Primordial background of cosmological axions

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Combined analysis of astrophysical and cosmological consequences of axion models permits rigorous selection of such models and identifications of the most self-consistent archion model based on models of quantum flavordynamics with a low ($\sim 10^6$ GeV) scale of generation symmetry breaking. In models of an invisible axion with a low ($< 10^8$ GeV) scale of the axion interaction there is predicted to be a primordial thermal axion background due to the establishment of equilibrium between the axions and the plasma in the early universe after a QCD phase transition due to the reaction $\pi N \to aN$. The spectrum of cosmic archions predicted in hierarchical decay models of the formation of structure in the universe is calculated. This will increase the possibilities for detailed testing of models of quantum flavordynamics with a low scale of the generation symmetry breaking.

1. INTRODUCTION

The existence of the axion[1-3] is a very stable prediction of the theory associated with the solution of the problem of the violation of *CP* invariance in QCD. However, the combination of the negative results of experimental searches for axions with astrophysical and cosmological bounds on their parameters leaves a fairly small choice of admissible models of an invisible axion, ruling out the possibility of direct detection of axions in experiments using existing and planned accelerators and leaving a hope merely of detection of fluxes of cosmological axions.[3,4] The aim of the present paper is to discuss possible manifestations of axions in the early universe and to analyze the predicted parameters of a primordial axion background as functions of the axion model that is chosen. In particular, an essentially new possibility for the existence of a primordial thermal axion background in models with a low (10^6—10^8 GeV) scale of the axion interaction is noted.

The archion model[5,6] occupies a special position. This model is attractive by virtue of being physically well founded in the framework of a gauge theory of broken symmetry of the quark and lepton generations, relating the axion hypothesis to the observed mass hierarchy and mixing of the fermion generations and the predicted mass spectrum of unstable neutrinos. A cosmological archion background responsible for the dark mass of the contemporary universe is one of the most important cosmological predictions of this theory, permitting not only an astronomical but also, in principle, a laboratory verification.

2. BOUNDS ON AXION MODELS

The complicated structure of the vacuum in QCD, due to instanton effects, gives rise to the appearance of the so-called θ term in the Lagrangian of this theory:

$$\Delta L = \frac{g^2}{16\pi}\,\theta_{QCD} G_{\mu\nu}\tilde{G}^{\mu\nu}, \tag{1}$$

Where $\theta_{\rm QCD}$ is an arbitrary constant that, in general, need not be small, g is the QCD gauge constant, $G_{\mu\nu}$ is the field-strength tensor of the gluon field, and $\tilde{G}_{\mu\nu} = \varepsilon_{\mu\nu\lambda\rho} G^{\lambda\rho}$.

The θ term violates *P* and *CP* invariance. In strong interactions, the effects of such violation have not been observed, and the negative results of the search for a dipole moment of the neutron, $d_n < 10^{-25}$ $e\cdot$cm, yields an upper bound on the constant, $\theta < 10^{-9}$. It is difficult to solve the *CP*-violation problem in strong interactions because even for a vanishing bare constant, $\theta_{\rm QCD} = 0$, an effective θ term can arise after symmetry breaking of the electroweak interaction, owing to the complexity of the mass matrices of the quarks in the standard model on the transition (by means of unitary transformations) to the physical basis of the quark fields, in which the mass matrices are diagonal.

As a consequence, the effective θ in the Lagrangian (1) is a sum of two terms of different natures:

$$\theta = \theta_{QCD} + \theta_{QFD}, \tag{2}$$

where

$$\theta_{QFD} = \det \hat{m}. \tag{3}$$

in which $\hat{m}$ is the total mass matrix, including all the colored fermions of the theory (both quarks and all possible hypothetical colored fermions). The different natures of the terms in (2), which, in general, are not small, do not allow us to expect *a priori* their cancellation to the level $\theta < 10^{-9}$. This is the essence of one of the "naturalness" problems in elementary-particle theory. Its solution is associated with the existence of a mechanism of natural suppression of θ, for which Peccei and Quinn[1] proposed the existence of an additional global chiral symmetry $U(1)_{\rm PQ}$ with uncompensated color anomaly. If such a symmetry is not broken—this requires the existence of at least one massless colored fermion, for example, a u quark[2]—then all the θ vacua are equivalent to the vacuum with $\theta = 0$. In other words, the θ term can be eliminated from the Lagrangian by a chiral phase transformation of the field of the u quark. Although the possibility of zero mass of the u quark is not, in general, excluded theoretically,[7] it lacks a fundamental justification in the standard model and encounters serious difficulties in a comparison with the results of current algebra and the PCAC hypothesis. However, the problem of the θ term can be solved natu-

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rally even in the case of broken $U(1)_{PQ}$ symmetry. In this case, the constant θ acquires the dynamical significance of the amplitude of the pseudo-Goldstone field of the broken $U(1)_{PQ}$ symmetry (which, in its turn, is explicitly broken by the color anomaly), so that $\langle\theta\rangle_{vac}$ automatically becomes zero, ensuring exact mutual compensation of the terms in (2). Such a pseudo-Goldstone field is called the axion.

The most important parameter that determines the axion properties is the energy scale F_a of the breaking of the $U(1)_{PQ}$ symmetry. This parameter determines both the mass of the axion ($m_a \propto F_a^{-1}$) and the strength of its interaction with the fermions ($g \propto F_a^{-1}$) and with the gauge bosons ($C \propto F_a^{-1}$), in particular, photons. The last of these interactions determines the lifetime of the axion with respect to decay into 2γ.

The simplest form of axion model (Weinberg–Wilczek axion[2]), in which the scale F_a is equal to the scale V of the electroweak symmetry breaking, was eliminated by a combination of experimental and astrophysical bounds,[8] which led to a very high lower bound on the scale: $F_a \gg V$. Thus, the scale F_a must be associated with a new fundamental high energy scale of the theory, for which the axion is invisible.

In all models of such type, the axion arises as a Goldstone boson corresponding to the phase of the complex $SU(2)\times U(1)$ scalar singlet, and generation of a coupling of the type $aF\tilde{F}$ occurs after the breaking of the $U(1)_{PQ}$ symmetry by mechanisms which ensure the presence of a $U(1)_{PQ}$–$SU(3)_C$–$SU(3)_C$ anomaly. In the most general case, the Lagrangian of the interaction of the axion with the fermions (quarks and leptons) and photons has the form

$$L = a g_{\alpha\beta}\bar{f}_\alpha(\sin\theta_{\alpha\beta} + i\cos\theta_{\alpha\beta}\gamma^5)f_\beta + C_{a\gamma\gamma}aF_{\mu\nu}\tilde{F}^{\mu\nu}. \qquad (4)$$

where $\alpha, \beta = 1, 2, 3$ are the indices of the generations of fermions f, and the constants $g \propto F_a^{-1}$, θ, and $C \propto F_a^{-1}$ depend on the choice of the axion model.

The general relation between the axion mass m_a and the scale F_a is given by

$$m_a = A_a m_\pi f_\pi / F_a, \qquad (5)$$

where m_π (f_π) are the pion mass and decay constant, and the parameter A_a depends on the choice of the axion model.

By a general theorem,[9] the diagonal couplings of the axion in Eq. (4) can only be pseudoscalar ($\theta_{\alpha\alpha} = 0$), so that there are no axionic long-range interactions if effects of higher orders are ignored. Depending on the choice of the model, the nondiagonal axion couplings, the existence of which cannot be ruled out, may be either scalar or pseudoscalar, or a combination of them ($\theta_{\alpha\beta} \neq 0$).

There are three main models of an invisible axion, which we shall now discuss.

1. The Dine–Fischler–Srednicki–Zhitnitskiĭ (DFSZ) axion.[10] This model retains the set of fermion fields (quark and lepton generations) of the standard axion, and additional fields are introduced only in the Higgs sector of the theory (the number of Higgs doublets φ is increased, and a singlet σ is added). The Lagrangian of the Yukawa interactions in the DFSZ model has the form

$$L_{Yuk} = g_{\alpha\beta}^{(f)}\bar{f}_{L\alpha}f_{R\beta}\varphi_f + \text{H.c.}, \qquad (6)$$

where $f_{L(R)}$ are the left (respectively, right) components of the quark and lepton fields, $f = u, d, e$; $\alpha, \beta = 1, 2, \ldots, N_g$ are the indices of the generations; φ_f are Higgs doublets $\varphi_u = \begin{pmatrix}\varphi_u^+\\ \varphi_u^0\end{pmatrix}$, $\langle\varphi_u^0\rangle = V_u/\sqrt{2}$, $\varphi_d = \varphi_l = \begin{pmatrix}\varphi_d^+\\ \varphi_d^0\end{pmatrix}$, $\langle\varphi_d^0\rangle = V_d\sqrt{2}$, $(V_u^2 + V_d^2)^{1/2} = V = (\sqrt{2}G_F)^{-1/2}$. In the most general case, the doublets φ_d and φ_c can be different (in particular, the choice $\varphi_e = \tilde{\varphi}_u$ is possible). However, the choice of these doublets given here is dictated by inclusion of the model in a grand-unification scheme. The Lagrangian (6) is invariant with respect to global chiral phase transformations of the symmetry group:

$$\begin{aligned} f_{L\alpha} &\to \exp(i\omega)f_{L\alpha},\\ f_{R\alpha} &\to \exp(-i\omega)f_{R\alpha},\\ \varphi_f &\to \exp(2i\omega)\varphi_f,\\ \sigma &\to \exp(-2i\omega)\sigma. \end{aligned} \qquad (7)$$

this symmetry being ensured by the presence, in the Higgs potential, of the term

$$L_H = \lambda\varphi_u\varphi_d\sigma^2 + \text{H.c.} \qquad (8)$$

Under the assumption $V \ll \langle\sigma\rangle = \tilde{V}/\sqrt{2}$, the axion field is mainly determined by the phase of the field σ, and the admixture of the phases of the fields φ_u^0 and φ_d^0 in the axion state is small, $\sim V_{u,d}/\tilde{V}$. For such an axion, there are no nondiagonal couplings with the fermions, and for the remaining DFSZ parameters of the axion in Eqs. (4) and (5) we have

$$\begin{aligned} A_a^{DFSZ} &= 2N_g\sqrt{z}/(1+z), \quad F_a = \tilde{V},\\ g_f &= \frac{m_f}{F_a}[X_f - (\delta_{fu} + \delta_{fd}z)N_g/(1+z)]\\ &= \frac{m_f m_a}{m_\pi f_\pi A^{DFSZ}}[X_f - (\delta_{fu} + \delta_{fd}z)N_g/(1+z)],\\ C_{a\gamma\gamma} &= \frac{m_a\alpha}{8\pi f_\pi m_\pi}\frac{1+z}{\sqrt{z}}\left(\frac{8}{3} - \frac{2}{3}(4+z)/(1+z)\right)\\ &= \frac{\alpha N_g}{4\pi F_a}\left(\frac{8}{3} - \frac{2}{3}(4+z)/(1+z)\right), \end{aligned} \qquad (9)$$

where $z = m_u/m_d \approx 0.56$; $X_{up} = 2/(x + 1/x)x$; $X_{lept} = X_{down} = 2x/(x + 1/x)$; $x = V_u/V_d$, f in g_f and δ_{fu} (δ_{fd}) takes the values of all physical states of the quarks and leptons $f = u, d, e, s, c, u, b, \tau \ldots$; the Kronecker delta symbol δ_{fu} (δ_{fd}) means that the correction corresponding to it is important only for the u and d quarks; α is the electromagnetic coupling constant.

2) Hadronic axion (Kim–Shifman–Vaĭnshteĭn–Zakharov: KSVZ).[11] In this model, there is an extension in the Higgs sector of the theory, as well as by the introduction of additional fermionic fields (a heavy quark Q of a color triplet). In this case, the Lagrangian of the Yukawa interactions is given by

$$L_{Yuk} = g\bar{Q}_L\sigma Q_R + \text{H.c.} \qquad (10)$$

This Lagrangian is invariant with respect to the global chiral phase transformations

$$\begin{aligned} Q_L &\to \exp(i\omega)Q_L,\\ Q_R &\to \exp(-i\omega)Q_R,\\ \sigma &\to \exp(-2i\omega)\sigma. \end{aligned} \qquad (11)$$

For such an axion, there are no couplings to the leptons at the tree level, and interaction with the light quarks occurs only through mixing with π^0. For the axion parameters in Eqs. (4) and (5), we obtain

$$A_a^{\mathrm{KSVZ}} = \sqrt{z}/(1+z),$$

$$g_{u,d} = \frac{m_u}{F_a(1+z)} = \frac{m_u m_a}{f_\pi m_\pi (1+z) A_a^{\mathrm{KSVZ}}},$$

$$C_{a\gamma\gamma} = \frac{m_a \alpha}{8\pi f_\pi m_\pi} \frac{1+z}{\sqrt{z}} \left(6Q_{em}^2 - \frac{2}{3}\frac{4+z}{1+z} \right)$$

$$= \frac{\alpha}{8\pi F_a} \left(6Q_{em}^2 - \frac{2}{3}\frac{4+z}{1+z} \right). \qquad (12)$$

For the values of $f = e, \mu, \tau, s, c,$...all the Yukawa constants are negligibly small; Q_{em} is the electric charge.

3) Archion model.[5,6] This model arises in the recently developed theory of broken symmetry of the quark and lepton generations. This theory naturally contains a global symmetry $U(1)_H$, the breaking of which leads to the appearance of a Goldstone boson of the invisible-axion type. In Ref. 5 this was called the archion, and it possesses interactions with the fermions that are both diagonal and nondiagonal with respect to the flavors. Identification of this global $U(1)_H$ symmetry with the PQ symmetry is possible in the presence of a triangle anomaly in the interaction of the axial $U(1)_H$ current with the gluons: $aG\tilde{G}$. In the simplest $SU(2)\times U(1)\times SU(3)_H \times U(1)_H$ form with the minimal set of heavy fermions, this anomaly cancels, and the archion remains almost massless, like the arion.[12] However, in contrast to the arion, the archion has no interaction with the photons, since there is a parallel cancellation of the triangle diagram $aF_{\mathrm{em}}\tilde{F}_{\mathrm{em}}$. However, in any realistic grand-unification extension of the model, for example, in $SU(5)\times SU(3)_H$, such cancellations no longer occur because of additional heavy fermions, and the archion becomes similar to the hadronic axion with a strongly suppressed interaction with the leptons.

The symmetry leads to the Yukawa couplings:

$$L_{Yuk} = g_f \bar{f}_{L\alpha} F_{R\alpha} \varphi^0 + g_{nF} \bar{F}_{R\alpha} F_L^{\beta} \xi_{\alpha\beta}^{(n)} + \mu_f \bar{F}_L^{\alpha} f_R^{\alpha} + \mathrm{H.c.}, \qquad (13)$$

where μ_f, with $f = u, d, e$, are mass parameters of the same order, and $F = U,D,E$ are additional heavy fermions; $\xi_{\alpha\beta}^{(n)}$ are horizontal scalars.

The Lagrangian (13) is invariant with respect to global chiral $U(1)_H$ transformations:

$$f_L \to \exp(i\omega) f_L, \quad f_R \to \exp(-i\omega) f_R,$$
$$F_L \to \exp(-i\omega) F_L, \quad F_R \to \exp(i\omega) F_R, \qquad (14)$$
$$\xi^{(n)} \to \exp(2i\omega)\xi^{(n)}, \quad n=0,\ 1,\ 2, \quad \varphi \to \varphi.$$

The parameters in (4) and (5) for the archion model have the form:

$$A_a = A_c \sqrt{z}/(1+z), \quad g_{33} = \frac{m_3^{(f)}}{V_1}, \quad g_{22} = \frac{m_2^{(f)}}{V_1} (S_{23}^{(f)})^2,$$

$$g_{11} = \frac{m_2^{(f)}}{V_1} (S_{13}^{(f)})^2,$$

$$g_{23} = \frac{m_3^{(f)}}{2V_1} S_{23}^{(f)}, \quad g_{13} = \frac{m_3^{(f)}}{2V_1} S_{13}^{(f)}, \quad g_{12} = \frac{m_2^{(f)}}{2V_1} S_{13}^{(f)} S_{23}^{(f)},$$

$$(15)$$

$$C_{a\gamma\gamma} = \frac{\alpha m_a}{4\pi f_\pi m_\pi} \frac{z}{z+1} \left(\frac{A_{em}}{A_c} - 2\frac{(4+z)}{3(1+z)} \right)^{-1}$$
$$= \frac{\alpha A_c z^{3/2}}{4\pi F_a (1+z)^2} \left(\frac{A_{em}}{A_c} - 2\frac{(4+z)}{3(1+z)} \right)^{-1},$$

where $s_{ij} = \sin\varphi_{ij}$ $(i,j = 1,2,3)$ are the angle parameters of the unitary rotation matrices V^f that diagonalized the fermion mass matrices, $V_{fL}^{+} m^f V_{fR} = \hat{m}^f_{\mathrm{diag}}$, and the phases are determined by the phase content of the V matrices:

$$\varphi_{12} = \arg(V_{23}V_{13}/V_{13}), \; \varphi_{13} = -\arg V_{13}, \; \varphi_{23} = -\arg(V_{12}V_{23}). \qquad (15a)$$

$V_1 = F_a$ is the scale of the $U(1)_H$ breaking, and A_c and A_{em} are the color and electromagnetic anomalies, respectively.

In all three models of the invisible axion, the Lagrangian of the interaction with the nucleons has the form

$$L_{NN'}^{a} = \bar{N}' i\gamma_5 (g^{(0)} + g^{(3)}\tau_3) N a, \qquad (16)$$

where for the DFSZ axion

$$g^{(0)} = \frac{M}{2F_a} G_A^{(0)}; \quad g^{(3)} = \frac{M}{2F_a}\frac{1-z}{1+z} G_A,$$

and for the hadronic axion and archion

$$g^{(0)} = \frac{M}{2N_g F_a}(X_u + X_d - N_g) G_A^{(0)};$$

$$g^{(3)} = \frac{M}{2N_g F_a}\left(X_u - X_d - \frac{1-z}{1+z} \right) G_A,$$

where M is the nucleon mass, and $G_A = G_A^{(0)} = 1, 2, 3$ is the isotriplet axial form factor of the nucleon, which is determined in experiments from measurement of the parameters of the weak charged currents of the nucleons.

There currently exist certain restrictions on the parameters of the models of an invisible axion. The best of them—lower bounds on the symmetry-breaking scale F_a (upper bounds on the axion mass m_a)—come from astrophysical arguments based on analysis of the influence of axion-radiation energy loss on the rate of evolution of stars of different types. The bounds[13] are as follows:

1) Main-sequence stars (like the sun):[14]

$$m_a^{\mathrm{DFSZ}} \leqslant 1 \text{ eV}, \quad F_a^{\mathrm{DFSZ}} \geqslant 3{,}7\cdot 10^7 \text{ GeV}.$$

$$m_a^{\mathrm{KSVZ}} \leqslant 20 \text{ eV}/\zeta, \quad F_a^{\mathrm{KSVZ}} \geqslant 2{,}9\cdot 10^5\, \zeta \text{ GeV}$$

2) Red giants:[15]

$$m_a^{\mathrm{DFSZ}} \leqslant 10^{-2} \text{ eV}, \quad F_a^{\mathrm{DFSZ}} \geqslant 3{,}7\cdot 10^9 \text{ GeV}$$

$$m_a^{\mathrm{KSVZ}} \leqslant 2 \text{ eV}/\zeta, \quad F_a^{\mathrm{KSVZ}} \geqslant 2{,}9\cdot 10^6\, \zeta \text{ GeV}$$

3) White dwarfs:[16]

$$m_a^{\mathrm{DFSZ}} \leqslant 3\cdot 10^{-2} \text{ eV}, \quad F_a^{\mathrm{DFSZ}} \geqslant 3{,}7\cdot 10^9 \text{ GeV}$$

Here, $\zeta = (A_{\mathrm{em}}/A_c - 1.95)/0.72$.

At the same time, using cosmological arguments that bound the energy density of the oscillations of the primordial axion field, we can obtain upper bounds on the scale F_a (lower bounds on the mass m_a)[17]

$$m_a^{\mathrm{DFSZ}} \gtrsim 7\cdot 10^{-6}\ \mathrm{eV}, \quad F_a^{\mathrm{DFSZ}} \lesssim 3\cdot 10^{12}\ \mathrm{GeV}.$$

$$m_a^{\mathrm{KSVZ}} \gtrsim 8\cdot 10^{-6}\ \mathrm{eV}, \quad F_a^{\mathrm{KSVZ}} \lesssim 8\cdot 10^{11}\ \mathrm{GeV}.$$

Use of experimental hints of the presence of a neutrino signal preceding the explosion of the supernova SN 1987A,[18] with identification of this signal with neutrino emission of the collapsing core of the star, makes it possible to obtain additional bounds on the parameters of the invisible axion. Analysis of the admissible influence of the radiation on the energy balance and time characteristics of the neutrino radiation in the formation of a neutron star rules out the range of values[18]

$$10^6 < F_a < 10^{10}\ \mathrm{GeV} \tag{17}$$

Thus, the use of the data of neutrino astronomy leaves very narrow intervals of allowed values for the parameters in the models of the invisible axion with $F_a \sim 10^6$ GeV and $F_a \sim 10^{10}$–10^{14} GeV.

We mention in this connection theoretical arguments for a "low" ($F_a \sim 10^6$ GeV) scale of the axion interaction. These arguments are associated with the problem of a hierarchy of breaking scales of the Peccei–Quinn and electroweak symmetries present in the models of the invisible axion. In such models, there is necessarily an interaction

$$\Delta L = \lambda(\varphi^+\varphi)(\sigma^+\sigma), \tag{18}$$

for which the "natural" value of the constant λ is determined by radiative effects. On breaking of the Peccei–Quinn symmetry, this interaction leads to generation of a correction and to a mass term of the Higgs field φ, $\Delta m^2 \sim \lambda F_a^2$. The hierarchy in the breaking of the Peccei–Quinn and electroweak symmetries requires that this correction be not greater than V^2, i.e., it requires

$$\lambda < V^2/F_a^2. \tag{19}$$

Such a condition appears "natural" in the case of a "low" ($F_a \sim 10^6$ GeV) scale but certainly requires additional special mechanisms of suppression in the case $F_a > 10^{10}$ GeV. Indeed, the interaction (18) can be induced by the diagrams of Fig. 1, giving in order of magnitude:

$$\text{a)}\ \lambda \sim g^2\alpha_c^2G^2/(2(4\pi^2)^2), \qquad \text{b)}\ \lambda \sim g_f^2g_{nf}^2/(8\pi^2). \tag{20}$$

Taking in order of magnitude $g \sim G$, $g_f \sim g_{nf} \gtrsim 10^{-2}$, we obtain: a) $\lambda \gtrsim 10^{-14}$, b) $\lambda \gtrsim 10^{-10}$. These results must be compared with $\lambda = 10^{-8}$ and $\lambda < 10^{-16}$ for the cases of "low" and "high" scales F_a, respectively.

3. THERMAL COSMOLOGICAL BACKGROUND OF AXIONS

The above analysis of the bounds on the axion model in the case of an axion of hadronic type allows a parameter value $F_a \sim 10^6$ GeV. We shall show here that this leads to a qualitatively new effect of establishment of equilibrium of the axions with matter and radiation in the early universe. Such equilibrium is realized after a QCD phase transition ($T_* \sim 160$ MeV) due to processes of the type $\pi N \leftrightarrow aN$.

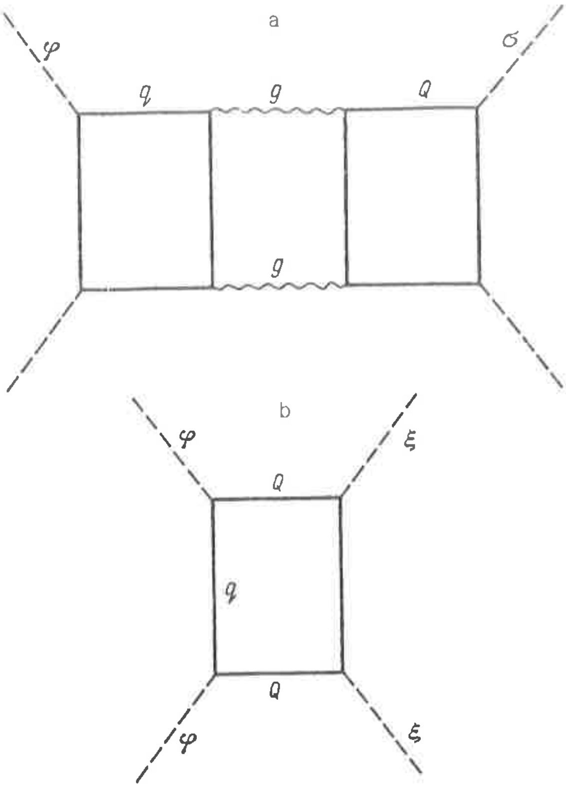

FIG. 1. Diagrams inducing the interaction (18) in the models of the hadronic axion (a) and the archion (b).

The kinetics of the establishment of equilibrium between the axions and the plasma in this process is described by the equation[19]

$$\frac{dn_a}{dt} = \sigma_0 c n_\gamma(\Delta n_N + 2n_N) + \psi(t) - \alpha n_a, \tag{21}$$

where n_a and n_N are the axion and nucleon concentrations, respectively; Δn_N is the excess nucleon concentration; σ_0 is the cross section of the process $\pi N \leftrightarrow aN$, $\psi(t)$ is a function that describes the axion production rate; αn_a is a term that takes into account the cosmological expansion; and c is the speed of light. Introducing the ratios $r_a = n_a/n_\gamma$, $\Delta r_N = \Delta n_N/n_\gamma$, and $r_N = n_N/n_\gamma$, where n_γ is the concentration of the light particles (e^+, e^-, γ), which satisfies the expansion law

$$\frac{dn_\gamma}{dt} = -\alpha n_\gamma, \tag{22}$$

we can rewrite Eq. (21) in the form

$$\frac{dr_a}{dt} = -\sigma_0 c n_\gamma r_a(\Delta r_N + 2r_N) + \psi/n_\gamma. \tag{23}$$

Under equilibrium conditions,

$$r_N' \approx \theta^{-3/2}\exp(-1/\theta), \tag{24}$$

where $\theta_N = kT/M_N$, the condition for decoupling of the axions has the form

$$\sigma_0 c n_\gamma t(\Delta r_N + 2r_N') = 1. \tag{25}$$

where the time is $t = 10^{-6}\theta_N^{-2}$ sec. The cross section of the process $\pi^0 N \leftrightarrow aN$ can be related to the cross section for $\pi^0 N \leftrightarrow \pi^0 N$ by

$$\sigma_0 = A_c^2 (f_\pi/F_a)^2 \sigma(\pi^0 N \leftrightarrow \pi^0 N). \qquad (26)$$

Since for models of the hadronic-axion type we have $A_c \approx 1$, we obtain

$$\sigma_0 = V_6^{-2} 10^{-14} \sigma(\pi^0 N \leftrightarrow \pi^0 N), \qquad (27)$$

where $V_6 = V_H/10^6$ GeV. The cross section of the process $\pi^0 N \leftrightarrow \pi^0 N$ depends weakly on the temperature at energies $\sim T_*$:

$$\sigma \approx 1{,}7 \cdot 10^{-27} (2+7\theta)^2 \ \text{cm}^2 \qquad (28)$$

Thus, using Eqs. (27) and (28), and bearing in mind that $c = 3 \times 10^{10}$ cm/s and $n_\gamma = 10^{41}\theta^3$, we rewrite the decoupling condition (25) in the form:

$$2 \cdot 10^5 (2+7\theta)^2 (\Delta r_N + 2\theta^{-3/2} e^{-1/\theta})\theta / V_6^2 = 1. \qquad (29)$$

Simplifying this expression, we can obtain an approximate dependence of the decoupling temperature on V_6:

$$15 - 1/\theta = \ln(V_6^2 - 8 \cdot 10^5 \Delta r_N \theta). \qquad (30)$$

The value of Δr_N is determined by the contemporary ratio of the baryon and photon concentrations and can be expressed in terms of the relative contribution Ω_B of the baryons to the total cosmological density[19]

$$\Delta r_N \approx \frac{\Delta n_N}{n_\gamma} \approx 7.5 \cdot 10^{-8} \Omega_B \left(\frac{2{,}7k}{T}\right)^3 \left(\frac{H}{50\ \text{km/sec} \cdot \text{Mpc}}\right) \qquad (31)$$

As can be seen from the expression (30), the term associated with the baryon excess can be ignored, and we then obtain a simple explicit dependence of the decoupling temperature on V_6:

$$\theta = \frac{1}{15 - 2\ln V_6}; \qquad T = \frac{M_N}{15 - 2\ln V_6}. \qquad (32)$$

Analysis of the expression (32) (Fig. 2) for various values of V_6 shows that in the interval $10^6 \leqslant V_H \leqslant 10^8$ GeV equilibrium of the axion gas with the matter and radiation is established. Moreover, the equilibrium is ensured by the presence of equilibrium $N\bar{N}$ pairs.

The phenomenon noted above gives grounds for the assertion that in the framework of models of the hadronic-axion type with a low energy scale $10^6 < V_H < 10^8$ GeV the presence of a thermal axion background in the universe is predicted.

The concentration of such axions relative to the neutrino concentration in the contemporary universe can be determined. At the time of decoupling of the axions, the plasma contained $e^\pm$, ν, $\bar{\nu}$, γ, $\pi^\pm$, π^0, μ^+, and μ^-. However, the π and μ equilibrium concentration was low compared with the $e^\pm$ concentration:

$$\frac{n_\pi}{n_e} \approx \left(\frac{m_\pi}{T}\right)^{3/2} \exp(-m_\pi/T); \qquad \frac{n_\mu}{n_e} \approx \left(\frac{m_\mu}{T}\right)^{3/2} \exp(-m_\mu/T). \qquad (33)$$

Before the decoupling of the axions, they had an equilibrium concentration relative to the neutrino concentration:

$$n_a/n_\nu^* = 2/3 = 0{,}67. \qquad (34)$$

After decoupling, there was heating through π and μ decays.

Using the conservation of the specific entropy of a relativistic gas with temperature T and Q internal degrees of freedom,

$$S = (\pi^2/30)(Q/\rho)T^3, \qquad (35)$$

we can find the change in the concentration of ν.

If before decoupling the number of degrees of freedom was

$$Q_0 = 2 + 3\frac{7}{4} + \frac{7}{2} + 3\frac{n_\pi}{n_e} + \frac{7}{8}\frac{n_\mu}{n_e} = 10.75 + 3\frac{n_\pi}{n_e} + \frac{7}{8}\frac{n_\mu}{n_e}, \qquad (36)$$

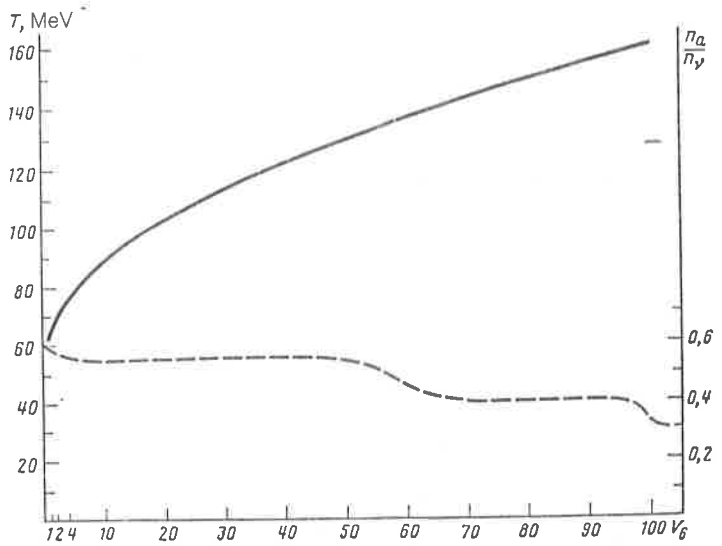

FIG. 2. Dependence of the axion decoupling temperature on the scale (continuous curve). Dependence of the relative concentration n_a/n_ν on the scale (broken curve).

then after the decoupling of ν we have

$$Q=2+3\frac{7}{4}+\frac{7}{2}=10,75. \quad (37)$$

Therefore, the ν concentration increased by a factor of k, where

$$k=Q_0/Q=(10,75+3n_\pi/n_e+3,5n_\mu/n_e)/10,75. \quad (38)$$

The dependence of the relative concentration of the thermal axions on the scale V_H [with allowance for (32)] is shown in Fig. 2. In particular, for $V_H = 10^6$ GeV,

$$n_a/n_\nu=0,57. \quad (39)$$

We note that in accordance with Ref. 20 the axionic Primakov effect in the $\gamma q \leftrightarrow aq$ reaction, and also the interaction of axions with diagonal and nondiagonal quark transitions, can ensure establishment of equilibrium for the axions at a temperature $T > T_{\rm QCD}$, the temperature of the QCD phase transition. However, making the corresponding estimates, we can establish that the concentration of the axions relative to the concentration of the neutrinos in this case will be

$$n_a/n_\gamma=0,29, \quad (40)$$

which is less, by a factor of 2, than the analogous quantity (39) for the process $\pi N \leftrightarrow aN$.

4. THE PROBLEM OF PRIMORDIAL SUPERHEAVY QUARKS IN THE MODEL OF THE HADRONIC AXION

As we have already noted, in the model of the hadronic axion there is not only an extension of the Higgs sector of the theory of the standard axion but also the introduction of the field of a superheavy quark Q, which possesses an arbitrary charge. The frozen concentration of such quarks can make an appreciable contribution to the contemporary density of the universe.

This concentration can be estimated by using the quenching equation[19]

$$\frac{dn_Q}{dt}=-\langle\sigma v\rangle n_Q^2+\psi(t)-3Hn_Q, \quad (41)$$

where $\langle\sigma v\rangle$ is the rate of annihilations of the type $Q\bar{Q}\to gg$, qq, n_Q is the concentration of the heavy quarks, and H is the Hubble constant.

One can also obtain the quenching condition for the superheavy quarks

$$4\nu_Q n_r\langle\sigma v\rangle(T/m_Q)t=1, \quad (42)$$

where n_r is the concentration of the relativistic particles at the time of Q quenching, and T is the quenching temperature corresponding to the time t. Bearing in mind that $n_r = (\pi^2/30)T^3$, $T \approx m_Q/3$, $t = M_{\rm P}/T^2$, $\langle\sigma v\rangle \approx \pi\alpha_s^2/m_Q^2$,[21] we obtain for the concentration of the remaining quarks Q the estimate:

$$\nu_Q=n_Q/n_r=0,7m_Q/\alpha_s^2 M_{\rm P}. \quad (43)$$

From the condition $\nu_Q m_Q < m_B \nu_B$, where $\nu_B \approx 7.5\times 10^{-8}\,\Omega_B$ is determined by the relation (31), we can obtain an upper bound for the mass of such exotic quarks: $m_Q \lesssim 5\times 10^4$ GeV.

After the QCD phase transition, hadronization occurs, and the heavy quarks enter superheavy hadrons. The Q and $\bar{Q}$ in such hadrons can annihilate by virtue of recombination reactions of the type $(q\bar{Q}) + (qqQ) \to 3q + (Q\bar{Q})$ accompanied by decay of the $(Q\bar{Q})$ quarkonium into ordinary hadrons. According to Ref. 21, the cross section of such reactions is parametrized by $\sigma_r \sim f\times 100$ mb, where $0<f<1$. For $f\sim 1$, the concentration of superheavy quarks in the hadron combustion stage can be reduced to $n_H/n_B \gtrsim 10^{-9}$, where n_H is the concentration of the superheavy hadrons, and n_B is the concentration of the baryons. This also raises the upper bound on the Q mass: $m_Q \leqslant 10^6$ TeV. In the course of further cosmological evolution, the Q hadrons can form anomalous nuclei of exotic isotopes. According to Ref. 21, the expected concentration of such nuclei contradicts the bounds on the admissible concentration of exotic isotopes even when $f\sim 1$. Thus, in the case of stable heavy quarks the model of the hadronic axion contradicts observational data, and the theory must contain a mechanism that ensures instability of such quarks. It should be noted that the inclusion of the model of the hadronic axion in grand-unification schemes necessarily leads in its realization to the appearance of a heavy lepton that interacts with the axion. Then mixing of the heavy Q quark with the light quarks, which would ensure instability of the Q, would lead to the appearance, at the tree level, of coupling of the axion to the leptons, so that the axion would cease to be hadronic. In the light of these difficulties of the model of the hadronic axion, particular importance must be accorded to the archion model, which ensures in a natural manner both instability of the heavy quarks and suppression of coupling of the archion to the leptons. We therefore turn to an analysis of the cosmological axion background predicted in the framework of this model.

5. COSMOLOGICAL BACKGROUND OF ARCHIONS

Models of quantum flavordynamics are associated with a unified theoretical scheme of various candidates for the role of the dark-mass particles discussed in the theory of formation of structure in the universe. In the case of a low (10^6–10^8 GeV) scale of generation symmetry breaking, which also determines the scale of the archion interaction, the model predicts not only a thermal archion background but also a further source of cosmological archions—decays of heavier neutrinos into lighter neutrinos and archions. Moreover, the model fixes the hierarchy of masses and the neutrino lifetimes.[5,6,22]

The relative contribution of the massive neutrinos and archions to the cosmological density is determined by the parameters of the model,[5] namely, by the scale V_1 of breaking of the horizontal symmetry and by the ratio of the Ykawa constants (g^2/G).

For fixed contemporary total density $\rho_{\rm tot}$ and baryon density ρ_B, the relation

$$\rho_a(V_1)+\rho_{\nu_\tau}(V_1)+\rho_{\nu_\mu}(V_1)+\rho_{\nu_e}(V_1)+\rho_B=\rho_{tot} \quad (44)$$

becomes an equation for V_1 and (g^2/G), whose solutions determine a discrete set of cosmological models with different types of dark matter forming the structure of the universe. In the framework of the archion model, there are, in the general case, six possible realizations of the dark matter.[5]

1. Cold dark matter. In this case, there is dominance of the density of the primordial axion field[23] and of axions from the decay of the structure of cosmic strings:[24]

$$\rho_a = A_r^2 (V_6/(2\cdot 10^4))\rho_{cr}. \tag{45}$$

2. Hot dark matter. There is dominance of the density ρ_{ν_τ} of the massive stable neutrinos ν_τ with mass $m_{\nu_\tau} \simeq 24$ eV and lifetime $\tau(\nu_\tau \to \nu_\mu a)$ greater than the age t_V of the universe:

$$\rho_{\nu_\tau} = 6{,}6\cdot 10^7/V_6 (g^2/G)\rho_{cr}. \tag{46}$$

3. Dominance of the densities of archions, ρ_a, and of muon neutrinos, ρ_{ν_μ}, which are relativistic products of decay of massive unstable ν_τ with mass $m_{\nu_\tau} = 50$–100 eV and lifetime $\tau(\nu_t \to \nu_\mu a) = 4\times 10^{16}$–$10^{16}$ sec; if $m_{\nu_\mu} \leqslant 5$ eV, then

$$\rho_{\nu_\mu + a} = \frac{V_6^{1/2}}{x 10^6 (g^2/G)^{1/2}}\rho_{cr}, \tag{47}$$

where $x \sim 1$ is the neutrino mixing parameter.

4. Dominance of the density ρ_{ν_μ} of nonrelativistic ν_μ from decay of ν_τ and primordial ν_μ if $m_{\nu_\mu} \simeq 10$ eV, $\tau_{\nu_\mu} > t_U$:

$$\rho_{\nu_\mu} = \frac{6{,}6\cdot 10^6}{V_6}(g^2/G)\rho_{cr}. \tag{48}$$

5. Dominance of the densities ρ_a of nonrelativistic archions and of electron neutrinos, ρ_{ν_e}, from the decay of ν_μ with mass $m_{\nu_\mu} = 50$–100 eV and lifetime $\tau(\nu_\mu \to \nu_e a) = 10^{16}$–$2\times 10^{15}$ sec under the condition of rapid decay of ν_τ with mass $m_{\nu_\tau} \sim 1$–10 keV, $\tau(\nu_\tau \to \nu_\mu a) \leqslant 10^8$–$10^{10}$ sec:

$$\rho_{\nu_e + a} = \frac{V_6^{3/2}}{x(g^2/G)\cdot 10^3}\rho_{cr}. \tag{49}$$

6. Dominance of the density ρ_a of nonrelativistic (or semirelativistic) archions, products of early decays of ν_μ (if $m_{\nu_\mu} > m_a$), or (if $m_{\nu_e} > m_a$) dominance of the density ρ_{ν_e} of nonrelativistic ν_e from decays of ν_τ and subsequent decays of ν_μ:

$$\rho_{\nu_e} = \frac{3{,}3\cdot 10^4}{V_6}(g^2/G)\rho_{cr}. \tag{50}$$

At the same time, for ρ_a the main contribution to the inhomogeneous dark matter is made, besides the primordial axions, by nonrelativistic axions from decays of ν_t:

$$\rho_a = \rho_{cr}/10V_6. \tag{51}$$

Thus, in the considered model there are physical justifications for the scenarios of cold (in the form of a primordial axion condensate), hot (in the form of stable neutrinos ν_τ), and unstable (in the form of unstable neutrinos ν_τ in cases 3 and 4, and ν_μ in cases 5 and 6) matter, these being realized as solutions of Eq. (44) of the form

$$\begin{aligned}
&\frac{V_6}{1{,}7\cdot 10^4} + \frac{7{,}4\cdot 10^7}{V_6}\left(\frac{g^2}{G}\right)\exp\left(-\frac{t_U}{\tau_{\nu_\tau}}\right)\\
&+\frac{V_6^{3/2}}{x(g^2/G)^{1/2}10^6}\exp\left(-\frac{t_U}{\tau_{\nu_\mu}}\right)\\
&+\frac{V_6^{3/2}}{x(g^2/G)\cdot 10^3}\left(1-\exp\left(-\frac{t_U}{\tau_{\nu_\mu}}\right)\right)\\
&+\frac{7{,}4\cdot 10^6}{V_6}\left(\frac{g^2}{G}\right)\exp\left(-\frac{t_U}{\tau_{\nu_\mu}}\right)\\
&+\frac{1{,}1}{10V_6}\left(1-\exp\left(-\frac{t_U}{\tau_{\nu_\tau}}\right)\right)\\
&+\frac{3{,}7\cdot 10^4}{V_6}\left(\frac{g^2}{G}\right)\left(1-\exp\left(-\frac{t_U}{\tau_{\nu_\mu}}\right)\right)=1.
\end{aligned} \tag{52}$$

where the lifetime of neutrinos relative to the decays $\nu_H \to \nu_L a$ is given by[5,6,22]

$$\tau(\nu_H \to \nu_L a) = 16\pi/(g_{HL}^2 m_H), \tag{53}$$

and g_{HL} are obtained from (15).

Equation (52) describes the phase curve in the space of model parameters $[V_1, (g^2/G)]$, the points of which correspond to all possible scenarios and their combinations. The curves are given in Figs. 3 and 4, respectively, for the massive and massless archions.

It is readily seen that it is only in the range of constants $(g^2/G) \approx 1.5\times 10^{-6}$–$1.5\times 10^{-4}$ (depending on the scale V_1) that there can be realization of the complete set of cosmological solutions of Eq. (44), since outside this range the dominant density exceeds (or does not reach) the critical

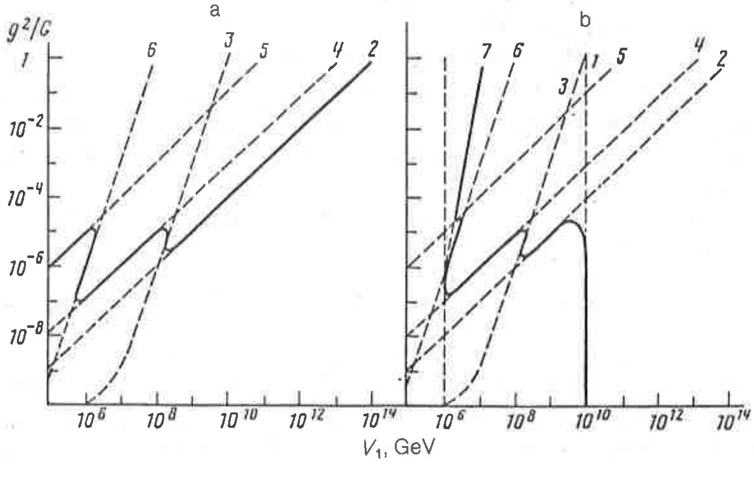

FIG. 3. Phase curve in the parameter space $[V_1, (g^2/G)]$ of the archion model: a) the case of the arion ($A_c = 0$); b) the case of the axion ($A_c \neq 0$). The number of each curve corresponds to the number of the cosmological scenario (7 corresponds to dominance of nonrelativistic axions).

density ρ_{cr}. On the other hand, astrophysical bounds leave only a narrow window around $V_l \sim 10^6$ GeV in which there may be a realization of scenarios 5 or 6 of unstable dark matter, and the broader interval $10^{10} < V_l < 10^{14}$ GeV in which scenarios 1 or 2 can be realized. Scenarios 3 and 4 are excluded by the bounds on V_l and by SN 1987A. Thus, the above range of variation of the constants (g^2/G) in conjunction with the bounds of the scale V_l indicate that only the unstable forms of dark matter in the form of scenarios 5 and 6 can actually exist.

We note in this connection that the hierarchical decay scenarios 5 and 6 combine the advantages of the models of cold, hot, and unstable dark matter, making these scenarios a very attractive, physically justified theoretical basis for detailed development of models of the formation of structure in the universe and comparison of the predictions of such models with the data of astronomical observations. Indeed, in these models the period of dominance of ν_τ of mass $\sim$1–10 keV introduces into the particle spectrum a short-wavelength part characteristic of models of cold dark matter, and the ν_μ from the decays of the ν_τ enhance, in the perturbation spectrum, the long-wavelength component inherent in the hot dark matter and needed for the formation of a clear cellular–mesh structure of "voids" and superclusters. Finally, ν_μ decays retard the rate of evolution of such structure and sustain it to the present time. At the same time, the primordial thermal archions, which currently form the coldest component of the dark matter, play an important part in the development of the part of the density perturbation spectrum with the shortest wavelengths, determining, in particular, the formation of massive halos outside the visible region. In this connection, it should be noted that in accordance with Ref. 25 bounds from the density of phase space on the mass of the halo particles[26] can be weakened or may even be absent for bosons.

In accordance with either of the scenarios 5 or 6, the universe must contain, besides a thermal archion background, a nonthermal archion background from ν_τ and ν_μ decays. The spectrum of the nonthermal archions will have the form

$$\frac{dn_a}{dk} = \frac{dn_a}{dk}\bigg|_{\nu_\mu} + \frac{dn_a}{dk}\bigg|_{\nu_\tau}, \tag{54}$$

where $dn_a/dk\,|_{\nu_\tau}$ is the spectrum from the decays $\nu_\tau \to \nu_\mu a$, and $dn_a/dk\,|_{\nu_\mu}$ is the spectrum from the decays $\nu_\mu \to \nu_e a$. In fact, the main contribution to the expression (54) is made by decays that occur at the epoch of nonrelativistic dominance.

We determine the form of the spectra of the decay axions for two values of the Yukawa constants for the case when $V_l = 10^6$ GeV. For this, we use the general expression for the spectrum of decay particles in an expanding universe at time t:[22]

$$\frac{dn_a}{dk} = 2n_t N \frac{t}{\tau_\nu} \frac{k^{N-1}}{k_0{}^N} \exp\left(-\left(\frac{k}{k_0}\right)^N \frac{t}{\tau_\nu}\right), \tag{54A}$$

where $n_t = n_0(t_U/t)^{3/2}$, $k_0 \simeq m_\nu/2$, $N = \frac{3}{2}$ for the stage of nonrelativistic dominance, $N = 2$ for the stage of relativistic dominance, and τ_ν is the neutrino lifetime. Taking into account the evolution of the spectrum from the time t to the contemporary epoch, we obtain the following dependences: for $(g^2/G) \approx 1.5 \times 10^{-6}$

$$\begin{aligned} \frac{dn_a}{dk}\bigg|_{\nu_\tau} &\approx 1{,}1\cdot 10^7\, k^{1/2}, \quad k_{max} \approx 6\cdot 10^{-5}\ \text{eV}, \\ \frac{dn_a}{dk}\bigg|_{\nu_\mu} &\approx 6{,}3\cdot 10\, k^{1/2}, \quad k_{max} \approx 3\ \text{eV}, \end{aligned} \tag{55}$$

and for $(g^2/G) \approx 1.5 \times 10^{-4}$,

$$\begin{aligned} \frac{dn_a}{dk}\bigg|_{\nu_\tau} &\approx 3{,}9\cdot 10^9 k^{1/2}, \quad k_{max} \approx 1{,}1\cdot 10^{-5}\ \text{eV}, \\ \frac{dn_a}{dk}\bigg|_{\nu_\mu} &\approx 3{,}1\cdot 10^4\, k^{1/2}, \quad k_{max} \approx 5\cdot 10^{-2}\ \text{eV}. \end{aligned} \tag{56}$$

The profiles of (55) and (56) are shown in Fig. 4.

It is obvious that after the occurrence of inhomogeneities and the formation of the halos of the galaxies the "homogeneous" spectra (55) and (56) change because the cold particles (axions and primordial ν_e), which for the chosen parameters have momentum $k \lesssim 10^{-2}$ eV, will be trapped in the halos. Taking into account this fact, we can determine the fraction of particles of the dark matter that at the present time form galactic halos. Assuming that $m_{\nu_e} \approx m_a$, and integrating the spectra in the appropriate ranges (the decay

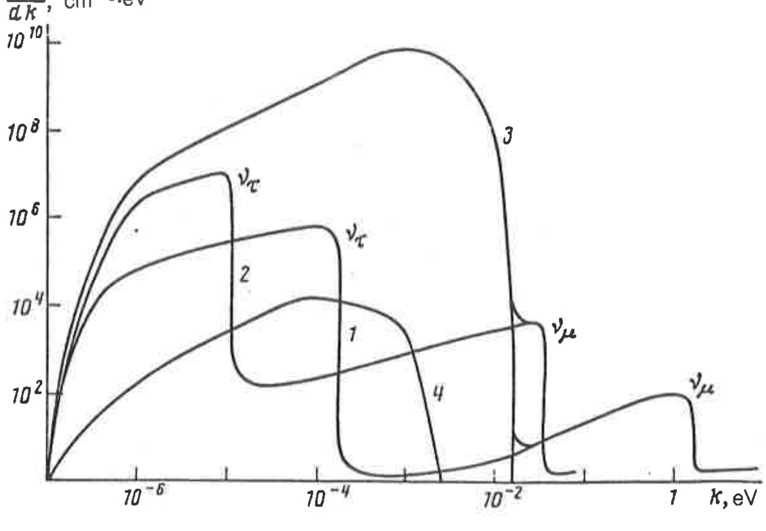

FIG. 4. Curves 1 and 2 are the spectra of axions from decays of ν_μ and ν_τ for the cases $g^2/G = 1.5 \times 10^{-6}$ and $g^2/G = 1.5 \times 10^{-4}$, respectively; curve 3 is the spectrum of axions in the halos of galaxies $(dn_a/dk\,|_G = 1.4 \times 10^{16}[k^2/(e^{k/T_a} - 1)])$; curve 4 is the spectrum of thermal axions $(dn_a/dk\,|_T = 2.5 \times 10^{12}[k^2/(e^{k/T} - 1)])$.

spectra of the axions and their thermal spectrum), we find that at the present time about 40% of the dark matter in the universe has gone into the inhomogeneities.

6. CONCLUSIONS

Our analysis shows that the combination of laboratory and astronomical searches for the effects of an invisible axion has made it possible to obtain very definite conclusions about its existence and admissible properties. Thus, the study of models of the invisible axion provides a rather convincing example of the effectiveness of the cosmological–microphysical approach to investigation of the physics of superhigh energies on the basis of a comprehensive analysis of the totality of its indirect manifestations. Analysis of all the effects predicted by the different models of the invisible axion permits a very well-founded selection of these models, distinguishing, as the most promising, the archion model with a low ($\sim 10^6$ GeV) scale of the symmetry breaking of the generations of quarks and leptons. In addition to the previously discussed predictions of such a model, we now turn to verification in experiments looking for archionic decays of τ, μ, and K and B mesons[27] and astronomical searches for a nonthermal electromagnetic background from $a \to 2\gamma$ decays.

In this paper, we analyze a further inescapable consequence of this model—the existence of a primordial thermal background of cosmic archions. This conclusion is also important for the development of the theory of the formation of large-scale structure of the universe on the basis of the cosmological scenarios, predicted in the framework of quantum flavordynamics with a low scale of the generation symmetry breaking of the hierarchy of decays of massive unstable neutrinos, for the search for a cosmological background of archions by the experimental methods of axion haloscopy, and for the astronomical search for electromagnetic radiation from decays of archions in the contemporary universe. The development of methods for detecting fluxes of cosmological axions have here made it possible to hope for manifestation of axion astronomy.

Attention should also be drawn to a further important epistemological aspect of the investigations made here—the possibility of detailed study of the many-parameter "hidden" sector of the theory of elementary particles. The example of quantum flavordynamics with a low energy scale encourages the hope that any well-developed many-parameter model of the physics of ultrahigh energies will lead to an abundance of indirect consequences capable of experimental or observational verification, these exceeding in number the number of independent parameters of the theory, so that the verification of such consequences can be reduced to an overdetermined system of equations for these parameters. Thus, in the framework of cosmomicrophysics it is possible to realize a general approach to experimental verification of a theory on the basis of an overdetermined system of equations for unknown parameters. Analysis of the combined indirect effects of the theory permits detailed study of the theory in a field not accessible to direct experimental verification.

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