$$I_0 = \frac{(2\pi)^6}{66}\alpha^2 G^2 (Ze)^2 m^6 n_0 \left(\frac{kT}{m}\right)^{10}. \quad (15)$$

The minimum proton density in the Oppenheimer–Volkoff model of a neutron star is $n_0 \sim n_e \sim 10^{-33}$ cm^{-3} (for the transitional matter density).[18] Using this value and taking $kT \sim m$ (which is the case at an early stage of the evolution) we obtain

$$I_0 \sim 10^{20-23} \text{ (erg/cm}^3\text{ sec)} \quad (16)$$

In view of the strong temperature dependence of I_0 [Eq. (15)] this result can be regarded only as a very rough estimate. Nevertheless, it indicates that under these conditions the effect under discussion may make the dominant contribution to the total neutrino luminosity of a magnetic star. In fact, the volume of a typical neutron star is $\sim 10^{18-19}$ cm^3, so that the neutrino luminosity due to the mechanism under consideration may be some 10^{38-39} erg/sec, and this is comparable to the contribution from the $n+n \to n+p+e^- + \bar{\nu}$ reaction, which is generally assumed to give the main contribution. Thus, the rather exotic nonlinear process $\gamma(Ze) \to \nu\bar{\nu}(Ze)$ may affect the evolution of macroscopic objects—neutron stars [we drew the same conclusions in Ref. 4 about another nonlinear process, $\gamma \to \nu\bar{\nu}$, which, however, can take place only under certain conditions (see the Introduction)]. The possibility, noted above, that the neutrino radiation may be channeled along the magnetic lines of force should be emphasized once more; in the final analysis such channeling is a consequence of parity nonconservation in the weak interactions. That this last phenomenon can be detected experimentally seems, unfortunately, problematical, although from a theoretical point of view it would not be without interest.

[1)]We note that this differs substantially from the conclusions drawn from "four-dimensional" quantum electrodynamics, in which the inclusion of a single pseudovector vertex may lead to a nonvanishing result.[8] In this sense an "expanded" analog of Furry's theorem holds in the "two-dimensional variant" of quantum electrodynamics.

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Astrophysical bounds on the mass of heavy stable neutral leptons

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Analytical and numerical calculations show that heavy neutral stable leptons are carried along by the collapsing matter during the formation of galaxies and possibly stars as well. The condensation in galaxies and stars results in appreciable annihilation of leptons and antileptons. Modern observations of cosmic-ray and γ-ray fluxes establish a limit $m_\nu \gtrsim 100$ GeV for the mass of neutral leptons, since annihilation of neutral leptons produces γ rays and cosmic rays. The obtained bound, in conjunction with ones established earlier, precludes the existence of stable neutral leptons (neutrinos) with $m_\nu > 30$ eV.

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1. INTRODUCTION

The possible existence of stable neutral leptons with nonzero rest mass has been widely discussed in recent years in connection with the development of the theory of elementary particles and the discovery of the τ lepton. In a number of theories, the neutrino can have a small nonzero mass,[1] and new, absolutely stable heavy neutral leptons also appear.[2]

According to the theory of the Hot Big Bang, the heavy neutral leptons and neutrinos must have been in equilibrium with matter during the early stages in the evolution of the Universe. During the cosmic expansion, the temperature decreased, and at a certain time the leptons came out of thermodynamic equilibrium, this determining the residual ("frozen") concentration of leptons in the Universe.[3-9] It was shown in Refs. 3-9 that this residual concentration of leptons contributes to the total matter density in the Universe, which makes it possible to obtain bounds on their mass. Leptons cannot have mass in the interval 30 eV $< m_\nu < 2$ GeV, since the existence of such leptons would lead to a high mean cosmological density and, therefore, to a contradiction with the observed age of the Universe. The upper limit arises because with increasing mass (for $m_\nu > 1$ MeV) of the leptons channels for their annihilation by means of neutral weak currents, $L^0\bar{L}^0 \to e^+e^-, \ldots$, are opened, the annihilation cross section increasing moreover with the lepton mass as $\sigma = G_F^2 m_\nu^2/\pi$ (in units of $\hbar = c = 1$), where $G_F = 10^{-5}/m_p^2$ is the weak coupling constant. The annihilation of leptons at the time when they come out of thermodynamic equilibrium leads[6-8] with this mass dependence of the annihilation cross section to a decrease of their residual concentration as $\propto m_\nu^{-3}$, which decreases the contribution of the leptons to the density of the Universe, which is $\propto m_\nu^{-2}$.

The weak interaction of nonrelativistic heavy leptons with matter depends on the kinetic energy of the leptons. After the time of "freezing" of leptons with $m_\nu > 1$ MeV the kinetic energy of the leptons satisfies $T \ll m_\nu$, but there is a large number of electrons and positrons ($n_e \approx n_e \approx n_\gamma$) in the matter. Therefore, weak scattering processes are possible, but when $t \gtrsim 0.1$ sec the characteristic scattering time becomes longer than the cosmological time, i.e., the leptons come out of thermal equilibrium. Thereafter, the lepton gas effectively interacts only gravitationally with the remaining matter.

The gravitational interaction of a nonrelativistic gas of massive neutrinos with matter during the galaxy-formation stage can lead to the condensation of neutrinos in galaxies and clusters of galaxies. According to Ref. 10, such condensation is ineffective for neutrinos with $m_\nu \lesssim 30$ eV and cannot explain the hidden mass in galaxies and clusters of galaxies. In what follows, we shall consider the astrophysical consequences of the existence of heavy neutral leptons with $m_\nu \gtrsim 2$ GeV.

Depending on the ratio of the contributions of the contributions of the leptons and ordinary matter to the cosmological density, two mass intervals can be distinguished: $2 \lesssim m_\nu \lesssim 10$ GeV, for which the contemporary cosmological density is determined by the density of the leptons, and $m_\nu \gtrsim 10$ GeV, for which the lepton density is lower than the density of ordinary matter.

Gravitational instability begins to develop in the lepton gas immediately after it has come out of thermal equilibrium with the matter; in the first case, the end of the radiation-dominated phase and the subsequent time of recombination are determined by the time when the radiation density is comparable with the lepton density. It is interesting that under conditions when the gravitational instability is determined by the spectrum of perturbations of the lepton gas the wavelengths of growing perturbations can be significantly shorter[1)] than for ordinary matter, since the neutrino and photon mechanisms of damping of short-wavelength perturbations are ineffective for the lepton gas. Therefore, besides the growing long-wavelength perturbations, which, according to modern ideas, result in the formation of the large-scale structure of the Universe, one can also have growth of perturbations which lead to an initial small-scale structure. However, since the expansion rate during the radiation-dominated stage is determined by the radiation density, the perturbations do not have sufficient time to increase appreciably before recombination.

In the case $m_\nu \gtrsim 10$ GeV, the presence of small-scale perturbations of the lepton gas cannot be expected to have a significant influence on the structure of the inhomogeneities of the Universe.

The condensation of heavy leptons with $m_\nu \gtrsim 2$ GeV in clusters of galaxies was considered in Ref. 8. It is asserted there that heavy leptons can be carried along by the collapsing matter only until the energy-dissipation mechanisms in the collapsing matter begin to operate. Since the collapsing heavy leptons do not radiate energy, their collapse would have to stop, and they would have persisted to the present time in the form of extended halos of the galaxies and clusters of galaxies.

It is shown in the present paper that the motion of the leptons in the nonstationary field of the contracting matter ensures an effective mechanism for dissipating their energy. Because of this, the contraction of the ordinary matter entails contraction of the lepton gas as well.

The increase in the density of leptons with $m_\nu > 2$ GeV during the galaxy-formation stage leads to an increase in their annihilation rate. Although only a small fraction of all the leptons participate in the annihilation, the observational consequences of such annihilation make it possible, on the basis of the date on cosmic-ray and γ-ray fluxes, to increase appreciably the number limit for the permitted mass of a heavy lepton compared with the cosmological limit $m_\nu > 2$ GeV. These bounds can be made much stronger (see also Ref. 6) if the leptons contract with the matter when stars are formed.

All the obtained results also hold for all other neutral stable massive particles, for example, for the lightest hadrons consisting of an absolutely stable new heavy quark.[11] Indeed, because of the asymptotic freedom of the strong interactions, annihilation of such hadrons via the strong interaction could be less probable than weak annihilation. The scattering of such hadrons would be determined by the strong interaction, which would significantly reduce their mean free path in matter compared with the case of heavy leptons and could lead to their condensation in the nuclei of galaxies or to the "adhesion" of the gas of such hadrons to the gas of the ordinary matter.

2. ENTRAINMENT OF LEPTONS BY CONTRACTING MATTER

When the ordinary gas ("baryons") radiates and contracts, the neutral heavy leptons ("leptons") move in a potential that depends on the time. Since the energy of a particle moving in a time-dependent potential is not in general conserved, the leptons have the possibility of decreasing their energy and, therefore, increasing their density, i.e., they can be carried along, or entrained, by the contracting baryons.

We consider the simplest case when the particles move along radial orbits in the central region, where the density does not depend (or depends only weakly) on the radius. If $\rho_\nu(t)$ and $\rho_b(t)$ are the densities of the leptons and baryons at the center, the motion of the leptons is determined by the equation

$$\ddot{r}=-\omega^2(t)r, \tag{1}$$

where

$$\omega(t)=\sqrt{{}^4/_3\pi G(\rho_\nu(t)+\rho_b(t))}.$$

Suppose the baryons slowly increase their density. For slowly varying ω, the amplitude of the oscillations is, as is well known, determined by the adiabatic invariant

$$A^2\omega(t)=A_0\omega(0). \tag{2}$$

With increasing ω, the amplitude of the lepton oscillations decreases and, therefore, their density increases:

$$\frac{\rho_\nu(t)}{\rho_\nu(0)}=\left(\frac{A_0}{A}\right)^3=\left(\frac{\omega(t)}{\omega(0)}\right)^{3/2}=\left(\frac{\rho_\nu(t)+\rho_b(t)}{\rho_\nu(0)+\rho_b(0)}\right)^{3/4}. \tag{3}$$

Equation (3) [$\rho_\nu(t)$ occurs on both sides of the equation] can be conveniently rewritten in somewhat modified form. We use the notation

$$x(t)=\rho_b(t)/\rho_\nu(t). \tag{4}$$

Then Eq. (3) has the form

$$x(1+x)^3=\frac{\rho_b(t)}{\rho_b(0)}x(0)(1+x(0)). \tag{5}$$

Two limiting cases are possible. In the first case, the expression on the right-hand side of (5) is small compared with unity. It is easy to show that this entails

$$\rho_b(t)\ll\rho_\nu(t), \tag{6}$$

i.e., the density of the baryons increases but still remains less than the initial density of the leptons. At the same time, the lepton density also increases:

$$\frac{\rho_\nu(t)}{\rho_\nu(0)}\approx1+3\frac{(\rho_b(t)-\rho_b(0))}{\rho_\nu(0)}. \tag{7}$$

It is noteworthy that in this case the lepton density increases by an amount that is greater than for the baryon density:

$$\Delta\rho_\nu=\rho_\nu(t)-\rho_\nu(0)=3(\rho_b(t)-\rho_b(0)), \tag{8}$$

$$\Delta\rho_b=\rho_b(t)-\rho_b(0). \tag{9}$$

However, with further increase in the baryon density ρ_b becomes equal to ρ_ν and can even exceed it somewhat, though the ratio ρ_b/ρ_ν increases very slowly with increasing ρ_b. Indeed, if the inequality (6) has the opposite sign, then

$$\frac{\rho_\nu(t)}{\rho_\nu(0)}=\left(\frac{\rho_b(t)}{\rho_\nu(0)+\rho_b(0)}\right)^{3/4} \tag{10}$$

and therefore

$$\frac{\rho_b(t)}{\rho_\nu(t)}\propto\rho_b^{1/4}(t). \tag{11}$$

The relations (10) and (11) also hold for motion of the leptons in circular orbits:

$$rv=\text{const},\quad v^2\approx G\rho r^2,\quad r^4\rho^{3/4}=\text{const}, \tag{12}$$

$$\frac{\rho_\nu(t)}{\rho_\nu(0)}=\left(\frac{r_0}{r}\right)^3=\left(\frac{\rho_b(t)}{\rho_b(0)+\rho_\nu(0)}\right)^{3/4}.$$

The case corresponding to the inequality (6) can be investigated even if the baryon density varies rapidly with the time. Solving Eq. (1) by successive approximations (taking at the start $\rho_\nu(t)=\rho_\nu(0)$, etc.), one can readily find that in the first order in

$$[\rho_b(t)-\rho_b(0)]/\rho_\nu(0) \tag{13}$$

the growth of the lepton density is determined by the relation (7), i.e., the relation (7) is also valid for the case of rapid growth of the baryon density.

Note that the approximation (6) and (7) can be realized only if the mean density in the Universe is determined mainly by the leptons. Otherwise, the asymptotic behavior (10) will be valid from the very start.

These analytic estimates show that with increasing baryon density the lepton density also increases by almost the same amount. It is important to establish that this conclusion is also valid in the general case of nonradial motion and inhomogeneous density (it is, in fact, easy to show that in the first approximation inhomogeneity does not change the result but leads only to a decrease in ω). To this aim, a number of numerical models were investigated. We used a method that makes it possible to study the evolution of a system of point self-gravitating masses. This method has been frequently used to simulate some plasma phenomena[12-14] and to simulate galaxies.[15,16] In its general features, the method is as follows.[17] If at a certain time the velocities and coordinates of the particles are given, their values at a subsequent time are found by a procedure which consists of three steps: 1) the density at the mesh sites is found by the clouds-in-cells method, 2) a fast Fourier transformation is used for the finite-difference Poisson equation, 3) the accelerations and new values of the coordinates and velocities are determined.

Energy dissipation by the baryons (cooling) was stimulated as follows. A group of particles was chosen from the complete system of particles that up to the onset of cooling was in equilibrium. If the velocity of any

TABLE I. Initial and final values of the central density in model 1.

Group of particles	Density at the center		Ratio of final density to initial density
	before cooling	after cooling	
1000 ("baryons")	2.8	20	7,1
4000 ("leptons")	11.2	22	2,0
5000 (total)	14	42	3.0

TABLE II. Change in the density in model 2.

Group of particles	Beginning		End	
	$\rho_0(r\leqslant 4)$	$\rho(9\leqslant r\leqslant 10)$	$\rho_0(r\leqslant 4)$	$\rho(9\leqslant r\leqslant 10)$
460 ("baryons")	0,0164	0,0055	0.90	0,008
9540 ("leptons")	0,222	0.057	0.99	0,268
Ratio of densities	13,5	10,4	1,1	33,5

of the chosen particles exceeded a definite limit, it was reduced abruptly to a definite value.

To show that the final result is independent of the employed method of energy dissipation, we investigated variants with different parameters of the cooling and with different initial stationary distributions. In all the considered variants, the calculations lead to equality of the densities of leptons and baryons to accuracy 20–50% in the central region of the system, the baryon density decreasing much more steeply than the lepton density with increasing distance from the center. As an illustration, we give the results of calculations of two models.

In the first model, we considered 5000 points and artificially cooled 1000 randomly chosen points. The values of the central densities before and after the cooling are given in Table I in relative units, and the nature of the distribution of the surface density is shown in Fig. 1. In the second model, 460 points out of 10 000 were effectively cooled. In Table II, we give the initial and final values of the density at the center and at a certain distance from the center (where the total density has fallen by approximately seven times).

Thus, the analytical and numerical calculations confirm that the heavy neutral leptons are carried along by the contracting, self-gravitational gas.

3. BOUNDS ON THE LEPTON MASSES

The entrainment of the leptons by the collapsing matter described in the previous section can occur both during the formation of galaxies and during the formation of stars. However, in the latter case the effectiveness of the mechanism is significantly determined by the details of the star-formation process. If entrainment of the leptons by the collapsing matter is to take place during the star-formation stage as well, it is necessary that at all stages the contraction of the protostars be governed by self-gravitation. If at some stage the principal part is played, not by self-gravitation, but, for example, by external pressure (which is very probable), then virtually all the leptons must remain in the interstellar medium. However, for condensation of leptons in the Galaxy the considered mechanism is fairly efficient.

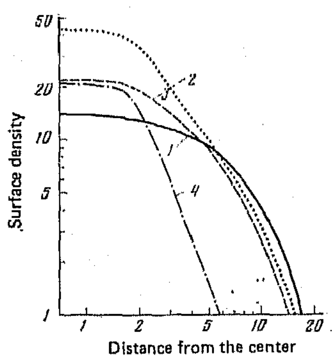

FIG. 1. Dependence of surface density on the radius (in relative units). Curve 1 is the initial distribution of the total density, curve 2 the final distribution of the total density, curve 3 the final distribution of the "leptons," and curve 4 the final distribution of the "baryons."

When the density of the leptons increases, the rate of their annihilation also increases. For $m_\nu \ll M_Z$, where M_Z is the mass of the heavy neutral intermediate vector boson for the neutral currents corresponding to the Weinberg–Salam model,

$$\sigma_{an}v = \frac{G_F^2 m_\nu^2}{2\pi} N, \tag{14}$$

where N is the number of possible weak-annihilation channels. For $m_\nu \lesssim 2$ GeV, $N = 14$ (Refs. 6–9); for $m_\nu > 2$ GeV, the channels of annihilation into $\tau^+\tau^-$ and charmed particles are opened; for $m_\nu > 5$ GeV, the channels with production of hadrons with a heavy b quark; at larger masses, channels associated with possible new leptons and quarks, which are predicted by modern theories of elementary particles, could be opened. The chains of decays of the particles produced by the annihilation process lead ultimately to γ rays, electrons and positrons, and nucleons and antinucleons. Therefore, an observable consequence of the annihilation of heavy leptons would be an increase in the cosmic-ray fluxes.

Let us estimate the yield of cosmic rays. The rate of annihilation of leptons in a unit volume is given by

$$\frac{dn_{an}}{dt} = n_{\nu G}^2 \sigma_{an} v, \tag{15}$$

where $n_{\nu_G} = n_{\bar\nu_G}$ is the density of leptons (and antileptons) in the Galaxy. The density n_{ν_G} is related to the cosmological residual-lepton concentration n_ν by (for $\rho_b > \rho_\nu$)

$$n_{\nu G} = n_\nu(\rho_G/\rho_b), \tag{16}$$

where ρ_G is the mean density of matter in the Galaxy and ρ_b is the density of matter in the Universe. For the residual lepton concentration, we have in accordance with Refs. 6–9

$$n_\nu = {}^4/_{11} n_\gamma r_\infty, \tag{17}$$

where[6]

$$r_\infty = 3\cdot 10^{-7}(m_p/m_\nu)^3. \tag{18}$$

Assuming that in each annihilation k_γ photons with energies $\gtrsim 100$ MeV are produced by the decay of π^0 and other particles, we obtain for the γ-ray flux

$$F(E_\gamma \geqslant 100\,\mathrm{MeV}) = \frac{n_{\nu G}^2 \sigma_{an} v k_\gamma V}{4\pi R_G^2}\frac{3}{8\pi} = \frac{n_\nu n_{\nu G}}{\rho_b}\frac{M_G}{4\pi R_G^2}\sigma_{an} v k_\gamma \frac{3}{8\pi}, \quad (19)$$

where $R_G \simeq 10$ kpc, and $M_G \simeq 10^{11}\,M_\odot = 2\times 10^{44}$g. For $k_\gamma \approx N$, $N=18$, $\rho_G \approx 10^{-23}$ g/cm^3, $\rho_b \simeq 2\times 10^{-30}$ g/cm^3 ($\Omega_b = 0.1$ for $H = 100$ km$\cdot$sec$^{-1}\cdot$Mpc^{-1}), $n_\nu \simeq 5\times 10^{-5}(m_p/m)^3$, and $\sigma_{an} v = 3.5\times 10^{-27}(m_\nu/m_p)^2$ cm^3/sec, we obtain

$$F(E_\gamma \geqslant 100\,\mathrm{MeV}) \simeq 0.77\,(m_p/m_\nu)^4\ (\mathrm{cm}^2\cdot\mathrm{sec}\cdot\mathrm{sr})^{-1}. \quad (20)$$

The experimentally observed isotropic background of γ rays at $E_\gamma \gtrsim 100$ MeV is[19]

$$F_{\exp}(E_\gamma > 100\,\mathrm{MeV}) = (1.0\pm 0.4)\cdot 10^{-5}\ (\mathrm{cm}^2\cdot\mathrm{sec}\cdot\mathrm{sr})^{-1}.$$

Thus, we obtain the bound

$$m_\nu > 15\ \mathrm{GeV}. \quad (21)$$

If we use data on the electron and proton components of the cosmic rays, we can obtain an even stronger bound on the mass of a heavy lepton. Taking the lifetime of cosmic rays in the Galaxy to be $T_{cr} \simeq 10^8$ years, we obtain for the flux of cosmic rays from lepton annihilation

$$F = \frac{dn_{cr}}{dt} T_{cr} \cdot \frac{c}{4}\frac{k_{cr}}{2\pi}, \quad (22)$$

where k_{cr} is the number of relativistic protons or electrons produced by one annihilation. Assuming approximately that in each annihilation there is produced an electron or proton with energy $\sim m_\nu$, we obtain

$$F = 800\,(m_p/m_\nu)^4\ (\mathrm{cm}^2\cdot\mathrm{sec}\cdot\mathrm{sr})^{-1}. \quad (23)$$

From the condition that there are no distortions in the region of $E = m_\nu$ in the observed cosmic-ray spectrum $J \approx E^{-2.6}$ particles$\cdot$cm$^{-2}\cdot$sec$^{-1}\cdot$sr$^{-1}\cdot$GeV^{-1} (in the range $10 < E < 10^6$, where E is measured in GeV) (Ref. 20) we obtain

$$m_\nu \geqslant 100\ \mathrm{GeV}. \quad (24)$$

This bound covers almost completely the admissible range of masses of heavy neutral leptons.

We note first of all that in unified theories of the weak and electromagnetic interactions based on the Higgs mechanism of spontaneous symmetry breaking there are upper bounds on the masses of heavy leptons: $m_\nu \lesssim M_Z$, where $M_Z \sim 100$ GeV is the mass of the intermediate boson of the weak interaction. This bound arises because of the condition that the presence of heavy leptons must not lead to a catastrophic rearrangement of the vacuum of the theory.[21] The bound is obtained under the assumption that are no superheavy Higgs mesons. If the main contribution to m_ν is determined by such Higgs mesons, the restriction of Ref. 21 is lifted and leptons with $m_\nu > M_Z$ could exist. However, formula (14) is valid only for contact weak interaction. When $m_\nu \gtrsim M_Z$, it is important to take into account the propagator of the Z boson, and the growth in the annihilation cross section with increasing m_ν must cease. It is to be expected that for $m_\nu \gg M_Z$, if such leptons exist at all, the weak-annihilation cross section begins to decrease with increasing mass of the leptons:

$$\sigma v \sim (g^2/m_\nu^2) N, \quad (25)$$

where $g^2/4\pi \sim \alpha$, in which g is the dimensionless (in units of $\hbar = c = 1$) coupling constant of the Z boson. Since the annihilation cross section (25) decreases with increasing lepton mass, the residual concentration of such leptons increases with their mass as $\propto m_\nu$. Therefore, the rate of their annihilation in the Galaxy will increase accordingly.

Because of this, for $m_\nu > 100$ GeV the contradiction with the observations remains and may even be stronger, since an increase in the contribution of heavy leptons to the matter density conflicts with data on the density of matter in the Universe.

It is interesting that in the case $m_\nu \sim M_Z/2$ the lepton-annihilation cross section increases strongly in resonance and $\sigma v \sim g^4/\Gamma^2 \sim \alpha^2/\Gamma^2$, where $\Gamma \sim g^2 M_Z \sim \alpha M_Z$ is the total width of the Z boson. Then a calculation analogous to the one in Refs. 5–8 gives

$$r_\infty \approx 10^{-10}(m_\nu/m_p)$$

and for the cosmic-ray flux we have $F \approx 10^{-5}$ particles $\cdot$cm$^{-2}\cdot$sec$^{-2}\cdot$sr^{-1}, which at cosmic-ray energies $E < 100$ GeV does not exceed the observed flux. Therefore, there remains the single interesting possibility that the mass of the heavy stable neutral lepton is $m_\nu = M_Z/2$.

4. CONCLUSIONS

The mechanism of entrainment of heavy neutral, absolutely stable leptons by collapsing matter proposed in this paper leads to severe restrictions on the mass of such leptons. In conjunction with bounds obtained earlier, they are an argument against the existence of heavy stable neutral leptons with mass $m_\nu > 30$ eV. This conclusion is also fully applicable for other neutral stable particles which interact weakly with matter; moreover, the weaker their interaction, the higher their residual concentration and, therefore, the stronger is the contradiction with the observations.

The results obtained here, when taken with the results of Ref. 10, preclude the possibility of explaining the hidden mass of galaxies and clusters of galaxies by the presence of a gas of massive neutrinos or any other particles that interact weakly with matter.

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[1)]Because the leptons do not get out of thermal equilibrium instantaneously, perturbations with wavelength shorter than the horizon distance ($M < (10^3 - 10^4)\,M_\odot$) could be suppressed at $t \sim 0.1$–1 sec.

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Production of charmed particles in νN collisions due to neutral weak currents

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A study is made of associated production of charmed particles in neutrino-nucleon interactions due to neutral weak currents. The angular distribution of the jets of charmed hadrons in νN interactions is determined in the lowest approximation in the quark-gluon coupling constant, according to which a charmed quark and antiquark are produced in an annihilation of a vector gluon and a virtual Z boson. It is shown that only a P-even dependence on the azimuthal angle φ occurs in the studied approximation, the P-odd dependence which is possible in the general case being equal to zero. The total cross section for charmed-particle production in neutrino-nucleon interactions is calculated, and the origin of the violation of scale invariance is demonstrated.

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1. In the present paper, the techniques of quantum chromodynamics are used to study associated production of charmed particles in the process $\nu+N\rightarrow\nu+$hadrons. It is assumed that the production of charmed particles is determined by the neutrino-gluon interaction $\nu+g\rightarrow\nu+c+\bar{c}$ (Fig. 1), where g is a gluon and c is a charmed quark. This model can be justified by reference to the theory of the strong interaction based on quantum chromodynamics,[1-3] according to which the production of heavy particles should be determined by the lowest order in the quark-gluon coupling constant (by virtue of the property of asymptotic freedom of the gauge theory). The assumed mechanism seems to be the most natural source of associated production of charmed particles (for example, $D+C$) or particles with hidden charm, such as the ψ particles. The presence of gluons in nucleons follows, in particular, from the fact that the valence quarks carry only half of the longitudinal momentum of the nucleon.[4] The proposed gluon mechanism does not require the introduction of pairs of heavy $c+\bar{c}$ quarks in the "sea" of quark-antiquark pairs.

As is well known, the gluon mechanism is most effective for the production of charmed particles in hadron-hadron,[5-8] eN,[9] and γN (Ref. 10) interactions, and also in the process $\nu+N\rightarrow\mu^-+X$.[11] Since the gluons have identical strong interactions with all the quarks (independently of their flavor), this mechanism will also lead (at appropriate energies) to the production of heavier quarks, for examble, $b+\bar{b}$ and $t+\bar{t}$. We note that in the process under consideration here, $\nu+N\rightarrow\nu+$charmed particles, the gluon mechanism makes it possible in principle to study the properties of the neutral weak current of the charmed quark (and also of the heavier quarks).

We make a detailed study of the angular distribution of charmed jets, which is of exceptional interest for testing the predictions of quantum chromodynamics. For this purpose, we calculate the standard contributions to the cross section for charmed-particle production when, in addition to neutrinos in the final state, quark and antiquark jets are detected. By carrying out an integration with respect to the momenta of the $c+\bar{c}$ pair, we find the inclusive cross section for charmed-particle production in neutrino-nucleon interactions.

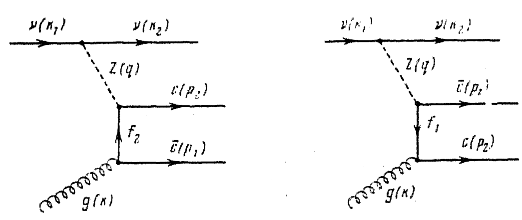

FIG. 1.

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