*CONFERENCES AND SYMPOSIA*

PACS numbers: 01.10.Fv, **12.60.–i**, **95.35.+d**, **98.80.–k**,

# In memory of V L Ginzburg
(Scientific session of the Physical Sciences Division of the Russian Academy of Sciences, 27 October 2010)

DOI: 10.3367/UFNe.0181.201106e.0655

The scientific session of the Physical Sciences Division (PSD) of the Russian Academy of Sciences (RAS), dedicated to the memory of V L Ginzburg, took place in the conference hall of the Lebedev Physical Institute, RAS on 27 October 2010.

The agenda of the session announced on the website www.gpad.ac.ru of the RAS Physical Sciences Division listed the following reports:

(1) **Mesyats G A** (Lebedev Physical Institute, RAS, Moscow) "Introductory word";

(2) **Rubakov V A** (Institute for Nuclear Research, RAS, Moscow) "Cosmology and the Large Hadron Collider";

(3) **Gurevich A V** (Lebedev Physical Institute, RAS, Moscow), **Zelenyi L M** (Space Research Institute, RAS, Moscow) "Intense gamma bursts in Earth's atmosphere (TGE) and the mission 'Chibis'";

(4) **Vasiliev M A** (Lebedev Physical Institute, RAS, Moscow) "Higher-spin theory";

(5) **Maksimov E G** (Lebedev Physical Institute, RAS, Moscow) "What is and what is not known about HTSC";

(6) **Pudalov V M** (Lebedev Physical Institute, RAS, Moscow, and Moscow Institute of Physics and Technology) "V L Ginzburg and the development of experimental work on high-temperature superconductivity at LPI: 'iron superconductors.'"

Papers based on talks 2, 4, and 6 are published below.

For several reasons, L P Pitaevskii was unable to attend the session. He presented a paper dedicated to the memory of V L Ginzburg, which is published in this issue of *Physics–Uspekhi* (p. 625).

*Uspekhi Fizicheskikh Nauk* **181** (6) 405 – 447 (2011)
DOI: 10.3367/UFNr.0181.201104g.0405
Translated by V A Rubakov, M A Vasiliev, E G Strel'chenko; edited by A M Semikhatov and A Radzig

PACS numbers: **12.60.–i**, **95.35.+d**, **98.80.–k**
DOI: 10.3367/UFNe.0181.201106f.0655

# Cosmology and the Large Hadron Collider

V A Rubakov

## 1. Introduction

It is conceivable that the Large Hadron Collider, LHC, will give solutions to, or at least shed light on, major cosmological problems, such as the nature and origin of dark matter and generation of matter–antimatter asymmetry. We give several examples showing the LHC potential to uncover the history of the early Universe: WIMPs as cold dark matter, gravitinos as warm dark matter, and electroweak baryogenesis as a mechanism for generating matter–antimatter asymmetry.

The startup of the LHC is a major event, not only in particle physics but also in cosmology. The Universe we know is full of mysteries. It hosts matter but not antimatter, and 40 years since it was understood that this presents a problem, we do not have an established theory explaining this asymmetry. The Universe hosts dark matter, and we do not know what it is made of. There is dark energy in the Universe whose nature is also obscure. The LHC may well shed light at least on some of these mysteries, which are important items in V L Ginzburg's famous list. Optimistically, the LHC experiments may discover dark matter particles and their companions, and establish the physics underlying the matter–antimatter asymmetry. Otherwise, they will rule out some very plausible scenarios; this will also have profound consequences for our understanding of the early Universe. There are also exotic hypotheses on the physics beyond the Standard Model, like the TeV-scale gravity; their support by the LHC will have a dramatic effect for cosmology and is hard to overestimate.

Here, we concentrate on a few examples highlighting the LHC cosmological potential. We first turn to dark matter, and present the WIMP scenario for cold dark matter, which is currently the most popular, and for good reason. We also consider a light gravitino scenario for warm dark matter. Both are to be probed by the LHC, as they require a rather particular new physics in the LHC energy range. We then discuss electroweak baryogenesis — a mechanism for the generation of matter–antimatter asymmetry that may have operated at a temperature of the order of 100 GeV in the early Universe. This mechanism also needs new physics at energies 100 – 300 GeV, and it will be definitely confirmed or ruled out by the LHC.

In our presentation, we necessarily omit numerous details, while trying to make the basic ideas and results clear. More complete accounts of the particle physics aspects of cosmology may be found in book [1] and reviews [2–7]. Dark matter, including various hypotheses about its particles, is reviewed in [8–11]. Electroweak baryogenesis is discussed in detail in reviews [12–15].

**V A Rubakov** Institute for Nuclear Research, Russian Academy of Sciences, Russian Federation. E-mail: rubakov@ms2.inr.ac.ru

*Uspekhi Fizicheskikh Nauk* **181** (6) 655 – 664 (2011)
DOI: 10.3367/UFNr.0181.201106f.0655
Translated by V A Rubakov; edited by A M Semikhatov

Before coming to our main topics, we overview the present energy balance in the Universe. It is characterized by the parameters $\Omega_i = \rho_{i,0}/\rho_c$, where $\rho_{i,0}$ is the present energy density of the type of matter $i$, and

$$\rho_c = 5 \times 10^{-6}\ \text{GeV cm}^{-3}$$

is the total energy density at the present epoch. This value is obtained from the measurement of the Hubble constant, assuming the validity of General Relativity and spatial flatness of the Universe (the precision at which the spatial flatness is established experimentally is less than 2 per cent deviation of the total energy density from $\rho_c$ [16]). Clearly, in a spatially flat universe, $\sum_i \Omega_i = 1$, where the sum ranges over all forms of energy. The known forms of matter in the Universe are mostly photons of the Cosmic Microwave Background (CMB), whose temperature is $T_0 = 2.726$ K, plus baryons and neutrinos. The present number density of CMB photons is $n_\gamma = 410\ \text{cm}^{-3}$, and their energy density is $\rho_{\gamma,0} = 2.7 \times 10^{-10}$ GeV cm$^{-3}$. We see that the energy density of photons is very small today compared to the total energy density. We note in passing that an important characteristic of the early Universe is the entropy density. It is of the order of the number density of photons. More precisely, in thermal equilibrium at temperature $T$, it is given by

$$s = \frac{2\pi^2}{45}\, g_* T^3\,, \tag{1}$$

where $g_*$ is the number of the degrees of freedom with $m \lesssim T$, that is, the degrees of freedom that are relativistic at the temperature $T$ (fermions contribute with a factor of $7/8$). The entropy stays *exactly* constant in a comoving volume, unless there are fairly exotic processes of entropy production. The present value of the entropy density (taking neutrinos into account as if they were massless) is

$$s_0 \approx 3000\ \text{cm}^{-3}\,. \tag{2}$$

There are two ways of measuring the mass density of baryons. One is related to the Big Bang Nucleosynthesis, the epoch of thermonuclear reactions ($T \sim 10^9$ K). The resulting light element abundances depend on the baryon-to-photon ratio in that epoch, which has stayed constant since then. Comparing theory with observations of light element abundances yields

$$\eta \equiv \frac{n_B}{n_\gamma} \approx 6 \times 10^{-10}\,. \tag{3}$$

The energy density of baryons in the present Universe is therefore equal to

$$\rho_{B,0} = m_B\, n_{B,0} \approx 2.5 \times 10^{-7}\ \text{GeV cm}^{-3}\,, \tag{4}$$

or, in terms of the proportion of the total energy density,

$$\Omega_B = 0.045\,.$$

The same value is independently obtained by the analysis of the CMB temperature anisotropy. By electric neutrality, the number density of electrons is nearly the same as that of baryons, and they therefore contribute a negligible portion of the total energy.

The remaining known stable particles are neutrinos. Their number density is calculable in the Hot Big Bang theory, and these calculations are nicely confirmed by Big Bang Nucleosynthesis. The present number density of each neutrino type is $n_{\nu_\alpha} = 115\ \text{cm}^{-3}$, where $\nu_\alpha = \nu_e, \nu_\mu, \nu_\tau$. The direct bound on the mass of the electron neutrino, $m_{\nu_e} < 2.6$ eV, along with the observations of neutrino oscillations, suggest that every type of neutrino has a mass smaller than 2.6 eV (neutrinos with masses above 0.05 eV must be degenerate, according to the neutrino oscillation data). The present energy density of all types of neutrinos is therefore smaller than $\rho_c$:

$$\rho_\nu^{\text{total}} = \sum_\alpha m_{\nu_\alpha} n_{\nu_\alpha} \lesssim 8 \times 10^{-7}\ \text{GeV cm}^{-3}\,,$$

which means that $\Omega_\nu^{\text{total}} < 0.16$. This estimate does not use any cosmological data. In fact, cosmological observations give a stronger bound:

$$\Omega_\nu^{\text{total}} \lesssim 0.014\,. \tag{5}$$

In terms of the neutrino masses, bound (5) is given by $\sum m_{\nu_\alpha} \lesssim 0.6$ eV [17–19], and hence every neutrino must be lighter than 0.2 eV. On the other hand, atmospheric neutrino data, as well as the K2K and MINOS experiments, show that the mass of at least one neutrino must be larger than 0.05 eV. Comparing these numbers, we see that it may be feasible to measure neutrino masses by cosmological observations (!) in the future.

We conclude that most of the energy density in the present Universe is not in the form of known particles; most energy in the present Universe must be in 'something unknown.' Furthermore, there is strong evidence that this something has two components: clustered (dark matter) and unclustered (dark energy).

Clustered dark matter presumably consists of new stable massive particles. These form clumps of energy density that constitutes most of the mass of galaxies and clusters of galaxies. There are a number of ways to estimate the contribution of nonbaryonic dark matter to the total energy density of the Universe (see Refs [8–11, 20] for details):

— The composition of the Universe affects the angular anisotropy of the cosmic microwave background. Quite accurate measurements of the CMB anisotropy, available today, allow estimating the total mass density of dark matter.

— Nonbaryonic dark matter is crucial for the structure formation of the Universe (see below). Comparison of the results of numerical simulations of structure formation with observational data gives a reliable estimate of the mass density of nonbaryonic clustered dark matter.

The bottom line is that the nonrelativistic component constitutes about 28 per cent of the total present energy density, which means that nonbaryonic dark matter has

$$\Omega_{DM} \approx 0.23\,; \tag{6}$$

the rest is due to baryons.

There is direct evidence that dark matter exists in the largest gravitationally bound objects, clusters of galaxies. It comes from the determination of gravitational potentials in clusters via measuring the velocities of galaxies, the X-ray temperature of the intracluster gas, gravitational lensing effects, and so on. These methods enable directly determining the mass-to-light ratio in clusters of galaxies. Assuming that this ratio applies to all matter in the Universe,[1] we arrive at the estimate for the mass density of clumped matter.

[1] This is a strong assumption, because only about 10 percent of galaxies are in clusters.

Remarkably, this estimate coincides with (6). Finally, dark matter also exists in galaxies. Its distribution is measured by observations of the rotation velocities of distant stars and gas clouds around a galaxy.

Nonbaryonic clustered dark matter is not the whole story. The above estimates yield an estimate for the energy density of all particles, $\Omega_\gamma + \Omega_B + \Omega_\nu^{total} + \Omega_{DM} \approx 0.3$. This implies that 70 percent of the energy density is unclustered. This component is called dark energy; it is responsible for the present accelerated expansion of the Universe. One candidate is the vacuum energy density, or the cosmological constant (see, e.g., Refs [20–26] for the reviews).

All this fits all cosmological observations nicely, but does not fit the Standard Model of particle physics. It is our hope that the LHC will shed light at least on some of the properties of the Universe.

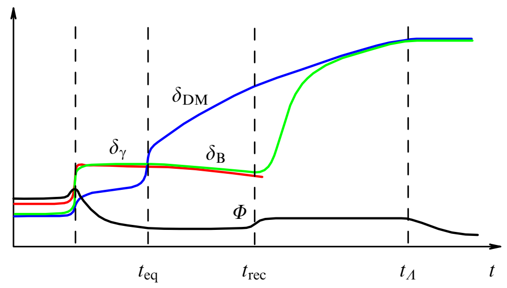


**Figure 1.** Time dependence, in the linearized theory, of density contrasts of dark matter, baryons, and photons, $\delta_{DM} \equiv \delta\rho_{DM}/\rho_{DM}$, $\delta_B$ and $\delta_\gamma$, as well as the Newtonian potential $\Phi$. $t_{eq}$ and $t_\Lambda$ correspond to the transitions from radiation domination to matter domination, and from decelerated expansion to accelerated expansion, and $t_{rec}$ refers to the recombination epoch.

## 2. Dark matter

Dark matter is characterized by the mass-to-entropy ratio

$$\left(\frac{\rho_{DM}}{s}\right)_0 = \frac{\Omega_{DM}\rho_c}{s_0} \approx \frac{0.23 \times 5 \times 10^{-6}\ \text{GeV cm}^{-3}}{3000\ \text{cm}^{-3}} = 4 \times 10^{-10}\ \text{GeV}\,. \quad (7)$$

This ratio has been constant in time since the freeze out of dark matter density: both the number of dark matter particles (and hence their mass) and the entropy are constant in a comoving volume.

Dark matter is crucial for our existence, for the following reason. Density perturbations in the baryon–electron–photon plasma before recombination do not grow because of high pressure, which is mostly due to photons; instead, perturbations produce sound waves of constant amplitude that propagate in the plasma. Hence, in a universe without dark matter, density perturbations in the baryonic component would start to grow only after baryons decouple from photons, i.e., after recombination. The physics behind the growth is quite simple: an overdense region gravitationally attracts surrounding matter; this matter falls into the overdense region, and the density contrast increases. In an expanding matter-dominated universe, this gravitational instability results in the density contrast growing as $(\delta\rho/\rho)(t) \propto t^{2/3} \propto T^{-1}$. Hence, in a universe without dark matter, the growth factor for baryon density perturbations is at most [2]

$$\frac{a(t_0)}{a(t_{rec})} = 1 + z_{rec} = \frac{T_{rec}}{T_0} \approx 10^3\,. \quad (8)$$

The initial amplitude of density perturbations is very well known from CMB anisotropy measurements: $(\delta\rho/\rho)_i = 1.5 \times 10^{-4}$. Hence, a universe without dark matter would still be rather homogeneous: the density contrast would be in the range of ten percent. No structure would have been formed, no galaxies, no life. No structure would be formed in the future either, because the accelerated expansion due to dark energy would soon terminate the growth of perturbations.

[2] Because of the presence of dark energy, the growth factor is even somewhat smaller.

Because dark matter particles decoupled from plasma much earlier than baryons did, perturbations in dark matter started to grow much earlier. The corresponding growth factor is larger than (8), and hence the dark matter density contrast at galactic scales reaches a value close to unity and perturbations enter the nonlinear regime and form dense dark matter clumps at the redshift $z \simeq 5$ and even earlier. After recombination, baryons fall into potential wells formed by dark matter, and then dark matter and baryon perturbations develop together. Galaxies form in those regions where dark matter was overdense originally. The development of perturbations in our Universe is shown schematically in Fig. 1. For this picture to hold, dark matter particles must be nonrelativistic early enough, because relativistic particles travel through gravitational wells instead of being trapped there. This means, inter alia, that neutrinos cannot make up a considerable part of dark matter, whence bound (5) follows.

Depending on the mass of dark matter particles and the mechanism of their production in the early Universe, dark matter may be *cold* (CDM) or *warm* (WDM). If dark matter particles were in thermal equilibrium with cosmic plasma in the early cosmological epoch, the CDM and WDM cases respectively correspond to heavy and light new particles:

$$m_{DM} \gtrsim 100\ \text{keV} \quad \text{for} \quad \text{CDM}\,, \quad (9)$$

$$m_{DM} \lesssim 100\ \text{keV} \quad \text{for} \quad \text{WDM}\,. \quad (10)$$

We discuss the warm dark matter option later, and now proceed with CDM.

### 2.1 WIMPS: Best guess for cold dark matter

There is a simple mechanism for the generation of dark matter in the early Universe. It applies to *cold* dark matter. Because of its simplicity and robustness, it is regarded by many as very likely, and the corresponding dark matter candidates — weakly interacting massive particles, WIMPs — are considered the best candidates. We describe this mechanism in general terms.

We assume that a heavy stable neutral particle Y exists and that Y particles can only be destroyed or created via their pair annihilation or creation, with annihilation products being the particles of the Standard Model. If the annihilation cross section is large enough, the overall cosmological behavior of the Y particles is as follows. At high temperatures, $T \gg m_Y$, the Y particles are in thermal equilibrium with the rest of the cosmic plasma; there are many Y particles in the plasma, and they are continuously being created and annihilated. As the temperature decreases below $m_{\rm Y}$, the equilibrium number density decreases. At some 'freeze-out' temperature $T_{\rm f}$, the number density becomes so small that the Y particles can no longer meet each other during the Hubble time, and their annihilation terminates. After that, the number of surviving Ys remains time independent in the comoving volume, and these relic particles contribute to the mass density in the present Universe.

It is straightforward to calculate the present mass density of Y particles in this scenario. The result is [3]

$$\frac{\rho_{\rm Y,0}}{s_0} = \frac{m_{\rm Y} n_{\rm Y,0}}{s_0} \simeq 7.6\,\frac{\ln\left(\sqrt{g_*}\,M_{\rm Pl} m_{\rm Y}\langle\sigma v\rangle\right)}{\sqrt{g_*}\,M_{\rm Pl}\langle\sigma v\rangle}\,, \qquad (11)$$

where $g_* = g_*(T_{\rm f})$ is the number of degrees of freedom upon freeze-out of the Y particles, $\sigma$ is the annihilation cross section, $v$ is the velocity of Y particles, and angle brackets denote thermal average, also at the freeze-out temperature:

$$T_{\rm f} \simeq \frac{m_{\rm Y}}{\ln\left(\sqrt{g_*}\,M_{\rm Pl} m_{\rm Y}\langle\sigma v\rangle\right)}\,. \qquad (12)$$

Formula (11) is quite remarkable. The mass density depends mostly on one parameter, the annihilation rate per particle, $\langle\sigma v\rangle$. The dependence on the Y-particle mass is through the logarithm and through $g_*(T_{\rm f})$; it is very weak. The value of the logarithm here is between 20 and 40, depending on the parameters (this means, in particular, that freeze-out occurs when the temperature decreases 20 to 40 times below the Y-particle mass). Substituting the numerical value $g_*(T_{\rm f}) \sim 100$ characteristic of the Standard Model and comparing with (7), we obtain the estimate

$$\langle\sigma v\rangle = (1-2)\times 10^{-36}~{\rm cm}^2\,. \qquad (13)$$

This is a weak-scale cross section, which tells us that the relevant energy scale is the TeV scale. We note in passing that estimate (13) is quite precise and robust.

The annihilation rate can be parameterized as $\langle\sigma v\rangle = \alpha^2/M^2$, where $\alpha$ is some coupling constant and $M$ is the mass scale (which may be higher than $m_{\rm Y}$). This parameterization is particularly appropriate for s-wave annihilation; it is suggested by the picture of pair annihilation of nonrelativistic Y particles via the exchange with another particle of mass $M$. With $\alpha \sim 10^{-2}$, the estimate for the mass scale is $M \sim 1$ TeV.

Thus, under very general assumptions, we find that the nonbaryonic dark matter may naturally originate from the TeV-scale physics. In fact, what we have found can be understood as an approximate equality between the cosmological parameter, the mass-to-entropy ratio of dark matter, and the particle physics parameters,

$$\left(\frac{\rho_{\rm DM}}{s}\right)_0 \simeq \frac{1}{M_{\rm Pl}}\left(\frac{\rm TeV}{\alpha_{\rm W}}\right)^2,$$

where $\alpha_{\rm W}$ is the electroweak gauge constant. The quantities in both sides of the above relations are both of the order of $10^{-10}$ GeV, and it is very tempting to think that this is not a mere coincidence. If it is not, the dark matter particle should be found at the LHC.

Of course, the most prominent candidate for the WIMP is the neutralino of the supersymmetric extensions of the Standard Model [27, 28]. The situation with the neutralino is somewhat difficult, however: pair annihilation of neutralinos often occurs in the p-wave, rather than in the s-wave. This results in a suppression factor in the annihilation rate proportional to $v^2 \sim T_{\rm f}/m_{\rm Y} \sim 1/30$. Hence, neutralinos tend to be overproduced in most of the parameter space of the MSSM and other models. Yet the neutralino remains a good candidate, especially at high $\tan\beta$.

### 2.2 Warm dark matter: light gravitinos

The cold dark matter scenario successfully describes the bulk of the cosmological data. Yet there are clouds above it. First, according to numerical simulations, the CDM scenario tends to overproduce small objects — dwarf galaxies: it predicts hundreds of dwarf galaxies in the vicinity of a large galaxy like the Milky Way, whereas only dozens of dwarfs have been observed so far (see, e.g., Ref. [30]). Second, again according to simulations, CDM tends to produce densities in galactic centers that are too high (cusps in density profiles); this feature is not confirmed by observations, either (see, e.g., Ref. [30] and the references therein). There is no crisis yet, but one may be motivated to analyze the possibility that dark matter is not that cold.

An alternative to CDM is warm dark matter, whose particles have energies of the order of $T$, where $T \gtrsim m$ after decoupling ($m$ is their mass). Then their spatial momenta decrease linearly with temperature, i.e., the momenta are approximately equal to $T$ all the time after decoupling. WDM particles become nonrelativistic at $T \sim m$. Only after that do the WDM perturbations start to grow: [4] as we mentioned above, relativistic particles escape from gravitational potentials, and hence the gravitational potentials are smeared out instead of becoming deeper. Before becoming nonrelativistic, WDM particles travel a distance of the order of the horizon size; the WDM perturbations are therefore suppressed at those scales. The horizon size at the time $t_{\rm nr}$ when $T \sim m$ is of the order of

$$l(t_{\rm nr}) \simeq H^{-1}(T\sim m) = \frac{M_{\rm Pl}}{\sqrt{g_*}T^2} \sim \frac{M_{\rm Pl}}{\sqrt{g_*}m^2}\,,$$

where $H(T)$ is the Hubble parameter at the temperature $T$. Due to the expansion of the Universe, the corresponding length at present is

$$l_0 = l(t_{\rm nr})\,\frac{a_0}{a(t_{\rm nr})} \sim l(t_{\rm nr})\,\frac{T}{T_0} \sim \frac{M_{\rm Pl}}{mT_0}\,, \qquad (14)$$

where we neglected the (rather weak) dependence on $g_*$. Hence, in the WDM scenario, objects smaller than $l_0$ in size are less abundant than in the CDM case. We point out that $l_0$ refers to the size of the perturbation as if it were in the linear regime; in other words, this is the size of the region from which matter clumps into a compact object.

The present size of a dwarf galaxy is a few kpc, and the density is about $10^6$ of the average density in the Universe. Hence, the size $l_0$ for these objects is around

[3] We omit irrelevant numerical factors in the arguments of the logarithm here and in (12).

[4] The situation is in fact somewhat more complicated, but this is irrelevant for our estimates.

100 kpc $\simeq 3 \times 10^{23}$ cm. From (14), requiring that perturbations of this size, but not much larger, be suppressed, we obtain the estimate for $m$ of a few keV for the mass of WDM particles. In fact, such a small mass of WDM particles is most likely inconsistent with the data. Hydrogen in the Universe was reionized at the redshift $z \sim 10$; this property leaves imprints on the CMB temperature anisotropy and, most notably, CMB polarization, and these imprints have been detected [16]. The reionization is attributed to the formation, burning, and explosions of the first stars, which are believed to have formed in dark matter halos with a mass around $10^5 M_{\odot}$. The initial comoving size of these halos is roughly $l_0 \sim 10$ kpc, and hence perturbations of this spatial size must not be strongly suppressed. According to (14), this implies a bound on the WDM particle mass around (a few) $\times 10$ keV.

Among WDM candidates, the light gravitino is probably the best. The gravitino mass is of the order of

$$m_{3/2} \simeq \frac{F}{M_{\mathrm{Pl}}} \,,$$

where $\sqrt{F}$ is the supersymmetry breaking scale. Hence, gravitino masses are in the right ballpark for rather low supersymmetry breaking scales, $\sqrt{F} \sim 10^7$ GeV. This is the case, for instance, in the gauge mediation scenario [31]. With such a low mass, the gravitino lifetime is much longer than the age of the Universe, and from this standpoint, gravitinos can indeed serve as dark matter particles. For what follows, the important parameters are the widths of decays of other superpartners into gravitino and Standard Model particles. These are approximately

$$\Gamma_{\tilde{\mathrm{S}}} \simeq \frac{M_{\tilde{\mathrm{S}}}^5}{F^2} \simeq \frac{M_{\tilde{\mathrm{S}}}^5}{m_{3/2}^2 M_{\mathrm{Pl}}^2} \,, \qquad (15)$$

where $M_{\tilde{\mathrm{S}}}$ is the superpartner mass.

One mechanism of gravitino production in the early Universe is decays of other superpartners. The gravitino interacts with everything else so weakly that, once produced, it moves freely, without interacting with the cosmic plasma. At production, gravitinos are relativistic; hence, they are indeed *warm* dark matter candidates. We assume that production in decays is the dominant mechanism and consider under what circumstances the present mass density of gravitinos coincides with that of dark matter (see, e.g., Ref. [32] for the details).

The rate of gravitino production in decays of $\tilde{\mathrm{S}}$-type superpartners in the early Universe is

$$\frac{\mathrm{d}(n_{3/2}/s)}{\mathrm{d}t} = \frac{n_{\tilde{\mathrm{S}}}}{s} \, \Gamma_{\tilde{\mathrm{S}}} \,,$$

where $n_{3/2}$ and $n_{\tilde{\mathrm{S}}}$ are number densities of gravitinos and superpartners and $s$ is the entropy density. For superpartners in thermal equilibrium, we have $n_{\tilde{\mathrm{S}}}/s = const \sim g_*^{-1}$ for $T \gtrsim M_{\tilde{\mathrm{S}}}$, and $n_{\tilde{\mathrm{S}}}/s \propto \exp(-M_{\tilde{\mathrm{S}}}/T)$ at $T \ll M_{\tilde{\mathrm{S}}}$. Hence, the production is most efficient at $T \sim M_{\tilde{\mathrm{S}}}$, when the number density of the superpartners is still large, while the Universe expands most slowly. The density of gravitinos produced in decays of $\tilde{\mathrm{S}}$s is therefore given by

$$\frac{n_{3/2}}{s} \simeq \left( \frac{\mathrm{d}(n_{3/2}/s)}{\mathrm{d}t} \, H^{-1} \right)_{T \sim M_{\tilde{\mathrm{S}}}} \simeq \frac{\Gamma_{\tilde{\mathrm{S}}}}{g_*} \, H^{-1}(T \sim M_{\tilde{\mathrm{S}}})$$
$$\simeq \frac{1}{g_*} \, \frac{M_{\tilde{\mathrm{S}}}^5}{m_{3/2}^2 M_{\mathrm{Pl}}^2} \, \frac{M_{\mathrm{Pl}}}{g_*^{1/2} M_{\tilde{\mathrm{S}}}^2} \,.$$

This gives the present mass-to-entropy ratio

$$\frac{m_{3/2} n_{3/2}}{s} \simeq \sum_{\tilde{\mathrm{S}}} \frac{M_{\tilde{\mathrm{S}}}^3}{g_*^{3/2} M_{\mathrm{Pl}} m_{3/2}} \,, \qquad (16)$$

where the sum ranges over all superpartner species *that have ever been relativistic in thermal equilibrium*. The correct value (7) is obtained for gravitino masses in range (10) at

$$M_{\tilde{\mathrm{S}}} = 100 - 300 \ \mathrm{GeV} \,. \qquad (17)$$

Hence, the scenario with the gravitino as the warm dark matter particle requires light superpartners, which are to be discovered at the LHC.

A few comments are in order. First, the decay of superpartners is not the only mechanism of gravitino production: gravitinos may also be produced in the scattering of superpartners. To avoid overproduction of gravitinos in the latter processes, we have to assume that the maximum temperature in the Universe (reached, say, after the post-inflationary reheating stage) is quite low, $T_{\max} \sim 1 - 10$ TeV. This is not a particularly plausible assumption, but it is consistent with everything else in cosmology and can indeed be realized in some models of inflation. Second, existing constraints on masses of strongly interacting superpartners (gluinos and squarks of the first and second generations) suggest that their masses exceed (17). Hence, these particles should not contribute to the sum in (16); otherwise, WDM gravitinos would be overproduced. This is possible if the masses of squarks and gluinos are larger than $T_{\max}$, such that they were never abundant in the early Universe. Third, the decay into gravitino and Standard Model particles is the only decay channel for the next-to-lightest superpartner (NLSP). Hence, with the estimate for the total width of the NLSP given by (15), we have

$$c\tau_{\mathrm{NLSP}} = 1 \ \mathrm{mm} - 100 \ \mathrm{m}$$

for $m_{3/2} = 1 - 10$ keV and $M_{\mathrm{NLSP}} = 100 - 300$ GeV. The NLSP should therefore either be visible in a detector or fly through it.

Needless to say, the outlined scenario is much more contrived than the WIMP option. It is reassuring, however, that it will be ruled out or confirmed at the LHC.

Finally, gravitinos can be much heavier than 100 keV, and still be the lightest supersymmetric particles. Then they serve as CDM candidates. Obviously, direct detection of CDM particles is hopeless in that case.

### 2.3 Discussion

If dark matter particles are indeed WIMPs, and the relevant energy scale is about 1 TeV, then the Hot Big Bang theory will be probed experimentally up to temperatures of $(10 - 100)$ GeV and down to the age of $10^{-9} - 10^{-11}$s in the relatively near future (which is to be compared with 1 MeV and 1 s accessible today through Big Bang Nucleosynthesis). With the microscopic physics to become known from collider experiments, the WIMP density will be reliably calculated and checked against data from observational cosmology. Thus, the WIMP scenario offers a window to a very early stage of the evolution of the Universe.

If dark matter is warm and its particles are gravitinos, then the prospect of quantitatively accessing such an early stage of the cosmological evolution is not so bright: it would be very difficult, if at all possible, to assess the gravitino mass experimentally; furthermore, the present gravitino mass density depends on an unknown reheat temperature $T_{\max}$. On the other hand, if this scenario is realized in Nature, the whole picture of the early Universe will be quite different from what we think today is the most likely early cosmology. Indeed, the gravitino scenario requires a low reheat temperature, which in turn calls for rather exotic mechanisms of inflation, etc.

The mechanisms discussed here are by no means the only ones capable of producing dark matter, and WIMPs and gravitinos are by no means the only candidates for dark matter particles. Other dark matter candidates include axions, sterile neutrinos, Q-balls, very heavy relics produced towards the end of inflation, and so on. Hence, even though there are grounds to hope that the dark matter problem will be solved by the LHC, there is no guarantee at all.

## 3. Baryon asymmetry of the Universe

In the present Universe, there are baryons and almost no antibaryons. The number density of baryons today is characterized by the ratio $\eta$ [see Eqn (3)]. In the early Universe, the appropriate quantity is

$$\Delta_{\mathrm{B}} = \frac{n_{\mathrm{B}} - n_{\bar{\mathrm{B}}}}{s} \,,$$

where $n_{\bar{\mathrm{B}}}$ is the number density of antibaryons and $s$ is the entropy density. If the baryon number is conserved and the Universe expands adiabatically, then $\Delta_{\mathrm{B}}$ is constant and its value is equal to $\eta$ up to a numerical factor,

$$\Delta_{\mathrm{B}} \approx 0.8 \times 10^{-10} \,.$$

At early times, at temperatures well above 100 MeV, the cosmic plasma contained many quark–antiquark pairs, whose number density was of the order of the entropy density, $n_{\mathrm{q}} + n_{\bar{\mathrm{q}}} \sim s$, while the baryon number density was related to the densities of quarks and antiquarks as $n_{\mathrm{B}} - n_{\bar{\mathrm{B}}} = (1/3)(n_{\mathrm{q}} - n_{\bar{\mathrm{q}}})$ (the baryon number of a quark is $1/3$). Hence, in terms of quantities characterizing the very early epoch, the baryon asymmetry can be expressed as

$$\Delta_{\mathrm{B}} \sim \frac{n_{\mathrm{q}} - n_{\bar{\mathrm{q}}}}{n_{\mathrm{q}} + n_{\bar{\mathrm{q}}}} \,.$$

We see that there was one extra quark for about every 10 billion quark–antiquark pairs. It is this tiny excess that is responsible for all the baryonic matter in the present Universe.

There is no logical contradiction in supposing that the tiny excess of quarks over antiquarks was built in as an initial condition. But this is not at all satisfactory for a physicist. Furthermore, the inflationary scenario does not provide such an initial condition for the hot stage; rather, inflation theory predicts that the Universe was baryon-symmetric immediately after inflation. Hence, we would like to explain the baryon asymmetry dynamically.

The baryon asymmetry can be generated from the initially symmetric state only if three necessary conditions, called Sakharov's conditions, are satisfied. These are:

(i) baryon number nonconservation;

(ii) *C*- and *CP*-violation; and

(iii) deviation from thermal equilibrium.

All three conditions are easily understood. (i) If the baryon number were conserved, and the initial net baryon number in the Universe was zero, the Universe today would still be symmetric. (ii) If C or CP were conserved, the rates of reactions with particles would be the same as the rates of reactions with antiparticles, and no excess of quarks over antiquarks would be generated. (iii) Thermal equilibrium is the most symmetric state of a system. If the baryon number were the only relevant quantum number, it would be washed out, rather than generated, as the system approaches thermal equilibrium. In fact, the baryon number is *not* the only relevant quantum number, and hence the last point requires qualification, which is not of importance for us.

There are two well-understood mechanisms of baryon number nonconservation. One of them emerges in hypothetical Grand Unified Theories and is due to the exchange by supermassive particles. It is very similar, e.g., to the mechanism of charm nonconservation in weak interactions, which occurs via the exchange by heavy *W*-bosons. The scale of these new, baryon-number-violating interactions is the Grand Unification scale, presumably of the order of $10^{16}$ GeV.

Another mechanism is nonperturbative [33] and is related to the triangle anomaly in the baryonic current. It already exists in the Standard Model and, possibly with slight modifications, operates in all its extensions. The two main features of this mechanism, as applied to the early Universe, are that it is effective over a wide range of temperatures [34], $100\ \mathrm{GeV} < T < 10^{11}$ GeV, and that it preserves $B - L$.

We pause here to discuss the physics behind electroweak baryon and lepton number nonconservation in a little more detail, although still at a qualitative level. The first object to consider is the baryonic current

$$B^{\mu} = \frac{1}{3} \sum_i \bar{q}_i \gamma^{\mu} q_i \,,$$

where the sum ranges over quark flavors. Naively, the baryonic current is conserved, but at the quantum level, its divergence is nonzero due to the triangle anomaly (similar effects go under the name of axial anomaly in the context of QED and QCD),

$$\partial_{\mu} B^{\mu} = \frac{1}{3} \times 3_{\mathrm{colors}} \times 3_{\mathrm{generations}} \times \frac{g^2}{32\pi^2}\, \epsilon^{\mu\nu\lambda\rho} F^{\mathrm{a}}_{\mu\nu} F^{\mathrm{a}}_{\lambda\rho} \,,$$

where $F^{\mathrm{a}}_{\mu\nu}$ and $g$ are the field strength of the $\mathrm{SU}(2)_{\mathrm{W}}$ gauge field and the $\mathrm{SU}(2)_{\mathrm{W}}$ coupling. Likewise, each leptonic current ($\alpha = \mathrm{e}, \mu, \tau$) is anomalous,

$$\partial_{\mu} L^{\mu}_{\alpha} = \frac{g^2}{32\pi^2}\, \epsilon^{\mu\nu\lambda\rho} F^{\mathrm{a}}_{\mu\nu} F^{\mathrm{a}}_{\lambda\rho} \,.$$

A nontrivial fact [35] is that there exist large field fluctuations, $F^{\mathrm{a}}_{\mu\nu}(\mathbf{x}, t) \propto g^{-1}$, for which

$$Q \equiv \int \mathrm{d}^3 x\, \mathrm{d}t\, \frac{g^2}{32\pi^2}\, \epsilon^{\mu\nu\lambda\rho} F^{\mathrm{a}}_{\mu\nu} F^{\mathrm{a}}_{\lambda\rho} \neq 0 \,.$$

Furthermore, for any such fluctuation, the value of $Q$ is an integer.

We now suppose that a fluctuation with a nonvanishing $Q$ has occurred. Then the baryon numbers at the end and beginning of the process are different:

$$B_{\mathrm{fin}} - B_{\mathrm{in}} = \int \mathrm{d}^3 x\, \mathrm{d}t\, \partial_{\mu} B^{\mu} = 3Q \,. \tag{18}$$

Similarly,

$$L_{n,\,\mathrm{fin}} - L_{n,\,\mathrm{in}} = Q\,. \tag{19}$$

This explains the selection rule mentioned above: $B$ is violated and $B - L$ is not.

At zero temperature, the large field fluctuations that induce baryon and lepton number violation are vacuum fluctuations, called instantons [35], which to a certain extent are similar to virtual fields that emerge and disappear in the vacuum of quantum field theory at the perturbative level. The difference is that instantons are *large* field fluctuations. This property results in a suppression of the corresponding probability, and hence the rate of baryon-number-violating processes, by the factor $\exp(-16\pi^2/g^2) \sim 10^{-165}$. On the other hand, at high temperatures, there are large *thermal* fluctuations, 'sphalerons' [36], whose rate is not necessarily small. And, indeed, $B$-violation in the early Universe is rapid compared to the cosmological expansion at sufficiently high temperatures, when

$$\langle\phi\rangle_T < T\,, \tag{20}$$

where $\langle\phi\rangle_T$ is the Higgs expectation value at temperature $T$.

One may wonder how the baryon number may be not conserved even though there are no baryon-number-violating terms in the Lagrangian of the Standard Model. This is discussed in detail, for instance, in Ref. [37]. In any case, it is tempting to use this mechanism of baryon number nonconservation for explaining the baryon asymmetry of the Universe. There are two problems, however. One is that $CP$-violation in the Standard Model is too weak: the CKM mechanism alone is insufficient to generate the realistic value of baryon asymmetry. Hence, we need extra sources of $CP$-violation. Another problem has to do with departure from thermal equilibrium, which is necessary for the generation of baryon asymmetry. At temperatures well above 100 GeV, the electroweak symmetry is restored, the expectation value of $\phi$ is zero,[5] relation (20) holds, and the baryon number nonconservation is rapid compared to the cosmological expansion. At temperatures around 100 GeV, relation (20) may be violated, but the Universe expands very slowly: the cosmological time scale at these temperatures is

$$H^{-1} \sim \frac{M_{\mathrm{Pl}}}{\sqrt{g_*}T^2} \simeq \frac{10^{19}\ \mathrm{GeV}}{10\times(100\ \mathrm{GeV})^2} \sim 10^{-10}\ \mathrm{s}\,, \tag{21}$$

which is very large by the electroweak physics standards. The only way in which a strong departure from thermal equilibrium at these temperatures may occur is through a first-order phase transition.

The property (valid within the perturbation theory only) that the expectation value of the Higgs field is zero at temperatures well above 100 GeV, while it is nonzero in the vacuum, suggests that there may be a phase transition from, crudely speaking, the phase with $\langle\phi\rangle = 0$ to the phase with $\langle\phi\rangle \neq 0$. The situation is pretty subtle here, because $\phi$ is not gauge invariant, and hence cannot serve as an order parameter, and therefore the notion of phases with $\langle\phi\rangle = 0$ and $\langle\phi\rangle \neq 0$ is vague. In fact, neither electroweak theory nor most of its extensions have a gauge-invariant order parameter, and hence there is no real distinction between these 'phases.' This situation is very similar to that in a liquid–vapor system, which does not have an order parameter and may or may not experience the vapor–liquid phase transition as the temperature decreases, depending on other parameters characterizing this system, e.g., pressure. In the Standard Model, the role of such a parameter is played by the Higgs self-coupling $\lambda$ or, in other words, the Higgs boson mass.

Continuing to use the somewhat inexact terminology, the interesting case for us is the first-order phase transition. Here, the effective potential $V_{\mathrm{eff}}(\phi)$ (free energy density as a function of $\phi$) has one minimum at $\phi = 0$ at high temperatures, and the expectation value of the Higgs field is zero. As the temperature decreases, another minimum appears at a finite $\phi$, and then becomes lower than the minimum at $\phi = 0$. However, the probability of the transition from the phase $\phi = 0$ to the phase $\phi \neq 0$ is very small for some time, and the system becomes overcooled.

The first-order phase transition occurs via spontaneous creation of bubbles of the new phase inside the old phase. These bubbles then grow, their walls eventually collide, and the new phase eventually occupies the entire space. The Universe boils. In the cosmological context, this process occurs when the bubble nucleation rate per Hubble time per Hubble volume is approximately 1. The velocity of the bubble wall in the relativistic cosmic plasma is roughly of the order of the speed of light (in fact, it is somewhat smaller, from $0.1c$ to $0.01c$), simply because there are no relevant dimensionless parameters characterizing the system. Hence, the bubbles grow large before their walls collide: their size at collision is roughly comparable to the Hubble size. While the bubble is microscopic at nucleation — its size is dictated by the electroweak scale and is roughly $(100\ \mathrm{GeV})^{-1} \approx 10^{-16}$ cm — its size is macroscopic at the time the walls collide, $0.1H^{-1} \sim 1$ mm, as follows from (21). Clearly, boiling is a highly nonequilibrium process, and it may be hoped that the baryon asymmetry may be generated at that time. And, indeed, there exist mechanisms for the generation of baryon asymmetry, which have to do with interactions of quarks and leptons with moving bubble walls. The value of the resulting baryon asymmetry may well be around $10^{-10}$, as required by observations, if there is enough $CP$-violation in the theory.

A necessary condition for the electroweak generation of the baryon asymmetry is that inequality (20) be violated *immediately after* the phase transition. Indeed, in the opposite case, electroweak baryon number violating processes are fast after the transition, and the baryon asymmetry generated during the transition is washed out afterwards. Hence, the phase transition must be of a strong enough first order. This is *not* the case in the Standard Model. To understand why, and to see in which extensions of the Standard Model the transition may be strong enough, we consider the effective potential in some detail. At zero temperature, the Higgs potential has the standard form

$$V(\phi) = -\frac{m^2}{2}\,|\phi|^2 + \frac{\lambda}{4}\,|\phi|^4\,,$$

where

$$|\phi| \equiv \left(\phi^\dagger\phi\right)^{1/2} \tag{22}$$

is the magnitude of the Higgs doublet $\phi$, $m^2 = \lambda v^2$, and $v = 247$ GeV is the Higgs expectation value in the vacuum.

[5] There are subtleties here; see below.

The Higgs boson mass is related to it as

$$m_{\rm H} = \sqrt{2\lambda}\,v\,. \tag{23}$$

In the leading order of the perturbation theory, finite-temperature effects modify the effective potential into

$$V_{\rm eff}(\phi, T) = \frac{\alpha}{2}|\phi|^2 - \frac{\beta}{3}\,T|\phi|^3 + \frac{\lambda}{4}\,|\phi|^4\,, \tag{24}$$

with $\alpha(T) = -m^2 + \hat{g}^2 T^2$ and $\beta = \tilde{g}^3/(2\pi)$, where $\hat{g}^2$ is a positive linear combination of squares of the coupling constants of all fields to the Higgs field (in the Standard Model, a linear combination of $g^2, g'^2$, and $y_i^2$, where $g$ and $g'$ are gauge couplings and $y_i$ are Yukawa couplings), while $\tilde{g}^3$ is a positive linear combination of cubes of coupling constants of all *bosonic* fields to the Higgs field. In the Standard Model, $\beta$ is a linear combination of $g^3$ and $g'^3$, i.e., a linear combination of $M_{\rm W}^3/v^3$ and $M_{\rm Z}^3/v^3$,

$$\beta = \frac{1}{2\pi}\,\frac{2M_{\rm W}^3 + M_{\rm Z}^3}{v^3}\,. \tag{25}$$

The cubic term in (24) is rather peculiar: in view of (22), it is not analytic in the original Higgs field $\phi$. Yet this term is crucial for the first-order phase transition: for $\beta = 0$, the phase transition would be of the second order. The origin of the nonanalytic cubic term can be traced back to the enhancement of the Bose–Einstein thermal distribution at low momenta, $p, m \ll T$,

$$f_{\rm Bose}(p) = \frac{1}{\exp\left(\sqrt{p^2 + m_{\rm b}^2}/T\right) - 1} \simeq \frac{T}{\sqrt{p^2 + m_{\rm b}^2}}\,,$$

where $m_{\rm b} \simeq g_{\rm b}|\phi|$ is the mass of the boson b that is generated due to the nonvanishing Higgs field and $g_{\rm b}$ is the coupling constant of the field $b$ to the Higgs field. Clearly, at $p \ll g_{\rm b}|\phi|$ the distribution function is nonanalytic in $\phi$,

$$f_{\rm Bose}(p) \simeq \frac{T}{g_{\rm b}|\phi|}\,.$$

This nonanalyticity gives rise to the nonanalytic cubic term in the effective potential. Importantly, the Fermi–Dirac distribution

$$f_{\rm Fermi}(p) = \frac{1}{\exp\left(\sqrt{p^2 + m_{\rm f}^2}/T\right) + 1}\,,$$

is analytic in $m_{\rm f}^2$, and hence in $\phi^\dagger\phi$, and therefore fermions do not contribute to the cubic term.

With the cubic term in the effective potential, the phase transition is indeed of the first order (within the approximation considered here): at high temperatures, the coefficient $\alpha$ is positive and large, and there is one minimum of the effective potential at $\phi = 0$, while at intermediate temperatures, $\alpha$ is small but still positive, and therefore there are two minima. The phase transition occurs at $\alpha \approx 0$; at that instant,

$$V_{\rm eff}(\phi, T) \approx -\frac{\beta}{3}\,T|\phi|^3 + \frac{\lambda}{4}\,|\phi|^4\,.$$

We find from this expression that immediately after the phase transition, the minimum of $V_{\rm eff}$ is at

$$\phi \simeq \frac{\beta}{\lambda} = \frac{\tilde{g}^3 T}{\lambda}\,.$$

Hence, the necessary condition for successful electroweak baryogenesis, $\phi > T$, translates into

$$\beta > \lambda\,. \tag{26}$$

According to (23), $\lambda$ is proportional to $m_{\rm H}^2$, whereas in the Standard Model, $\beta$ is proportional to $2M_{\rm W}^3 + M_Z^3$. Therefore, relation (26) holds for small Higgs boson masses only; in the Standard Model, using (23) and (25) shows that that would happen for $m_{\rm H} < 50$ GeV, which is ruled out.[6]

This discussion indicates a possible way to make the electroweak phase transition strong. We need the existence of new bosonic fields that have large enough couplings to the Higgs field(s), and hence make large contributions to $\beta$. To produce an effect on the dynamics of the transition, the new bosons must be present in the cosmic plasma at the transition temperature, $T \sim 100$ GeV, and therefore their masses should not be too high, $M \lesssim 300$ GeV. In supersymmetric extensions of the Standard Model, the natural candidate is the scalar partner of the top quark, whose Yukawa coupling to the Higgs field is the same as that of the top quark, that is, large. The light stop scenario for electroweak baryogenesis indeed works, as has been shown by the detailed analysis in Refs [41–43].

Yet another issue is *CP*-violation, which has to be strong enough for the successful electroweak baryogenesis. Because the asymmetry is generated in the interactions of quarks and leptons (and their superpartners in supersymmetric extensions) with the bubble walls, *CP*-violation must occur at the walls. We now recall that the walls are made of the Higgs field(s). This points towards the necessity of *CP*-violation in the Higgs sector, which may only be the case in a theory with more than one Higgs fields.

To summarize, electroweak baryogenesis requires a considerable extension of the Standard Model, with masses of new particles in the range $100-300$ GeV. Hence, this mechanism will definitely be ruled out or confirmed by the LHC. We emphasize, however, that electroweak baryogenesis is not the only option: an elegant and well-motivated competitor is leptogenesis, and there are several other mechanisms that have been proposed in the literature.

## 4. Concluding remarks

The ideas we have discussed may not be the right ones: we can only hypothesize on physics beyond the Standard Model and its role in the early Universe. The TeV-scale physics may be dramatically different from the physics we are used to. As an example, it cannot be ruled out that TeV is not only an electroweak scale but also a gravitational scale. This is the case in models with large extra dimensions, in which the Planck scale is related to the fundamental gravity scale in a way that involves the volume of extra dimensions, and hence the fundamental scale can be much below $M_{\rm Pl}$ (for a review, see, e.g., Ref. [44]).

If the LHC finds that, indeed, the fundamental gravity scale is in the TeV range, this will have very profound

[6] In fact, in the Standard Model with $m_{\rm H} > 114$ GeV, there is no phase transition at all [38–40]; the electroweak transition is a smooth crossover instead. This fact is not visible from expression (24), but that expression is the lowest-order perturbative result, while the perturbation theory is not applicable in describing the transition in the Standard Model with large $m_{\rm H}$.

consequences for both microscopic physics and cosmology. On the microscopic physics side, this would allow studying quantum gravity and its high-energy extension (possibly string theory) at colliders, while on the cosmological side, the entire picture of the early Universe would have to be revised. Inflation, if any, would have to occur either at low energy density or in the state of strong quantum gravity effects. The highest temperatures in the usual expansion history would be at most in the TeV range, such that dark matter and baryon asymmetry would have to be generated either below TeV temperatures or in the quantum gravity mode. Even more intriguing would be the study of the quantum gravity cosmological epoch, with hints from colliders gradually coming in. This, probably, is too optimistic an outlook to be realistic.

It is more likely that the LHC will find something entirely new, something theorists have not thought about, or, conversely, find so little that we will have to get serious about the anthropic principle. In any case, the LHC results will definitely change the landscape of fundamental physics, cosmology included.

## References


1. Gorbunov D S, Rubakov V A *Vvedenie v Teoriyu Rannei Vselennoi. Teoriya Goryachego Bol'shogo Vzryva* (Introduction to the Theory of the Early Universe. Hot Big Bang Theory) (Moscow: URSS, 2008) [Translated into English (Singapore: World Scientific, 2011)]
2. Brandenberger R H, hep-ph/9701276
3. Freedman W L, Turner M S *Rev. Mod. Phys.* **75** 1433 (2003); astro-ph/0308418
4. Rubakov V *PoS* (RTN2005) 003 (2005); http://pos.sissa.it/archive/conferences/019/003/RTN2005_003.pdf
5. Peacock J A, in *Proc. 1998 European School of High-Energy Physics, St. Andrews, Scotland, 23 Aug. – 5 Sep. 1998*
6. Shaposhnikov M, in *Proc. 2000 European School of High-Energy Physics, Caramulo, Portugal, 20 Aug. – 2 Sep. 2000*
7. Tkachev I I, in *Proc. 2003 European School on High-Energy Physics, Tsakhkadzor, Armenia, 24 Aug. – 6 Sep. 2003*; hep-ph/0405168
8. Jungman G, Kamionkowski M, Griest K *Phys. Rep.* **267** 195 (1996); hep-ph/9506380
9. Bottino A, Fornengo N, hep-ph/9904469
10. Olive K A, astro-ph/0301505
11. Bertone G, Hooper D, Silk J *Phys. Rep.* **405** 279 (2005); hep-ph/0404175
12. Dolgov A D, hep-ph/9707419
13. Rubakov V A, Shaposhnikov M E *Usp. Fiz. Nauk* **166** 493 (1996) [*Phys. Usp.* **39** 461 (1996)]; hep-ph/9603208
14. Riotto A, Trodden M *Annu. Rev. Nucl. Part. Sci.* **49** 35 (1999); hep-ph/9901362
15. Trodden M *Rev. Mod. Phys.* **71** 1463 (1999); hep-ph/9803479
16. Komatsu E et al. *Astrophys. J. Suppl.* **192** 18 (2011); arXiv:1001.4538
17. Seljak U et al. *Phys. Rev. D* **71** 103515 (2005); astro-ph/0407372
18. Komatsu E et al. *Astrophys. J. Suppl.* **180** 330 (2009); arXiv:0803.0547
19. Malinovsky A M et al. *Pis'ma Astron. Zh.* **34** 490 (2008) [*Astron. Lett.* **34** 445 (2008)]
20. Sahni V, Starobinsky A *Int. J. Mod. Phys. D* **9** 373 (2000); astro-ph/9904398
21. Weinberg S, astro-ph/0005265
22. Chernin A D *Usp. Fiz. Nauk* **171** 1153 (2001) [*Phys. Usp.* **44** 1099 (2001)]
23. Padmanabhan T *Phys. Rep.* **380** 235 (2003); hep-th/0212290
24. Peebles P J E, Ratra B *Rev. Mod. Phys.* **75** 559 (2003); astro-ph/0207347
25. Sahni V *Lect. Notes Phys.* **653** 141 (2004); astro-ph/0403324
26. Copeland E J, Sami M, Tsujikawa S *Int. J. Mod. Phys. D* **15** 1753 (2006); hep-th/0603057
27. Vysotskii M I, Nevzorov R B *Usp. Fiz. Nauk* **171** 939 (2001) [*Phys. Usp.* **44** 919 (2001)]
28. Krasnikov N V, Matveev V A *Usp. Fiz. Nauk* **174** 697 (2004) [*Phys. Usp.* **47** 643 (2004)]; hep-ph/0309200
29. Tikhonov A V, Klypin A *Mon. Not. R. Astron. Soc.* **395** 1915 (2009); arXiv:0807.0924
30. Salucci P, arXiv:1008.4344
31. Gorbunov D S, Dubovskii S L, Troitskii S V *Usp. Fiz. Nauk* **169** 705 (1999) [*Phys. Usp.* **42** 623 (1999)]; Dubovskii S L, Gorbunov D S, Troitskii S V, hep-ph/9905466
32. Gorbunov D, Khmelnitsky A, Rubakov V *JHEP* (12) 055 (2008); arXiv:0805.2836
33. 't Hooft G *Phys. Rev. Lett.* **37** 8 (1976)
34. Kuzmin V A, Rubakov V A, Shaposhnikov M E *Phys. Lett. B* **155** 36 (1985)
35. Belavin A A et al. *Phys. Lett. B* **59** 85 (1975)
36. Klinkhamer F R, Manton N S *Phys. Rev. D* **30** 2212 (1984)
37. Rubakov V *Klassicheskie Kalibrovochnye Polya* (Classical Theory of Gauge Fields) (Moscow: URSS, 1999) [Translated into English (Princeton, NJ: Princeton Univ. Press, 2002)]
38. Kajantie K et al. *Phys. Rev. Lett.* **77** 2887 (1996); hep-ph/9605288
39. Karsch F et al. *Nucl. Phys. B* **474** 217 (1996); hep-lat/9603004
40. Karsch F et al. *Nucl. Phys. B Proc. Suppl.* **53** 623 (1997); hep-lat/9608087
41. Carena M, Quirós M, Wagner C E M *Phys. Lett. B* **380** 81 (1996); hep-ph/9603420
42. Carena M et al. *Nucl. Phys. B* **503** 387 (1997); hep-ph/9702409
43. Carena M et al. *Nucl. Phys. B* **650** 24 (2003); hep-ph/0208043
44. Rubakov V A *Usp. Fiz. Nauk* **171** 913 (2001) [*Phys. Usp.* **44** 871 (2001)]; hep-ph/0104152


PACS numbers: **04.65.+e**, **11.10.–**, 11.30.Ly
DOI: 10.3367/UFNe.0181.201106g.0665

# V L Ginzburg and higher-spin fields

M A Vasiliev

## 1. Higher-spin fields yesterday and today

Relativistic fields are characterized by two types of quantum numbers: mass $m \geqslant 0$, and spin $s = 0, 1/2, 1, 3/2, \ldots, \infty$. To date, only two types of particles have been observed experimentally: particles of spin $s = 1/2$ — that is, $\mathrm{e}, \nu, \mu, \mathrm{u}, \mathrm{d} \ldots$, which describe matter fields, and those of spin $s = 1$, like photons, gluons, and W and Z bosons, which serve as mediators of interactions.

The main goal of the LHC is to find the hypothetical particle of spin 0, the Higgs boson H. Massless particles of spin 2 (graviton) and spin 3/2 (gravitino) remain to be discovered, although gauge theories related to them, namely gravity and supergravity, are well known, at least at the classical level.

The theory of free fields of any spin and mass is perfectly defined at the Lagrangian level. A nontrivial and highly interesting problem arises once the question of the structure of the theory of interacting fields of spins $s > 2$ is raised.

The foundation of the theory of free higher-spin fields was laid in the classical work of Dirac [1] and Fierz and Pauli [2]. The history of the development of higher-spin theory can be roughly split into two stages. Before the creation of supergravity [3], i.e., approximately from 1936 till 1976, the main goal was to describe higher-spin resonances. During this period, the main focus was on the study of massive particles in four dimensional spacetime. After the creation of super-

**M A Vasiliev** Tamm Theory Department, Lebedev Physical Institute, Russian Academy of Sciences, Moscow, Russian Federation
E-mail: vasiliev@lpi.ru

*Uspekhi Fizicheskikh Nauk* **181** (6) 665 – 672 (2011)
DOI: 10.3367/UFNr.0181.201106g.0665
Translated by M A Vasiliev; edited by A Radzig