et al. [18] have adopted the argument that all primordial D which is burnt in stars creates $^3$He, a fraction of which survives stellar processing. Using studies of meteorites and the solar wind, they concluded that $[N(D+{}^3He)/N(H)]_{prim} < 10^{-4}$ . If this condition is imposed, the theoretical correlation in Fig. 3 shows that the limit (19') is quite save. Ellis et al. have weakened the latter condition on the basis of new data of $^3$He in galactic HII regions. These data have, however, apparently been withdrawn. [21].

At any rate, it is perhaps wiser to give a more conservative cosmological limit

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*) $\Omega_B$ is related to the baryon-photon number density ratio $\eta = n_B/n_\gamma$ by

$$\Omega_B = 3.53 \times 10^{-3} \, h_0^{-2} \, (T_0/2.7K)^3 \quad (\eta/10^{10}) \; .$$

$$N_\nu < 5 - 6 \,, \qquad (19'')$$

in order to be on the safe side. Even such a weaker bound is quite impressive.

It is very important that the limits (1) and (19) are so close together. However, it must be emphasized in this context that $N_\nu$ in (1) and (19) is not necessarily the same number. In the bound (19) also other possible light particles which have been relativistic at $T \sim 1$ MeV, are counted, even if they would interact only superweakly. On the other hand, very massive neutrinos ($\sim 1$ GeV) are not included in (19), but are counted in the CERN-experiments, provided they couple with full-strength.

The limits (1) and (19) provide important restrictions for phenomenological models motivated by superstring theory, which contain additional neutral particles in each fermion generation [22] (notably right-handed neutrinos which may be light).

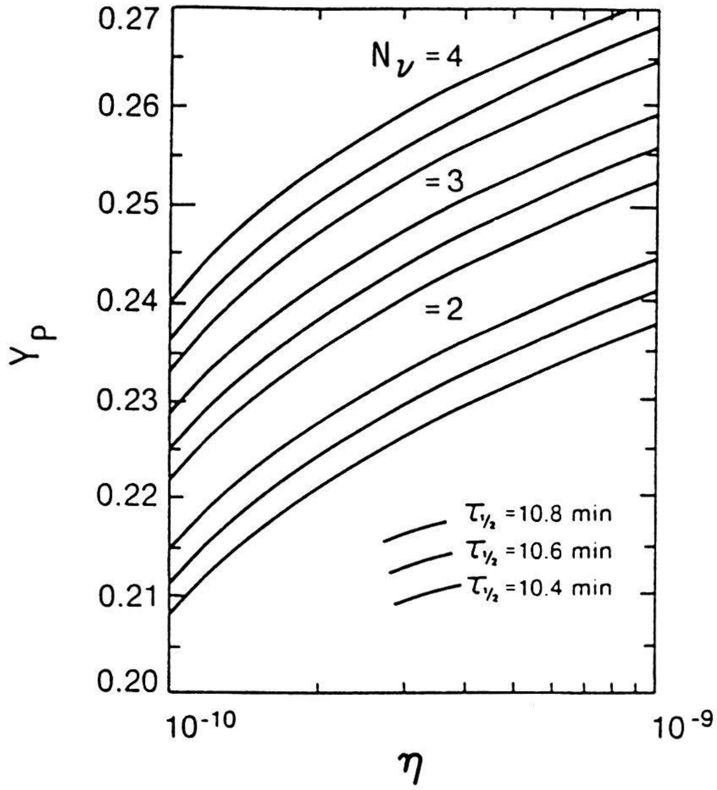


Fig. 2 The predicted [18] primordial abundances (by mass) of $^4He$ ($Y_p$) as a function of $\eta = n_B/n_\gamma$ for $N_\nu = 2,3,4$ and various values for $\tau_{\frac{1}{2}}$.

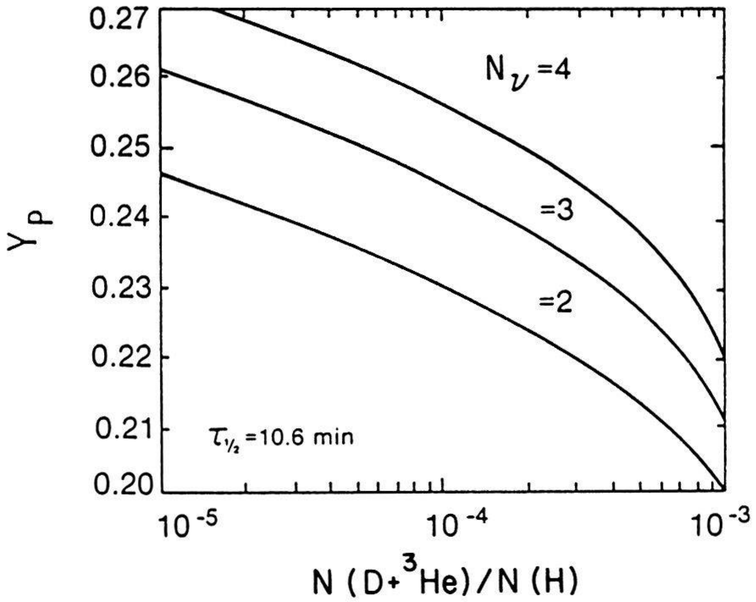


Fig. 3 The predicted [18] abundances (by mass) of $^4$He ($Y_p$) versus the predicted abundance (by number relative to H) of D plus $^3$He for $N_\nu$ = 2,3,4 and $\tau_{\frac{1}{2}}$ = 10.6 min.

## I.3 Plausible Extrapolation

Further back in time, hadrons begin to overlap above $T \sim 200$ MeV and we expect that all hadrons coalesce into each other and a plasma of the constituents of hadrons - quarks and gluons - was formed. Crude estimates lead also to the expectation that the latent heat of this transition might be considerable.

Since most of you are solid state physicists, it might be interesting to elaborate a bit on this. In principle, the pressure $p(T,\mu)$ as a function of $T$ and the chemical potential $\mu$ can be computed from

$$p(T,\mu) = \frac{T}{V} \ln \left(\mathrm{Tr}\, e^{-(H-\mu N)/kT}\right) ,$$

where $H$ is the Hamiltonian of QCD. The entropy density $s(T,\mu)$, the baryon number density $n(T,\mu)$ and the energy density $\varepsilon(T,\mu)$ are then derived from its partial derivatives

$$s(T,\mu) = \frac{\partial p}{\partial T} , \quad n(T,\mu) = \frac{\partial p}{\partial \mu} , \quad \varepsilon(T,\mu) = Ts - p + \mu n .$$

Simple phenomenological expressions for $p(T,\mu)$ in the hadron and in the quark-gluon phases give for $\mu = 0$ the curves sketched in Fig. 4.

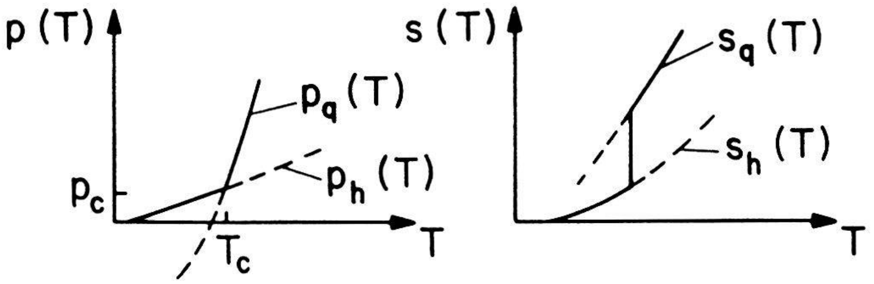


Fig. 4 Schematic picture for $p(T)$ and $s(T)$ for hadron matter and quark-gluon plasma ($\mu = 0$).

This deconfinement transition is clearly predicted by lattice QCD calculations. For pure glue the transition is a first order phase transition [23]. With quarks, it is, for $\mu = 0$, conclusively known only that the entropy density changes rapidly within a narrow temperature range.

It is not clear whether the deconfinement transition had important cosmological effects. One possibility is that it produced an inhomogeneous distribution of nucleons, but most probably these inhomogeneities were again smeared out long before nucleosynthesis. If not, they would affect the cosmic abundance of light elements.

A very speculative effect has been envisaged by Witten [24], who considered the possibility that the baryon inhomogeneities could have caused the production of nuggets of "strange quark matter", which might possibly be stable. Such strange quark balls belong now also to the long list of candidate constituents of the dark mass of the Universe. De Rújula [25] has studied how such nuggets could be searched for. Most people consider it as unlikely that they exist. The physics involved is, however, too difficult to make reliable predictions about this interesting possibility.

Lattice QCD calculations also predict that the condensates $\langle \bar{q}q \rangle$, $\langle G_{\mu\nu} \cdot G^{\mu\nu} \rangle$ of the quark and gluon fields "melt away" at $T \simeq 200 \pm 50$ MeV. Apart from quark-mass effects the chiral symmetry is thereby restored.

Another symmetry restoration in a phase transition is expected around the Fermi-scale $G_F^{-\frac{1}{2}} \simeq 300$ GeV. The effective Higgs potential (the free energy density) of the electroweak theory is temperature dependent; above a sufficiently high temperature, the minimum occurs for a vanishing expectation value $\langle\phi\rangle$ of the Higgs field $\phi$. In the high-temperature phase all gauge bosons should then be masseless (and additional scalar bosons should appear). This is very analogous to the transition of a type II superconductor to its normal state and the disappearance of the Meissner-Ochsenfeld effect (which is just the Higgs mechanism). This electroweak phase-transition has probably no important cosmological effects.

## I.4 Open Questions and Mysteries in the Standard Hot Big Bang-Model

The standard models of particle physics and cosmology are both incomplete in many respects. Let me briefly recall some of the main questions and puzzles, which cannot be answered in the standard hot big bang model.

### A. Why is the cosmological constant so small ?

During the electroweak and QCD (phase-) transitions the vacuum energy changes roughly by $(10^2\mathrm{GeV})^4$, resp. $(10^{-1}\mathrm{GeV})^4$. These are enormous changes in comparison to the vacuum energy of the present Universe, which is certainly less than*) $10^{-46}(\mathrm{GeV})^4$. This discrepancy constitutes one of the most profound problems

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*) The present vacuum energy cannot be much larger than the critical density (5):

$$\rho_c \simeq 8\times10^{-47}\ h_o^2\ \mathrm{GeV}^4 .$$

For the cosmological constant $\Lambda$ this implies

$$\Lambda \lesssim 8\pi G\,\rho_c = 0.9\times10^{-121}\ h_o^2\ m_{Pl}^2$$

of present day field theory and cosmology. Its solution is probably deeply rooted. At the same time, the problem of the cosmological constant indicates that inflationary models, which are based on dominant vacuum energies in very early epoches, might be much too naiv. So far the cosmological constant problem remains unsolved also in the superstring theory.

B. The dark matter problem

The discussion in section I.2 has shown that there is much more to the Universe than meets the eye. At most 10 % of the mass-energy is in known forms. It is even not excluded that the Universe contains dark matter near the critical density, if it is more uniformly distributed than the luminous matter (remember the new IRAS data). In this case, most astronomical tests do not take into account such an inert background. Many forms of dark matter have been proposed [26].

The most conservative (and perhaps most reasonable) attitude is, that $\Omega_0 = \Omega_B \simeq 0.1$ and that most of the baryons might have been incorporated in a pregalactic population of stars [27]. The most likely unseen remnants of such stars would be low mass objects $M < 0.1\ M_\odot$, like "Jupiters", Brown Dwarfs, etc. (see Fig. 5), or massive black holes of perhaps $10^3\ M_\odot$. Such objects would satisfy several constraints: Masses above $\simeq 0.1\ M_\odot$ would contribute too much background light and the remnants of "ordinary" massive stars of 10-100 $M_\odot$ would produce too much heavy elements. Very massive objects of $10^3\ M_\odot$ would not violate the nucleosynthesis constraint and would probably terminate their evolution by a collapse and thereby swallowing most of their mass. On the other hand, much heavier objects $\gtrsim 10^6\ M_\odot$ would have detectable effects, such as dynamical friction or accretion [28].

Gravitational lensing [29] will perhaps one day enable us to discover these baryonic remnants.

The list of proposed non-baryonic dark matter constituents is long; it includes massive neutrinos, axions, and many other "...inos", suggested by supersymmetric models.

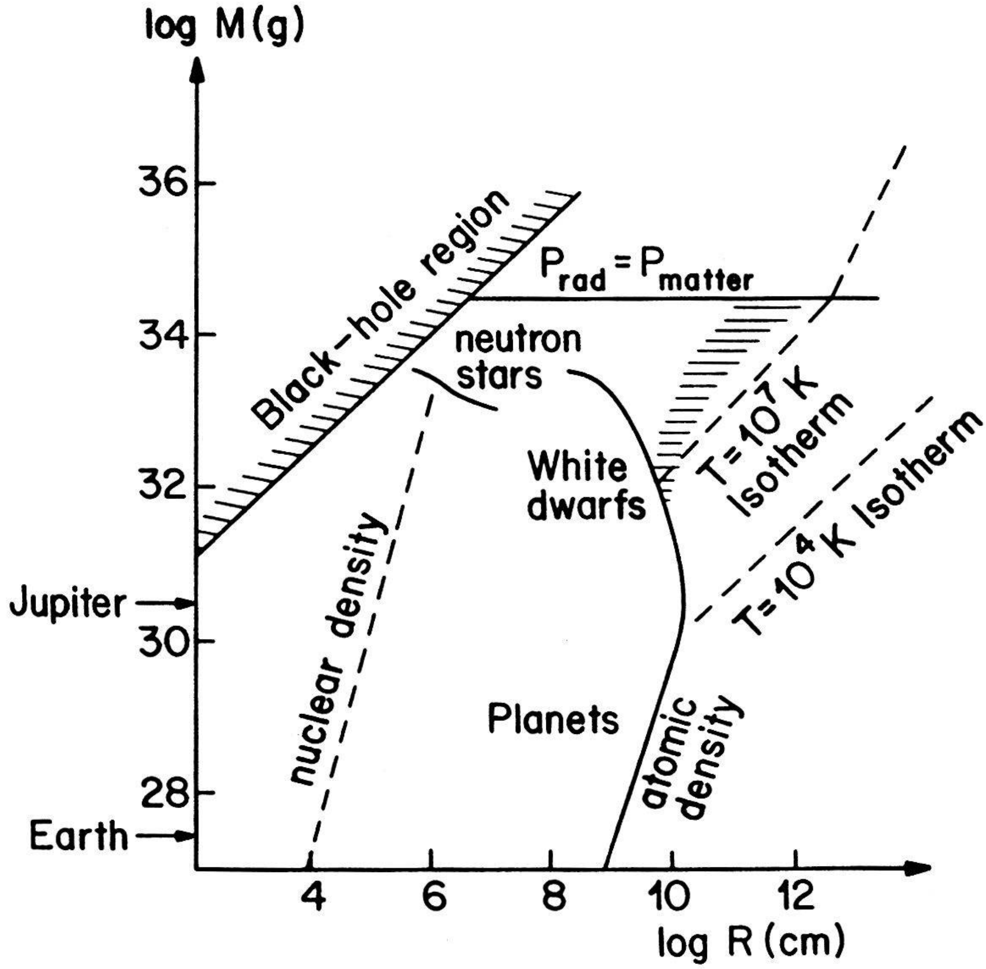


Fig. 5 Mass and length scales for various structures, in particular for cold objects. Also shown is the $T \simeq 10^7$ nuclear ignition isotherm.

Such exotic possibilities have been considered more and more serously in view of the difficulties to understand the formation of galaxies and their associations from small density fluctuations, which are compatible with the isotropy of the 3K background radiation *). We turn next to this topic.

### C. Primordial fluctuations, formation of galaxies and large-scale structures

The early Universe cannot have been completely smooth. Galaxies and their associations in clusters and superclusters

*) Possible experiments for direct detection of particle candidates for galactic dark matter are planned [30]. This is difficult, but some of the schemes look feasible.

have evolved by gravitational instabilities from small density fluctuations. The origin of the initial fluctuations and its spectrum is one of the major unresolved cosmological problems.

Two interesting proposals have so far been suggested. In inflationary scenarios the density fluctuations arise from initial quantum correlations of the Higgs field [31]. Another possibility is that cosmic strings produced the density fluctuations. This alternative will be discussed in section II.

The theory of galaxy formation is in a very fluent and uncertain state. We still do not know why galaxies exist !

On the observational side we see large-scale structures over tens or hundreds of megaparsecs. A recent survey [32] on the large-scale galaxy distribution shows that large, apparently empty, quasispherical voids dominate space, and that galaxies are concentrated in thin sheets and ridges between the holes (see Fig. 6). No satisfactory explanation for this foam-like appearance of the galaxy distribution exists.

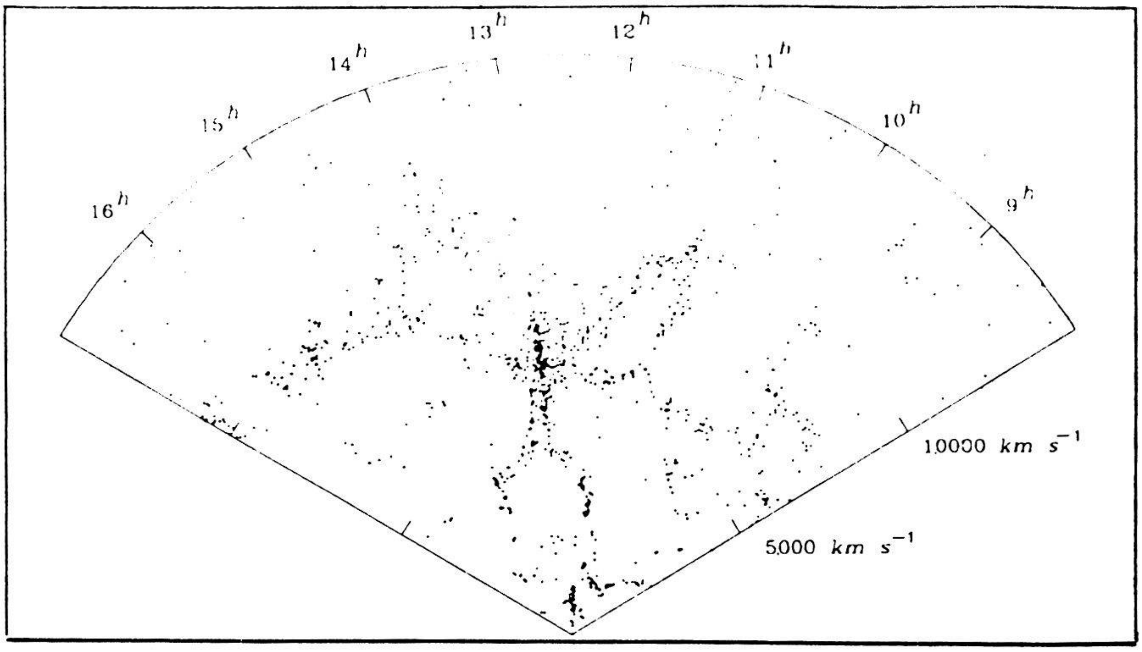


Fig. 6 Velocity distribution [32] in the declination wedge $2.65 \leq \delta \leq 32.5^{0}$ of 1'061 objects with magnitude $m_B \leq 15.5$ and velocity $v \leq 15'000$ km/s.

A crucial ingredient may be the concept of biased galaxy formation, which has been introduced recently and which has changed the discussion very much [33]. It presumes that galaxies only form in the rare peaks of an initial (gaussian) distribution of density perturbations. This biasing hypothesis enhances the large-scale structure. Voids develop, containing a large number of 'failed' galaxies, and galaxies form only in the dense regions between voids. The question what causes this biasing remains.

Furthermore, it is difficult to reproduce the most massive aggregates of galaxies and to account for the clustering of the great galaxy clusters. Positive spatial correlations have even be observed [34] among superclusters on a scale of $\sim 100\ h_0^{-1}$ Mpc, which appear to be stronger than those of galaxies and galaxy clusters. 'Cold' and 'hot' dark matter scenarios [26] have both their difficulties. One speculative way out invokes vacuum strings left over from a phase transition in the very early Universe and which may have induced non-gaussian fluctuation distributions. This possibility will be discussed later (section II).

D. The problem of initial conditions

It is by now well-known that the standard model has to be supplemented by a number of initial conditions, which are highly unnatural. Let us recall, as an example, the flatness (or entropy) problem. We have seen that the present density parameter is in the interval $0.1 \lesssim \Omega_0 \lesssim 2$ . In Friedman models $\Omega = 1$ is an unstable fixed point under time evolution. As a consequence, $\Omega$ should already be very close to one at the time of recombination: $|\Omega - 1| \lesssim 5\times10^{-4}$ at $t \sim 10^5$ yr. Shortly before nucleosynthesis starts, we find even the well-tuned value $|\Omega - 1| < 10^{-15}$ at $T = 1$ MeV.

Another naturalness problem is the horizon puzzle. These puzzles, as well as their speculative solutions by inflation [31], have been discussed so much in recent years [35] that I will not further elaborate on them. The original beauty of inflation is, lost, however, and among the many proposed schemes there is no generally accepted one (see also table I).

Another initial condition of the standard model, restricted to $T \lesssim 100$ GeV, is the baryon asymmetry, reflected in the small baryon/entropy ratio

$$n_B/s_\gamma \simeq 0.8\times10^{-8}\Omega_B\, h_o^2 \,. \tag{20}$$

Attempts to explain this small ratio within the framework of unified theories as a result of baryon-number-nonconserving processes in the very early Universe, lead us to the next topic.

## I.5 Speculations in Particle Physics and Cosmology

Particle physicists working on the unification program have only very few possibilities - such as proton decay - to test their theories in laboratory experiments. For this reason, cosmological considerations play an important role in constraining theoretical attempts. This has led to a symbiotic relationship between high energy physics and cosmology. Most of the exciting activity in cosmology during the last ten years has been in this direction.

The original excitement has, for various reasons, somewhat declined. The first unified theories are all beset with profound difficulties, like the "hierarchy problem" in GUTS (which requires a very fine tuning of parameters which should not be spoiled by radiative corrections). Proton decay has not be seen, ruling out the minimal SU(5) model of Giorgi and Glashow. Clearly, we are lacking crucial ingredients in our approach to grand unification.

Almost everybody's hope is now that superstring theory[36] will solve all the problems and will lead to a coherent unified theory of gravity and all other (gauge) interactions. Previous attempts (see table I) would then just be "effective low energy approximations" of the new ("final") theory, which are obtained through compactification and dimensional reductions. (See the invited lecture of A. Neveu at this meeting.) This new "theory of everything" would, of course, help us to understand the very early Universe, including the Planck era. But at the time of this conference, the theory is still in its infancy. Its phenomenological implications can only be worked out in detail, once the basic theoretical studies have further advanced. One cosmological implication is that there may exist another form of matter ('shadow matter') in the Universe, which only interacts with ordinary matter through gravity. Some constraints of this hidden form of matter (e.g. from nucleosynthesis) have already been discussed [37].

We supplement these brief remarks in table I, which summarizes the present situation. One particular aspect, namely the cosmic strings, will be discussed in the next section in some detail.

Table I. Unification program and the very early Universe

| Particle theories | Cosmological aspects | Remarks |
|---|---|---|
| (sypersymmetric) GUTs | Baryon asymmetry | no quantitative prediction of $n_B/s$ |
| (supers.) Kaluza-Klein (particularly in 11 and 10 dimensions) | Topological defects (domain walls, string, monopoles) | monopoles and domain walls should be avoided (inflation); strings may be great |
| Superstrings ("theory of everything") | exotic matter: massive neutrinos, axions, WIMPS (weakly interacting massive particles) | strongly restricted by cosmological arguments; may be needed for galaxy formation |
| | primordial fluctuations (quantum induced, strings) | power spectrum ? gaussian ? |
| ? <br> ? | Inflation | original beauty is lost, many schemes, non of which is accepted that would naturally give: <br> - sufficient inflation <br> - small density fluctuations <br> - sufficient reheating for baryosynthesis |
| | shadow matter | constraint by nucleosynthesis, etc. |

## II. COSMIC STRINGS

I have indicated already several times that vacuum strings might have very interesting cosmological consequences. In particular, they can generate the necessary density fluctuations for galaxy formation.

In this section, we describe briefly the properties, the formation and evolution of these topological defects and discuss some of their cosmological consequences.[38]

### II.1 The Physics of Vacuum Strings

#### A. Strings in gauge theories

Strings are topologically stable structures which may appear in the vacuum as a result of domains formed during the spontaneous breaking of a global or a gauge symmetry. This phenomenon is well-known from superfluidity and superconductivity (Helium II, type II - superconductors). We consider in the following only 'local' strings, arising through breakings of local symmetries. The prototyp are the magnetic flux vortices in type II - superconductors, which are solutions to the Ginzburg-Landau equations of superconductivity, first discovered by Abrikosov[39]. The appearance of the same objects in relativistic field theories was demonstrated by Nielson and Olesen [40] in their treatment of the Abelian Higgs model, where the Higgs field plays the role of the superconducting order parameter [41]. I recall briefly the string-solutions in the context of this field theoretical model.

The Lagrangian density of the Abelian Higgs model, with gauge group $U(1)$, is given by

$$\mathcal{L} = D_\mu \phi^* D^\mu \phi - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2}\lambda\left(\phi^*\phi - \eta^2\right)^2 , \qquad (21)$$

where $D_\mu = \partial_\mu - ieA_\mu$, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ and $e$ is the coupling constant. The minimum of the Higgs potential $V(\phi) = \frac{1}{2}\lambda(\phi^*\phi - \eta^2)^2$ is assumed on the vacuum manifold $M_0 = \{\phi : |\phi| = \eta\}$, which is in this simple example a circle.

We are interested in configurations for which the fields are static and constant in one direction taken to be the $x_3$-direction. Then we can chose a gauge with $A_0 = A_3 = 0$ and are thus led to look for regular finite energy solutions of the 2-dimensional Yang-Mills-Higgs-Lagrangian (21).

For a string solution $\phi$ is approaching the "Higgs-vacuum" far away from the vortex tube, i.e. $|\phi| \longrightarrow \eta$ and $|D\phi| \longrightarrow 0$ for $|x| \longrightarrow \infty$. The latter condition, together with $F_{\mu\nu} \longrightarrow 0$, guarantees that the energy density vanishes asymptotically outside the string core. Since $\phi$ lies asymptotically in the vacuum manifold $M_0$, the Higgs field defines a mapping $f: S^1 \longrightarrow M_0$, which associates to each direction in (ordinary) space the asymptotic value of the Higgs field in the vacuum manifold $M_0$, which in our example is again a circle $S^1$. The winding number (or degree) $n$ of the map $f$ is an integer and characterizes the homotopy class (i.e. the element in $\pi_1(M_0)$) of the mapping. This is a simple example of a topological 'quantum number' (charge).

It is easy to show [41] that the magnetic flux of the string is n times the elementary magnetic flux $2\pi/e$. An elementary string corresponds to $n = \pm 1$ and is topologically stable. Strings with $|n| \geqslant 2$ are probably unstable and decay into elementary ones.

The thickness of the string is determined by the Compton wavelengths of the Higgs particle and the vector mesons, $m_\phi^{-1}$, $m_A^{-1}$, where $m_\phi = \sqrt{2\lambda}\,\eta$, $m_A = \sqrt{2}\,e\,\eta$. Usually $m_\phi < m_A$ and the string has an inner core of 'false vacuum' with linear mass density $\mu_\phi \sim (\lambda\eta^4)\, m_\phi^{-2} \sim \eta^2$, and a tube of magnetic field of radius $m_A^{-1}$ with $\mu_A \sim B^2 m_A^{-2} \sim \eta^2$ (because the flux is $2\pi/e$). Thus, the total mass of the string per unit length is roughly

$$\mu \sim \eta^2 \ (\sim 10^{-6}\, m_{Pl}^2 \text{ in GUTS}) . \qquad (22)$$

(A more accurate treatment can be found in books on superconductivity.)

The internal structure of the string is often unimportant and quantities like the energy momentum tensor $T_{\mu\nu}$ can be averaged over cross sections. For a static straight string along the z-axis one finds from Lorentz invariance and the conservation law $T^{\mu\nu}{}_{,\nu} = 0$ easily [38]

$$\int T_{\mu\nu}\, dxdy = \mu(1,0,0,1) \; . \tag{23}$$

Note that

$$\mu = G\mu \frac{m_{Pl}}{\ell_{Pl}} = 2.2 \left(\frac{G\mu}{10^{-6}}\right) 10^{10}\, M_{\odot}/\text{kpc}. \tag{24}$$

More complicated strings are expected to occur, whenever $\pi_1(M_0)$ is nontrivial. In the general case of a symmetry breaking $G \longrightarrow H$ the vacuum manifold is $G/H$. If $\pi_1(G) = \pi_0(G) = I$ ($G$ is simply connected) one knows from homotopy theory that $\pi_1(G/H) = \pi_0(H)$. So $H$ should contain a discrete symmetry. This shows that string solutions are expected to occur for example in the scheme $SO(10) \longrightarrow SU(5) \times Z_2$. Strings form in many grand unified theories [38], as well as in the low energy sector of superstring theories [42].

At finite temperatures the effective potential for $\phi$ is temperature dependent and its minimum is at $\langle\phi\rangle = 0$ for $T$ larger than some critical temperature $T_c$. Thus, the symmetry is restored for $T > T_c$.

In the cosmological context, as the Universe cools through the critical temperature, fluctuation regions with $\langle\phi\rangle \neq 0$ develop, in which the directions of $\langle\phi\rangle$ are correlated over a correlation length $\xi$. For a second-order transition, the correlation length is $T_c^{-1}$, but can be much larger for a first-order transition. At any rate, $\xi$ is always smaller than the horizon length.

Since the free energy tends to be minimized, a slowly varying $\langle\phi\rangle$ is preferred. Much of the initial chaotic variations will therefore quickly die out in the course of further evolution. However, if $\pi_1(M_0)$ is non-trivial, the formation of topological defects is unavoidable.

## B. Dynamics of strings

We consider now the dynamics of macroscropic strings, whose dimensions are much larger than the thickness.

The motion of the string defines a 2-dimensional world sheet $x^{\mu}(\tau,\sigma)$, where $\tau$ is a timelike and $\sigma$ is a spacelike parameter. The action $S$ must be a functional of the dynamical variables $x^\mu(\tau,\sigma)$, which has the following properties: (i) S is invariant under general coordinate transformations; (ii) S should be invariant with respect to reparametrizations of the world sheet; (iii) the action should be an integral over the 2-dimensional world sheet $x^\mu(\tau,\sigma)$.

This determines S uniquely, up to an irrelevant factor:

$$S = -\mu\int \sqrt{-g^{(2)}}\, d\tau\, d\sigma \quad . \tag{25}$$

Here $g^{(2)}$ denotes the determinant of the induced metric on the world sheet. (In our applications, the metric of spacetime will be Minkowski or Friedman.) With the notation

$$\dot{x}^\mu = \frac{\partial x^\mu}{\partial \tau}, \qquad x'^\mu = \frac{\partial x^\mu}{\partial \sigma} \tag{26}$$

we have

$$g^{(2)} = \dot{x}^2\, x'^2 - (\dot{x}\cdot x')^2 \quad . \tag{27}$$

Thus, the Lagrangian for the equation of motion for a string is

$$L = -\mu\sqrt{(\dot{x}\cdot x')^2 - \dot{x}^2 x'^2} \quad . \tag{28}$$

The corresponding nonlinear equations of motion can be simplified by suitable choices of the parameters $(\tau,\sigma)$ and exact solutions can be found [38]. For instance, a circular loop will rapidly collapse, reaching the velocity of light as it shrinks to a "point". (When the size of the loop becomes comparable to the string width $\sim\eta^{-1}$, it decays into elementary particles.) Explizit solutions [43] suggest that a substantial fraction of loop trajectories never self-intersect.

## C. Gravitational radiation from oscillating loops

For non-intersecting loops the dominant energy-loss mechanism is gravitational radiation. Very roughly this can be estimated as follows. Loops of size R have frequencies $\omega \sim R^{-1}$ and the quadrupole formula gives for the energy loss

$$\dot{M} \sim -G(MR^2\omega^3)^2 \sim -G\mu^2 \quad , \tag{29}$$

where $M \sim \mu R$ is the mass of the loop. The lifetime of the loop is thus

$$\tau = M/|\dot{M}| \sim \frac{R}{G\mu} \tag{30}$$

and the energy loss fraction in one oscillating period is

$$\frac{|\dot{M}|}{M}\ \frac{2\pi}{\omega} \sim G\mu \ . \tag{31}$$

The dimensionless number $G\mu$ is crucial in all that follows. According to (22) we have

$$G\mu \sim G\eta^2 \sim \left(\frac{\eta}{m_{Pl}}\right)^2 \sim 10^{-6} \tag{32}$$

for a typical grand unification scale $\eta \sim 10^{16}$ GeV.

More accurate calculations can be performed for explicit loop solutions. One finds [44] instead of (29)

$$\dot{M} = -\gamma G \mu^2 \ , \tag{33}$$

where $\gamma$ depends on the solution, but is typically $\sim 100$ .

It turns out [45] that the gravitational radiation is more important than electromagnetic radiation for macroscopic loops ($M \gg m_{Pl}$) .

D. Intercommuting

When two strings intersect, they can change partners (intercommute), as shown in Fig. 7a. This phenomenon is also known in superconductivity. Double intersections and self-intersections can result in the formation of closed loops (Figs. 7b, c).

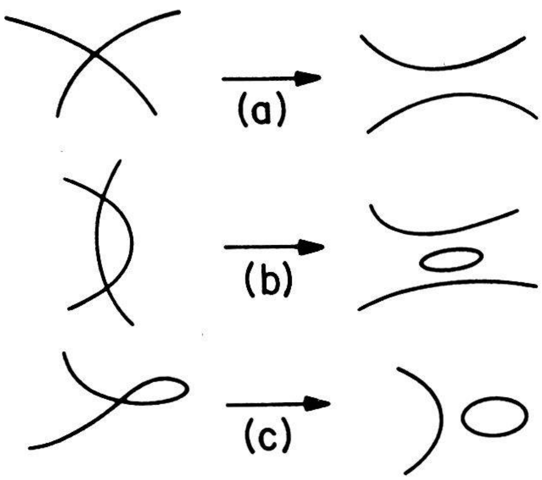


Fig. 7 Intercommuting of strings.

Numerical solutions [46] of the nonlinear field equations describing colliding strings suggest that intercommuting occurs with high probability.

## E. Gravitational field of strings

Eq.(23) shows that the string tension $\mu$ is equal to the energy per unit length of the string. This implies that the gravitational field is very different from that of massive rods. Indeed, the correct Newtonian limit for a quasistatic matter distribution is [47]

$$\Delta \varphi = 8\pi G \left(T_{00} - \tfrac{1}{2} \operatorname{tr} T\right)$$

and the right-hand side vanishes for (23).

It is not difficult to find [48] the exact solution of the Einstein equations for a string. For a straight string along the z-axis the metric outside the string core is in cyclindrical coordinates

$$ds^2 = dt^2 - dz^2 - dr^2 - (1-4G\mu)^2\, r^2 d\varphi^2 \ . \qquad (34)$$

This metric is locally flat, because the coordinate transformation $\varphi' = (1-4G\mu)\,\varphi$ brings it locally to Minkowskian form. Since $\varphi'$ varies from 0 to $(1-4G\mu)\,2\pi < 2\pi$, the metric (34) describes a "conical space", that is, a flat space with a wedge of angular size $8\pi G\mu$ taken out (see Fig. 8), whereby the two faces of the wedge must be identified.

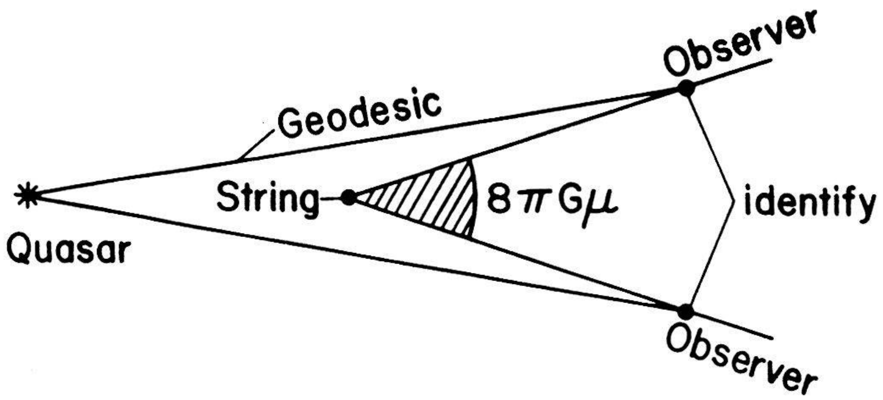


Fig. 8 The conical space outside a straight string core.

It is obvious from Fig. 8 that a string would produce double images of equal brightness of quasars. The typical separation angle between two images is $\sim 4\pi G\mu \sim 3$–$30$ arcsec for $G\mu \sim 10^{-6}$–$10^{-5}$.

This is a kind of classical analogue of the Aharanov-Bohm effect. There is no space-time curvature outside the string, but the metric is not globally Minkowskian.

The conical nature of space around a string has another interesting cosmological consequence. A wake is formed behind a relativistically moving string that has the shape of a wedge with an opening angle $\sim 8\pi G\mu$ and a density contrast $\delta\rho/\rho \sim 1$ . This effect may be important for the formation of large scale structures in 'cold' matter scenarios (see section II.2).

In order to avoid possible misunderstandings, we emphasize at this point that the long range time averaged field of a loop is just that of a point mass with mass equal to that of the loop. Loops could thus have served as seeds for the primordial density fluctuations.

F. Strings in the expanding Universe

If $\tau$ denotes the conformal time variable $(d\tau = dt/a(t))$, then the metric of a flat Friedman Universe has the form

$$ds^2 = a^2(\tau)\,[d\tau^2 - d\underline{x}^2]$$

and thus the Lagrangian for a string is just the one for Minkowski space, multiplied by $a^2(\tau)$ . Analytic and numerical analysis of the equations of motion show the following important facts [38]: (i) Waves <u>bigger</u> than the horizon are conformally stretched (both amplitude and wavelength grow like $a(\tau)$ , the shape of the string remaining unchanged). (ii) Irregularities on scales <u>smaller</u> than the horizon are smoothed out.

G. Interaction with particles

For the cosmological evolution of strings, one needs also the force of friction due to their interaction with light particles. The scattering cross section has been estimated by Everett [49] and this leads to the following force of friction per unit length acting on a moving string with velocity $v$ relative to the radiation

$$F_s \sim N_s\ T^3 v/\ln^2(T\delta)\ , \tag{35}$$

where $N_s$ is the number of light particles interacting with the fields of the string and $\delta$ is its width.

## II.2 Formation and Evolution of Strings

Initially, strings are formed as defect lines in the orientation of the Higgs field. One expects that they have the shape of random walks of step given by the correlation length $\xi$ and typical distances between neighboring string segments also of this magnitude. This is confirmed by Monte Carlo simulations [50] which show also: (i) About 80 % of the total string length is due to 'infinite' strings. (ii) The remaining strings are closed loops with a scale-invariant distribution $dn \propto R^{-4} dR$, where $R$ is the characteristic size of the loop.

The evolution of a network of cosmic strings is complicated in view of the physics discussed in the previous section. Qualitative discussions (see, for instance, the excellent review article of Vilenkin [38]) and numerical studies [51] lead to the following conclusions:

(1) Expansion of the Universe straightens out long strings on scales smaller than the horizon and conformally stretches them on scales greater than the horizon.

(2) There are just a few segments of 'infinite' strings per horizon volume at any time.

(3) Loops surviving the decay due to gravitational radiation at time $t$ have sizes greater than $\sim G\mu$ times the horizon length. In particular, the smallest surviving loops today have a size $\sim G\mu\, t_0 \sim$ few kpc, for $G\mu \sim 10^{-6}$.

(4) During the radiation and matter dominated eras, the length distributions of closed loops are, respectively

$$\frac{dn}{dR} \sim t^{-3/2} R^{-5/2} \ , \quad \frac{dn}{dR} \sim (tR)^{-2} \ .$$

Today, the smallest surviving loops have masses $M \sim \mu R \sim$ $\sim (G\mu/10^{-6})\, 10^{10}\, M_\odot$ and they are typically separated by $\sim 10$ Mpc.

(5) During all evolutionary phases, the energy density due to strings is always much smaller than the total energy density.

A discussion [38] of the spectra of density fluctuations generated by strings shows interesting properties. For example, for a baryon-dominated Universe matter starts to accrete on loops at the time of recombination and one finds a spectrum corresponding to the hierarchical gravitational clustering picture: the fluctuations increase towards smaller scales with a lower cut-off of a galactic mass $\sim 10^{12}\ M_{\odot}$ .

The string model for galaxy formation may help to solve the problems discussed in section I.4C. For example, loops of similar size are correlated [52] and this leads to matter perturbations which are highly non-Gaussian and have non-random phases [53]. Turok and Brandenberger [54] have discussed in particular the correlation function of Abell clusters.

Another interesting aspect of the string scenario is the following. Closed loops at the time of galaxy formation have in general much smaller sizes than those of the galaxies condensing around them. Therefore, accretion of matter onto loops may lead to the formation of massive compact objects (quasars and active galactic nuclei).

Even if strings have nothing to do with galaxy formation, they might produce some unique observational effects:

(1) One expects [55] that out of $\sim 10^4$ quasars, one is doubled by a string. It is, however, difficult to prove that a particular pair of quasar images is formed by a cosmic string. Paczynski [56] has, therefore, suggested that observations of very distant galaxies with the Space Telescope may be a better way for discovering cosmic strings. A galaxy may appear cut by a sharp edge if there is a cosmic string between the galaxy and the observer. Since the probability for this is quite small, an extensive observational program would be necessary to see this.

(2) Kaiser and Stebbins [57] have pointed out that rapidly moving strings would produce steplike discontinuities in the 3K-background with $\Delta T/T \sim 10\ G\mu$ . Present observational limits are consistent with $G\mu < 10^{-5}$ .

(3) Gravitational waves produced by oscillating loops add up to a stochastic gravitational background which would, for

example, induce a "timing noise" of the millisecond pulsar[58]. With improved observations, this effect may become detectable within several years, if $G\mu \sim 10^{-6}$. If not, cosmic strings will be pretty well ruled out as candidate sources for any large-scale fluctuations having a significant impact on galaxy clustering.

## References

1) W. Buchmüller and D. Wyler, Nucl.Phys. B268, 621 (1986).
2) G.F.R. Ellis, Quarterly Journal of the Royal Astronomical Society 16, 245 (1975).
3) Talks given by L. Maiani and L. Di Lella at the High Energy Physics Conference of the European Physical Society, Bari - 1985.
4) L. Di Lella, Proc.Intern.Symp. on lepton and photon interactions at high energies (Kyoto, 1985).
5) S. Komamiya, Proc.Intern.Symp. on lepton and photon interactions at high energies (Kyoto, 1985).
6) M. Fritschi, E.Holzschuh, W.Kündig. J.W.Petersen, R.E. Pixley, and H.Stüssi, Zürich Univ. preprint (1986).
7) V.A. Lubimov et al., Phys.Lett. 159B, 217 (1985), and references therein.
8) R. Abela et al., Phys.Lett. 146B, 431 (1984); Argus Collaboration, Phys.Lett. B163, 404 (1985).
9) For a review, see K. Winter, 2nd ESO/CERN Symposium on "Cosmology, Astronomy and Fundamental Physics", ESO (Garching b. München), 1986.
10) S.P. Mikheyev and Y.Yu. Smirnov, in Proceedings of the Tenth International Workshop on Weak Interactions, Savonlinna, Finland, 1985.
11) L. Wolfenstein, Phys.Rev. D17, 2369 (1978); Phys.Rev. D20, 2634 (1979).
12) See also: H.A. Bethe, Phys.Rev.Lett. 56, 1305 (1986).

13) A. Sandage and G.A. Tammann, in Proceedings of the First ESO/CERN Symposium on "Large-Scale Structure of the Universe, Cosmology and Fundamental Physics, ed. G. Setti, L. Van Hove (ESO and CERN, 1984); M.R. Robinson, "The Cosmological Distance Ladder", Freeman and Company, 1985.
14) G.A. Tammann and A. Sandage, Ap.J. 294, 81 (1985).
15) A. Sandage and G.A. Tammann, Ref. 13).
16) M.R. Robinson, in 2nd ESO/CERN Symposium on "Cosmology, Astronomy and Fundamental Physics", ESO (Garching b. München), 1986; A. Yahil et al., Ap.J. 301, L1 (1986).
17) F. Melchiori, in 2nd ESO/CERN Symposium, loc.cit.
18) For a recent recalculation and references, see J. Yang et al., Ap.J. 281, 492 (1984).
19) A. Boesgaard and G. Steigman, Ann.Rev.Astr.Astrophys. 23, 319 (1985).
20) J. Ellis et al., Phys.Lett. 167B, 457 (1986).
21) D.N. Schramm, private communication.
22) P. Candelas et al., Nucl.Phys. B258, 46 (1985)
23) For a review, see H. Satz, Ann.Rev.Nucl.Part.Sci., 35, 245 (1985).
24) E. Witten, Phys,Rev. D30, 272 (1984).
25) A. De Rujula, Nucl.Phys. A434, 605c (1985).
26) For a review, see J. Primack, Lectures at the International School of Physics Enrico Fermi, Varenna 1984 (SLAC-preprint 3387, 1984).
27) M.J. Rees, in Proceedings of Ref. 9) and references therein, to appear.
28) B.J. Carr, J.R. Bond and W.D. Arnett, Ap.J. 277, 445 (1984).
29) P. Schneider, Ap.J. 30, L31 (1986), and references therein.
30) P.F. Smith, in Proceedings of Ref. 9).
31) For detailed surveys, see R. Brandenberger, Rev.Mod.Phys. 57, 1 (1985); A.D. Linde, Rep.Progr.Phys. 47, 925 (1984).
32) M. Geller, J. Huchra and V. de Lapparent, Ap.J. 302, L1 (1986).
33) J.A. Peacock and A.F. Heavens, Mon.Not.R.astr.Soc. 217, 805 (1985); J.Bardeen et al., Ap.J. in press.

34) N.A. Bahcall and W.S. Burgett, Ap.J. 300, L 35 (1986).
35) G. Börner and N. Straumann, Phys.Bl. 41, 146 (1985).
36) A. Neveu, talk at this conference; M. Green, in Proceedings of Ref. 9), to appear.
37) E. Kolb, D. Selzel and M. Turner, Nature 314, 415 (1985).
38) For an excellent review of cosmic strings, see A. Vilenkin, Phys.Rep. 121, 263 (1985). Cosmic strings were first discussed by Y.B. Zel'dovich, Mon.Not.R.Astron.Soc. 192, 663 (1980), and by A. Vilenkin, Phys.Rev.Lett. 46, 1169 (1981).
39) A.A. Abrikosov, Zh.Eksp.Teor.Fiz. 32, 1442 (1957) [Sov.Phys. JETP 13, 451 (1961)].
40) H.B. Nielsen and P. Olesen, Nucl.Phys. B61, 45 (1973).
41) See, for example, N. Straumann, in Proceedings of the 1982 SIN Spring School: Beyond the Standard Model.
42) E. Witten, Phys.Lett. 153B, 243 (1985).
43) N. Turok, Nucl.Phys. B242, 520 (1984); T.W.B. Kibble and N. Turok, Phys.Lett. 116B, 141 (1982).
44) T. Vachaspati and A. Vilenkin, Phys.Rev. D31, 3052 (1985).
45) T. Vachaspati, A.E. Everett, and A. Vilenkin, Phys.Rev. D30, 2046 (1984).
46) P. Shellard, unpublished.
47) See e.g., N. Straumann, "General Relativity and Relativistic Astrophysics", Texts and Monographs in Physics, Springer (Berlin), 1984.
48) J.R. Gott, Ap.J. 288, 422 (1985).
49) A.E. Everett, Phys.Rev. D24, 858 (1981).
50) T. Vachaspati and A. Vilenkin, Phys.Rev. D30, 2036 (1984).
51) A. Albrecht and N. Turok, Phys.Rev.Lett. 54, 1868 (1985).
52) N. Turok, UCSB preprint (1985).
53) N. Turok and D.N. Schramm, Nature 312, 598 (1984)
54) N. Turok and R.H. Brandenberger, UCSB preprint (1985); N. Turok, in Proceedings of Ref. 9).
55) A. Vilenkin, Ap.J. 282, L51 (1984).
56) B. Paczynski, Nature 319, 567 (1986).
57) N. Kaiser and A. Stebbins, Nature 310, 391 (1984).
58) C.J. Hogan and M.J. Rees, Nature 311, 109 (1984).