Bounds on very heavy relic neutrinos by their annihilation in the Galactic halo

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Taking into account neutrino condensation in the gravitational field of collapsing matter, we investigate the annihilation of heavy relic neutrinos in the Galaxy resulting in the generation of cosmic rays. The main neutrino annihilation processes are considered: i.e., $\nu\bar{\nu} \to f\bar{f}$ and $\nu\bar{\nu} \to W^+W^-$. The condensation mechanism allows one to get information on the density distribution in the Galaxy halo without any recourse in an explicit dynamical halo model, and the resulting cosmic ray spectrum provides constraints on the heavy neutrino mass. The comparison of the predicted cosmic ray flux with the observed one excludes the heavy neutrino mass range 60 GeV $< m_\nu <$ 115 GeV. Such a restriction leads to a bound on the present energy density of very heavy neutrinos which may be comparable to the corresponding baryonic one only in the range 115 GeV $< m_\nu <$ 300 GeV. Our approach is valid for multicomponent dark matter and can be used for species that give even a negligible contribution to the critical cosmological density.

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I. INTRODUCTION

By measuring the width and height of the Z-boson peak CERN e^+e^- collider LEP experiments [1] tell us that there are three neutrino species. However, this constraint applies only to light neutrinos with a mass $m_\nu < M_Z/2$, where M_Z is the mass of the Z boson, and therefore does not forbid the existence within the framework of the standard model of very heavy neutrinos, so that Z-boson decays into them are prohibited by phase space. In particular, the results of modern experiments are not inconsistent with the existence of heavy Dirac neutrinos with $m_\nu > 44$ GeV [2].

The important additional source of information on neutrino parameters is cosmology. According to the theory of the hot universe, in the Universe there should exist a background of relic neutrinos, whose concentration is related to that of relic photons. The energy density ρ_ν of massive stable neutrinos should be smaller than the total critical energy density ρ_c of the Universe. Assuming that the cosmological density does not exceed the critical one,

$$\rho_\nu \cong \sum_\nu (m_\nu/100 \text{ eV}) H^{-2} \rho_c < \rho_c \ ,$$

where H is the present Hubble expansion rate in units of 100 km s^{-1} Mpc^{-1}, it was found [3] that the allowable ranges for the neutrino mass are

$$m_\nu < 30 \text{ eV and } 3 \text{ GeV} < m_\nu < 3 \text{ TeV} \ .$$

In these calculations a range of large neutrino masses is available as a result of the rising cross section of neutrino annihilation with an increase of the neutrino mass as $\sigma \sim m_\nu^2$ for 1 GeV $\leq m_\nu < \frac{M_Z}{2}$. This leads to a decrease of both the residual neutrino concentration as $n_\nu \sim 1/(m_\nu \sigma v) \sim m_\nu^{-3}$ and of the neutrino energy density in the Universe. But beyond the Z pole, the annihilation cross section for the process $\nu\bar{\nu} \to f\bar{f}$ (where f is a fermion) starts decreasing as m_ν^{-2} as a result of the momentum dependence of the Z-boson propagator. In this region the relic number is proportional to m_ν and the neutrino energy density increases as m_ν^2, reaching again the critical value ρ_c for a few TeV.

However, these calculations were carried out without taking into account the important annihilation channel $\nu\bar{\nu} \to W^+W^-$. As was shown in Ref. [4], in the region of neutrino masses $m_\nu > M_W$ this process defines the fastest rate of the neutrino annihilation with a cross section $\sigma \sim m_\nu^2$. This leads to a corresponding decrease in the neutrino energy density $\rho_\nu \sim m_\nu^{-2}$, and therefore there is no longer a cosmological upper bound on the stable neutrino mass based on the consideration of the total energy density in the Universe. It should be noticed that above 10 TeV (possibly within 100 TeV), the unitarity principle should manifest somehow so that the cross section should reach an upper bound and deviate from the analysis considered below. This extreme region will be discussed elsewhere [5].

Nevertheless, the nonuniform distribution of massive

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neutrinos in the Universe and the local increase of the neutrino density in the formation of the galactic halo lead to considerable sensitivity of astrophysical data to the existence of heavy neutrinos in the Universe. As was shown first in Ref. [6], the condensation of neutrinos in the Galaxy should speed up their annihilation, thus resulting in the generation of cosmic rays. Therefore the comparison of the predicted cosmic ray flux due to such annihilation with the observed flux allows one to obtain nevertheless some restrictions on the heavy neutrino mass. The effect of the neutrino condensation caused by the gravitational binding of heavy neutrinos in the gravitational field of collapsing gas at the stage of Galaxy formation (which was discovered in Ref. [6]) was used in Ref. [7] for an analysis of the implications of the weak annihilation of supersymmetric relic particles in the Galaxy.

In the present paper we carry out a detailed analysis of the influence of the effects of the very heavy neutrino annihilation ($m_\nu > 44$ GeV) on cosmic ray production in the Galaxy. To be explicit, we consider the standard electroweak model, including, however, one additional family of fermions. Then the heavy neutrino ν and heavy charged lepton L form a standard $SU(2)_L$ doublet. In order to ensure the stability of the heavy neutrino ν, we assume that $m_\nu < M_L$ and that the heavy neutrino is a Dirac neutrino. The organization of the paper is the following. In Sec. II we compute the residual relic concentration of heavy neutrinos taking into account the main processes of their annihilation, i.e., $\nu\bar{\nu} \to W^+W^-$, $\nu\bar{\nu} \to f\bar{f}$. In Sec. III the distribution of collisionless particles in galaxies is considered. In Sec. IV we present the computation of the cosmic ray spectra due to neutrino annihilation and investigate the possible constraints on the mass of stable heavy neutrinos. In Sec. V we analyze the role of the Higgs meson in the cosmological constraint on the neutrino mass. In Sec. VI we summarize the main astrophysical and cosmological consequences of the present paper.

II. RESIDUAL RELIC CONCENTRATION OF HEAVY NEUTRINOS IN THE UNIVERSE

In the early Universe at high temperatures ($T \gg m_\nu$), heavy neutrinos (if they exist) should be in thermal equilibrium with other kinds of particles and their concentration should be compatible with that of photons. As the temperature in the Universe drops, neutrinos become nonrelativistic at $T \sim m_\nu$ and their abundance falls off rapidly according to $n_\nu \sim \exp(-m_\nu/T)$, although they are still in thermal equilibrium with the other particles. However, in the further expansion of the Universe, as the temperature drops below the freeze-out temperature T_f, the weak interaction processes become too slow to keep neutrinos in equilibrium with the other particles. The equilibrium is destroyed when the rate of change of the equilibrium concentration due to the temperature decrease turns out to be comparable to the rate of the equilibrium reactions. As a consequence, the concentration of heavy neutrinos fails to follow the equilibrium concentration and the exponential drop of the concentration becomes much slower.

The residual relic concentration of heavy neutrinos in the Universe is given, according to Ref. [8], by

$$n_\nu = (4/g_*)(2Tr_f/m_\nu)n_\gamma \;, \tag{1}$$

where $n_\gamma = 0.24T^3$ is the equilibrium photon concentration, $T = 2.7$ K is the present photon temperature, the factor $4/g_*$ takes into account the increase of the photon temperature due to the annihilation of the particles after the quenching of heavy neutrinos, $g_* = g_*(T) = N_b + \frac{7}{8}N_f$ is the number of effective degrees of freedom (the photon contribution to g_* is 2), the quantity

$$r_f = (g_s/4)(T_f/m_\nu)^{-3/2}\exp(-m_\nu/T_f) \tag{2}$$

is the relative equilibrium concentration of heavy neutrinos at the moment of quenching, and g_s is the number of particle spin states (for photons and massive fermions, $g_s = 2$). The freeze-out temperature is defined by

$$m_\nu/T_f - \tfrac{1}{2}\ln(T_f/m_\nu) = 40 + \ln(\sigma v M_p^2) + \ln(m_\nu/M_p) + \ln(g_s/g_*^{1/2}) \;, \tag{3}$$

where M_p is the proton mass and v is the neutrino relative velocity (in units of the velocity of light).

Taking into account Eq. (3), expression (1) for the relic concentration of heavy neutrinos can be written as

$$n_\nu \cong 2\times 10^{-18} g_*^{-1/2}(M_p/m_\nu)(\sigma v M_p^2)^{-1}[40 + \ln(\sigma v M_p^2) + \ln(m_\nu/M_p) + \ln(g_s/g_*^{1/2})]n_\gamma \;, \tag{4}$$

where the condition $T_f \ll m_\nu$ is assumed.

In order to calculate n_ν, we have to know the total annihilation cross section of neutrinos at freeze-out. At that moment the heavy neutrinos are nonrelativistic [8], and so in the annihilation their total energy is $\sqrt{s} \cong 2m_\nu$. The analysis of the annihilation reactions of heavy neutrinos shows the following.

(1) In the region of neutrino masses $m_\nu > 100$ GeV, the reaction

$$\nu\bar{\nu} \to Z \to W^+W^- \;, \tag{5}$$

through s-channel Z-boson exchange, starts dominating the total cross section of the neutrino annihilation [4] only if the total energy $\sqrt{s}$ is not in the vicinity of the resonant Higgs boson peak.

(2) The relative contributions to the process (5) of the diagrams with heavy lepton exchange and of the interference diagrams are suppressed in comparison with the s-channel Z-boson exchange if ($\hat{m}_\nu^2 \gg 10$) or ($\hat{m}_\nu^2 \gg 1$, $M_l^2 \gg 10M_W^2$), where $\hat{m}_\nu = m_\nu/M_W$.

(3) In the region $\hat{m}_\nu > 2$, the annihilation cross section for reaction (5) is given approximately by

$$\sigma\beta \cong 6.8(\hat{m}_\nu^2 + 4) \text{ (pb)} \;, \tag{6}$$

where β is the neutrino velocity.

(4) The solution of Eq. (3) by an iterative procedure [4] yields for the freeze-out temperature the value

$T_f \cong m_\nu/30$.

We note that in Eq. (4) the expression in square brackets depends very weakly on the neutrino mass m_ν and can be replaced by a constant value. As a result of this approximation, we can write the residual relic concentration of heavy neutrinos in the form

$$n_\nu \cong 2 \times 10^{-9}/[\hat{m}_\nu(\hat{m}_\nu^2 + 4)] \ (\mathrm{cm}^{-3}) \ , \qquad (7)$$

if $\hat{m}_\nu > 2$. Therefore, in the region of large neutrino masses, because of the process (5), the residual neutrino concentration decreases fast enough as $n_\nu \sim m_\nu^{-3}$.

In the region of neutrino masses $m_\nu < 100$ GeV, the dominant annihilation process is neutrino annihilation into a fermion pair

$$\nu\bar{\nu} \to Z \to f\bar{f} \qquad (8)$$

through Z-boson exchange in the s channel. Let us write the total cross section for this process:

$$\sigma = \frac{NG^2M_W^4D_Z}{32\pi\,\cos^4\theta_W}\frac{1}{s}(\beta_f/\beta)\left[\frac{s^2}{2}(g_v^2+g_a^2)(1+\beta^2\beta_f^2/3) + 2(g_v^2-g_a^2)m_f^2(s-2m_\nu^2) + 4g_a^2 s m_\nu^2 m_f^2 M_Z^{-2}(sM_Z^{-2}-2)\right] , \qquad (9)$$

where θ_W is the Weinberg angle, G is the Fermi constant, N is a color factor ($N = 1$ for leptons and $N = 3$ for quarks), m_f is the fermion mass, $\beta = (1-4m_\nu^2/s)^{1/2}$, $\beta_f = (1-4m_f^2/s)^{1/2}$, g_v and g_a are the standard vector and axial vector constant, $D_Z = [(s-M_Z^2)^2 + \Gamma^2 M_Z^2]^{-1}$, and Γ is the Z-boson width. Here we neglected the Higgs boson contribution, which is reasonable either at $\sqrt{s} \cong 2m_\nu$, as a result of the suppression of Higgs boson exchange, or in the region of large Higgs boson mass $m_H \gg \sqrt{s}$.

The appearance in expression (9) of the term proportional to m_ν^2 is inevitable because of the axial vector part of the weak neutral current. However, the role of the additional terms in formula (9) is important only for heavy fermions, which can be produced in the annihilation of neutrinos with large masses. But in this case the dominant reaction of annihilation is the process $\nu\bar{\nu} \to W^+W^-$, so that the contribution of these terms to the residual neutrino concentration is not sensible. Thus, in particular, the actual values of the top quark mass and the masses of new quarks of fourth generation are not very important in the calculation of n_ν.

The results of the numerical calculations of the residual concentration of heavy neutrinos on the basis of Eqs. (3), (4), and (9) [and the formulas of Ref. [4] for reaction (5)] are presented in Fig. 1. The calculations were carried out at heavy lepton mass $M_L = 1$ TeV (however, the actual value of M_L is not very important) and at the top quark mass $m_t = 200$ GeV. The masses of the other heavy quarks and the Higgs boson mass were assumed to be above 1 TeV. The annihilation cross section was evaluated at

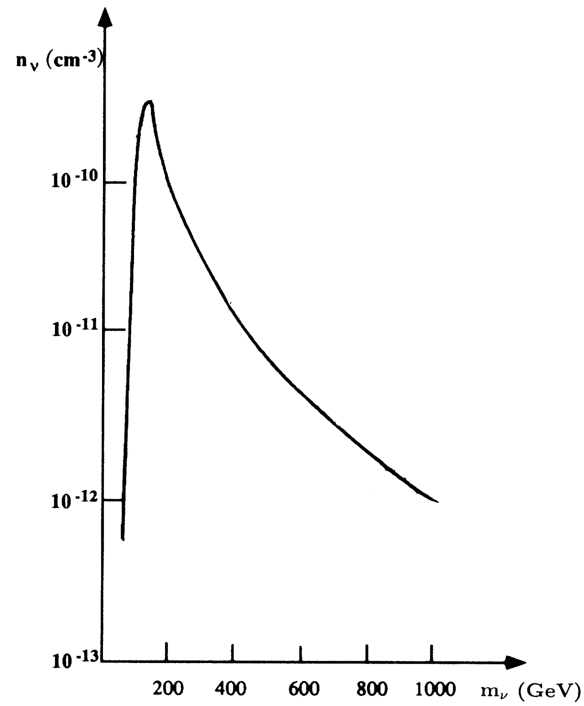

FIG. 1. Relic concentration of heavy neutrinos.

$$s = \langle s_f \rangle \cong 4m_\nu^2 + 6m_\nu T_f \ , \qquad (10)$$

since the annihilating neutrinos are nonrelativistic and constitute a Boltzmann-distributed gas, so that $\langle p^2 \rangle = 3m_\nu T_f$ [the results of the numerical calculations depend very weakly on the value of the second term in Eq. (10)].

Qualitatively, the behavior of n_ν (as shown in Fig. 1) is easily understood. In the region $m_\nu \sim M_Z/2$, the residual relic neutrino concentration is small as a result of the huge value of the cross section at resonance in the s channel. With an increase of the neutrino mass, the cross section for neutrino annihilation into fermions drops and this leads to an increase of the residual concentration. But at $m_\nu > M_W$ the additional annihilation channel $\nu\bar{\nu} \to W^+W^-$ opens and gradually becomes the dominant one, since its cross section grows like m_ν^2 (until $m_\nu < M_L$) and the residual concentration drops again.

III. DISTRIBUTION OF COLLISIONLESS PARTICLES IN GALAXIES

The very heavy neutrinos decoupled from radiation at very early epochs ($t < 10^{-8}$ sec), contrary to baryons, which remain in thermal equilibrium with photons until the last scattering at redshift $z \sim 1500$. However, as soon as the cosmological expansion is dominated by matter (for instance, by heavy neutrino or baryonic densities at $z \sim 10^4$–10^5), heavy neutrinos may gravitationally cluster into clouds, thus forming a primordial density seed where galaxies may later condense. After the recombination, once baryons become neutral ($z \sim 1500$), baryonic matter also plays a role in clustering and baryons "feel" the already existing primordial gravitational seeds in a coupled gravitational system [9]. Such a seed role (of cold dark matter) is able to speed up baryonic clustering and to reduce the primordial density contrast for baryons up to the observed bounds ($\Delta T/T \lesssim 10^{-4}$–$10^{-5}$). The subsequent baryonic galaxy formation is due to self-gravity and energy dissipation. The energy dissipation may amplify the galactic density contrast (with respect to the cosmological one) by many orders of magnitude.

At the stage of the formation of the Galaxy, neutrinos can interact with matter by gravitation only. Therefore no energy dissipation due to radiation takes place in the gas of heavy neutrinos as it is in the case of ordinary matter. Nevertheless, as was shown in Ref. [6], the motion of neutrinos in the nonstatic gravitational field of ordinary matter, which contracts as a result of energy dissipation via radiation, provides an effective mechanism of energy dissipation for neutrinos too. As a consequence, contracting ordinary matter induces the collapse of the neutrino gas and leads to the following significant increase in the neutrino (antineutrino) density in the Galaxy [6]:

$$n_{\nu G} = n_\nu \rho_G/\rho \ , \tag{11}$$

where $\rho_G \cong 5 \times 10^{-24}$ g/cm^3 is the average density of matter in the Galaxy and $\rho \cong 4 \times 10^{-31}$ g/cm^3 is the density of matter in the Universe.

For future applications let us consider the mechanism, suggested in [6], in more detail.

When ordinary matter ("baryons") contracts as a result of energy dissipation via radiation, neutral heavy leptons ("neutrinos") move in a potential, which varies with time. Since the energy of a particle moving in a time-variable potential is generally not conserved, neutrinos can reduce their energy and, consequently, increase their density.

To illustrate this mechanism, let us treat, following [6], the simplest case of particle motion along radial orbits in the central part of the contracting baryon system, where the density is independent of radius. If $\rho_\nu(t)$ and $\rho_b(t)$ are, respectively, the central densities of neutrinos and baryons, the motion of neutrinos is determined by the equation

$$d^2r/dt^2 = -\omega^2(t) r \ , \tag{12}$$

where

$$\omega(t) = \{(4\pi/3)G[\rho_\nu(t) + \rho_b(t)]\}^{1/2} \ .$$

Let baryons increase slowly their density. For slowly varying ω the amplitude of oscillations is provided by the adiabatic invariant

$$A^2\omega(t) = A_0^2\omega(0) \ . \tag{13}$$

As ω grows, the oscillation amplitude of neutrinos decreases and their density increases, respectively, according to the equation

$$\begin{aligned}\rho_\nu(t)/\rho_\nu(0) &= (A_0/A)^3 \\ &= [\omega(t)/\omega(0)]^{3/2} \\ &= \{[\rho_\nu(t) + \rho_b(t)]/[\rho_\nu(0) + \rho_b(0)]\}^{3/4} \ .\end{aligned} \tag{14}$$

Introducing the variable

$$x(t) \equiv \rho_b(t)/\rho_\nu(t) \ , \tag{15}$$

one can conveniently rewrite Eq. (14) in the form

$$x(1+x)^3 = [\rho_b(t)/\rho_b(0)]x(0)[1+x(0)]^3 \ . \tag{16}$$

Two limiting cases are possible. In the first one the expression on the right-hand side of Eq. (16) is small, which corresponds to the condition

$$\rho_b(t) \ll \rho_\nu(t) \ . \tag{17}$$

In this case the neutrino density grows as

$$\rho_\nu(t)/\rho_\nu(0) = 1 + 3[\rho_b(t) - \rho_b(0)]/\rho_\nu(0) \ , \tag{18}$$

whereas baryon density while growing remains still smaller than the initial density of neutrinos. However, because of radiation energy losses, baryon density can grow to make the condition (17) invalid, so that the opposite condition

$$\rho_b(t) \gg \rho_\nu(t) \tag{19}$$

holds. Then one obtains from Eq. (16) that the density of neutrinos grows with time as

$$\rho_\nu(t)/\rho_\nu(0) = \{\rho_b(t)/[\rho_\nu(0) + \rho_b(0)]\}^{3/4} \tag{20}$$

and therefore

$$\rho_b(t)/\rho_\nu(t) \sim \rho_b(t)^{1/4} \ . \tag{21}$$

The same Eqs. (20) and (21) hold for the motion of neutrinos along circular orbits, when [6]

$$\begin{aligned} rv &= \text{const} \ , \\ v^2 &= G\rho r^2 \ , \\ r^2\rho^{1/2} &= \text{const} \ , \end{aligned}$$

and

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

The analytical treatment given above was completed in [6] by numerical models, proving the same law of condensation [Eq. (21)] to be valid in the cases of nonradial and noncircular motions and also for inhomogeneous density. The result was shown to be independent of the details of the numerical methods used. So both analytical and numerical calculations [6] prove the conclusion on the condensation of collisionless gas in self-gravitating systems, while contracting as a result of radiation energy losses.

The considered process of condensation of a collisionless gas may take place in any collapsing system of ordinary matter, provided that at all stages contraction is dominantly supported by self-gravity. It is generally assumed that this condition is not satisfied at the initial stages of formation of objects smaller than globular clusters, at which the development of thermal instability and the effects of the outer pressure of the hot gas are dominant. So the considered mechanism should be effective in the course of galaxy formation, but does not seem to work in the process of the formation of globular clusters and smaller astronomical objects (stars, in particular).

If heavy neutrinos (or some other hypothetical particles) dominate the cosmological density, such a mechanism provides an explanation for the formation of massive halos of galaxies by these particles. It was used, e.g., in [10] for the scenario of a neutrino-dominated Universe to explain why massive neutrinos remain at the periphery of galaxies and do not contribute much to the density in the central parts of galaxies. In this case the assumption that hypothetical particles dominate in the galactic halo allows one, without any recourse to an explicit dynamical model of halo formation, to get estimates for the density distribution in a halo to analyze the particle distribution in the Galaxy and to evaluate possible effects of their weak annihilation (as it was done in [7] for supersymmetric particles). However, the universality of the mechanism [6] permits its application to the case of a small contribution of the hypothetical particles to the total density, thus providing a reasonable estimate for the expected distribution of the particles in the Galaxy and their possible effects, even if they do not play any significant role in the dynamics of halo formation. The actual distribution of collisionless particles deserves special analysis for this case, which will be done in a separate work. As an estimate, following [6], we used above, for the averaged central density of massive neutrinos, Eq. (11), linearly relating it to the averaged baryonic density in the Galaxy. One can check that such a relationship corresponds to the law of condensation (21) for the ratio of the central densities of baryons and neutrinos.

IV. SPECTRUM OF COSMIC RAYS

The condensation of heavy neutrinos in the Galaxy leads to an increase in the rate of the neutrino annihilation, resulting in a copious production of cosmic rays. The most stringent limit on the mass of heavy neutrinos can be therefore obtained by considering the electronic component of cosmic rays. Then, in this section, we shall evaluate the output of relativistic electrons by exploiting the following formula of the flux [6]:

$$J = (dn/dt)T_e(c/4)(\delta/2\pi)\ (\mathrm{cm^{-2}\,s^{-1}\,sr^{-1}})\ , \qquad (22)$$

where

$$dn/dt = n_{\nu G}^2 \sigma_G v \qquad (23)$$

is the rate of neutrino annihilation in the Galaxy per unit volume, σ_G is the cross section of the neutrino annihilation in the Galaxy, $T_e \cong 10^7$ yr is the lifetime of cosmic rays in the Galaxy, c is the velocity of light, and δ is the number of relativistic electrons with energy in the interval $E_e - \Delta E/2$ to $E_e + \Delta E/2$, which are produced in one act of the neutrino annihilation.

Substituting formulas (11) and (23) into Eq. (22) gives us a general expression for the flux of cosmic rays:

$$J \cong 3.5 \times 10^{12} n_\nu^2 \delta(\sigma_G \beta/\mathrm{pb})\ . \qquad (24)$$

In order to obtain constraints on the heavy neutrino mass, we consider first the annihilation channel $\nu\bar{\nu} \to e^+e^-$ in the Galaxy. Since heavy neutrinos in the Galaxy are nonrelativistic, i.e.,

$$E_\nu \cong m_\nu + m_\nu (v/c)^2/2\ , \qquad (25)$$

where $v = 300$ km/sec is the velocity of neutrinos in the Galaxy, then the ultrarelativistic electrons in the annihilation reaction (8) are produced practically monochromatic with $E_e \cong m_\nu$ (even for $m_\nu = 1$ TeV the energy spread is $\Delta E \cong 1$ MeV). Electrons can also be produced in the secondary processes from the decays of μ and τ leptons and quarks, but these processes contribute only to the soft part of the cosmic ray spectrum and therefore we shall neglect them.

The annihilation cross section for the process $\nu\bar{\nu} \to e^+e^-$ can be written, according to formula (9), as

$$\sigma_G \beta \cong 2.9 s M_W^2 D_Z \quad (\mathrm{pb})\ , \qquad (26)$$

and therefore we have, for the flux of cosmic electrons,

$$J_e \cong 10^{13} n_\nu^2 s M_W^2 D_Z\ , \qquad (27)$$

where n_ν is calculated at s_f [see Eq. (11)] and $\delta = 1$ is used. The experimental energy spectrum of cosmic electrons [11] integrated over the energy resolution ΔE of the detector is given by

$$J^{\mathrm{expt}} = 1.16 \times 10^{-2} E_e^{-2.6} \Delta E\ , \qquad (28)$$

where 3 GeV $\leq E_e \leq$ 300 GeV and $\Delta E \ll E_e$.

Results of numerical calculations of the electron flux (27) are presented in Fig. 2. Also, the experimental flux (28) is shown in this figure for the electron energy $E_e = m_\nu$ and energy resolution $\Delta E = 1$ GeV. As we can see, in this case the existence of heavy neutrinos is forbidden in the mass range

$$60\ \mathrm{GeV} < m_\nu < 115\ \mathrm{GeV}\ . \qquad (29)$$

The absence of the constraint in the mass range 44–60 GeV is due to the smallness of the relic neutrino concentration (the annihilation cross section of the neutrinos in the early Universe is very large in the vicinity of the Z-boson peak). The constraint in the region of large neutrino masses is a consequence of the rapid decreasing of the flux (27) $J_e \sim m_\nu^{-4}(\hat{m}_\nu^2 + 4)^{-2}$ in comparison with the experimental flux (28) $J^{\mathrm{expt}} \sim m_\nu^{-2.6}$.

One might argue that the bounds (29) are approximate, because of the approximate nature of the mechanism of neutrino clustering discussed in the previous section. Indeed, some care and a more cautious attitude must be exercised because of the complexity of the phenomenon of multifluid clustering [5–8] (and the related

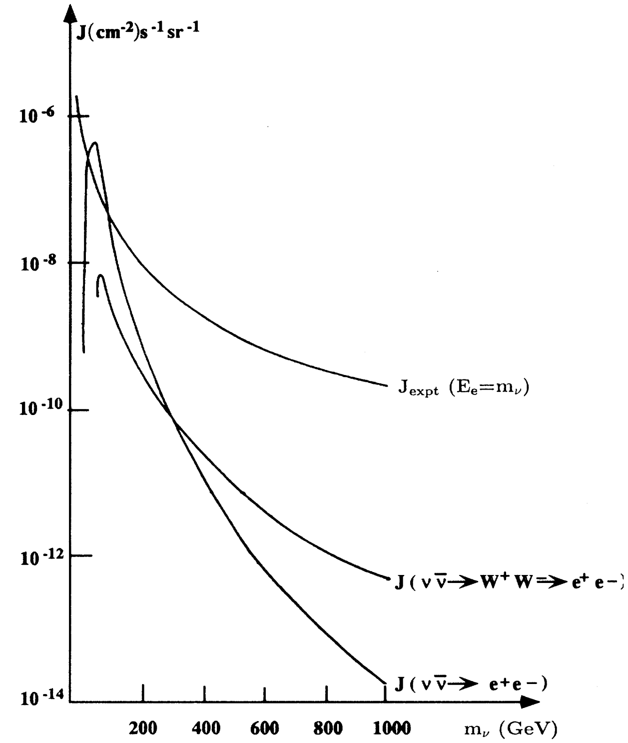

FIG. 2. Flux of cosmic electrons from the neutrino annihilation in the Galaxy as function of the neutrino mass.

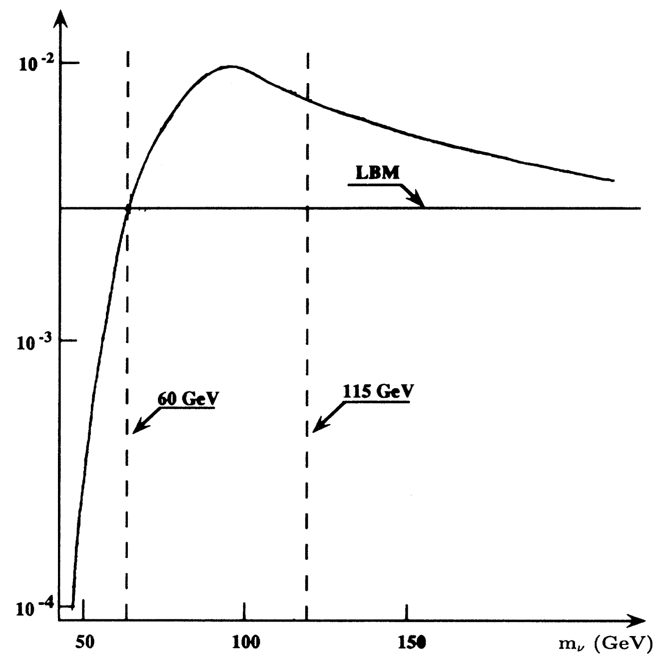

FIG. 3. Relic neutrino energy density versus neutrino mass, where $\rho_c = 10^{-29}$ g/cm^3 and LBM is the luminous baryonic matter contribution ($\Omega_b = 3 \times 10^{-3}$).

dissipation) in the process of galaxy formation. However, in the extreme case of a totally inefficient neutrino clustering (contrary to our expectation), the above bounds do not apply, but on the other side there is no longer any role for heavy neutrinos in dark galactic halos.

In our view, in most realistic scenarios, the neutrino clustering mechanism does constitute a reliable process, and in general it might be only partially corrected, thus leading to a slight enlarging (or narrowing) of the forbidden mass window given in Eq. (29).

Let us note that any bound on the neutrino mass does imply a corresponding limit on the present energy density in the Universe. This point can be seen from Fig. 3, where results of numerical calculations of the dimensionless neutrino energy density $\Omega_\nu = \rho_\nu/\rho_c$ at $\rho_c = 10^{-29}$ g/cm^3 are plotted. In the neutrino mass region below 60 GeV, the relic neutrino energy density ($\rho_\nu = 2m_\nu n_\nu$) is too small (below the luminous baryonic density $\Omega_b > 3 \times 10^{-3}$) to be of observational interest. Only in the region above 115 GeV could we expect significant contributions and cosmological implications of the neutrino component (comparable to the baryonic one).

We should note that the constraint (29) is very sensitive to the value of the energy resolution ΔE. In particular, the improvement of the resolution from 1 GeV to 1 MeV ($\Delta E = 1$ MeV is the width of the energy distribution of electrons produced in neutrino annihilation in the Galaxy) would give a relative gain of the order 10^3, thus allowing the removal of the constraint on m_ν in the region of 1 TeV and leading to a relic neutrino energy density below the baryonic one.

An important constraint on m_ν could be obtained, at least in principle, by investigation of a high-energy positron line from the annihilation reaction $\nu\bar{\nu} \to e^+e^-$ in the Galaxy. The flux of positrons in this case would be the same as for electron production (27). However, in the energy region where separate measurements of electrons and positrons are available ($E_e < 50$ GeV), a significant excess of electrons was found [12]. At higher energies we can use only theoretical predictions. If we assume, for instance, the validity of the model of dynamical halo, then it gives us the ratio $N(e^+)/N(e^-) < 10^{-2}$ at $E_e > 50$ GeV [12]. It follows from Fig. 2 that in this case it would be possible to search for heavy Dirac neutrinos with a mass up to 350 GeV. The detection of an anomalous output of positrons with energy above 50 GeV would be a clear signature of the annihilation of Dirac neutrinos in the Galaxy halo, because the annihilation of massive Majorana neutrinos into light fermions in the Galaxy is severely suppressed at low energies since, in the last case, it is required that the annihilation be in the p wave as a consequence of Fermi statistics. Let us note that for neutrino masses below M_W the positron line radiation (as a signature for heavy neutrino annihilation in the halo) was considered in detail in Ref. [13].

The other possibility to obtain a constraint on m_ν is to consider the process of neutrino annihilation $\nu\bar{\nu} \to W^+W^-$ and electron production from W-boson decays $W^- \to e^-\nu_e$. However, in this case there is a strong suppression factor

$$\delta \cong \Delta E/(9q) \tag{30}$$

of the flux of cosmic rays. The appearance of this factor is a consequence of the flat energy distribution of the electrons due to the relativistic motion of W bosons in reaction (5):

$$(q_0 - q)/2 \leq E_e \leq (q_0 + q)/2 , \tag{31}$$

where q_0 and q are the energy and the momentum, respectively, of the W boson. The coefficient $\frac{1}{9}$ in Eq. (30) is the W-branching ratio into the channel $W^- \to e^- \bar{\nu}_e$.

So we obtain the following approximate expression for the flux of cosmic rays (electrons) from the cascade $\nu\bar{\nu} \to W^+W^- \to e^-\bar{\nu}_e$:

$$J_{We} \cong 3.3 \times 10^{10} n_\nu^2 (\hat{m}_\nu^2 + 4)\hat{m}_\nu^{-1}(\Delta E/\text{GeV}), \quad \hat{m}_\nu^2 > 2 . \tag{32}$$

The result of the numerical calculations of the flux J_{We} is also presented in Fig. 2 at $\Delta E = 1$ GeV. As one can see from Fig. 2, the flux (32) is lower than the experimental flux (28) as a result of the fast decreasing of $J_{We} \sim m_\nu^{-3}(\hat{m}_\nu^2 + 4)^{-1}$, and therefore in this case there is no additional constraint on m_ν.

The annihilation reaction $\nu\bar{\nu} \to W^+W^-$ with positron production from W-boson decays $W^+ \to e^+\nu_e$ could also give an important constraint on m_ν. The flux of positrons in this case would be the same as for electron production (32). However, the energy resolution of the positron detection is very poor (for example, at the positron energy $E_e \cong 30$ GeV the resolution is $\Delta E \cong 20$ GeV [12]), but the experimental flux J^{expt}, as well as the flux J_{We} (in contrast with J_e for the reaction $\nu\bar{\nu} \to e^+e^-$), is proportional to the energy resolution and its relative value does not depend on ΔE (at $\Delta E \ll Ee$). If we assume, for instance, the validity of the model of dynamical halo, then it follows from Fig. 2 that in this case the limit on the heavy neutrino mass would be $m_\nu > 500$ GeV.

V. ROLE OF HIGGS MESON IN THE COSMOLOGICAL CONSTRAINTS ON THE NEUTRINO MASS

It was assumed in previous sections that the Higgs meson is heavy ($m_H \gg m_\nu$) so that the contribution of Higgs meson exchange,

$$\nu\bar{\nu} \to H \to W^+W^- , \tag{33}$$

to the process of neutrino annihilation can be neglected.

In this section we will take into account the finite mass of the Higgs meson and show how reaction (33) modifies the restriction (29), which was obtained by us in the limit case $m_H \to \infty$. In the case of a finite mass of the Higgs meson, the cross section for the process (33) is given by

$$\sigma_H = \sigma_{HH} + \sigma_{LH} = \frac{G^2 M_W^6}{16\pi} \frac{\beta_W}{\beta} \frac{m_\nu^2}{s} D_H [G_{HH} + 2(\hat{s} - \hat{m}_H^2) G_{LH}] , \tag{34}$$

$$G_{HH} = (\hat{s} - 4\hat{m}_\nu^2)(\hat{s}^2 - 4\hat{s} + 12) , \tag{35}$$

$$\begin{aligned} G_{LH} = {} & -\hat{M}_L^2 2(\hat{s}+2) + 2(\hat{s}+2)\hat{m}_\nu^2 - \hat{s}^2 + 2\hat{s} - 8 + \frac{2}{\hat{s}\beta\beta_W}\{-\hat{M}_L^4(\hat{s}+2) + \hat{M}_L^2[2(\hat{s}+2)\hat{m}_\nu^2 - \hat{s}^2 + \hat{s} - 2] \\ & -\hat{m}_\nu^4(\hat{s}+2) + \hat{m}_\nu^2(3\hat{s}-2) - 2\hat{s} + 4\} \ln\left[\frac{\hat{s}(\beta_W\beta - 1) + 2(1 + \hat{m}_\nu^2 - \hat{M}_L^2)}{-\hat{s}(\beta_W\beta + 1) + 2(1 + \hat{m}_\nu^2 - \hat{M}_L^2)}\right] , \end{aligned} \tag{36}$$

where $\beta_W = (1 - 4M_W^2/s)^{1/2}$, $D_H = [(s - m_H^2)^2 + \Gamma_H^2 m_H^2]^{-1}$, Γ_H is the width of the Higgs meson, $\hat{s} = s/M_W^2$, and $\hat{M}_L = M_L/M_W$. The second term (σ_{LH}) in Eq. (34) is the result of the interference between the diagrams of the s-channel H boson exchange and the t-channel charged L-lepton exchange. The interference between the diagrams of Z and H boson exchanges in the total cross section does not occur. The cross section σ_H was calculated for the first time in Ref. [4], but we note that there is a difference between the expression for G_{LH} in formula (36) and the corresponding expression of Ref. [4].

As the temperature in the Universe drops, very heavy neutrinos quickly become nonrelativistic and we can put $\hat{s} \cong 4\hat{m}_\nu^2 + \hat{\Delta}$, $\hat{\Delta} \ll 4\hat{m}_\nu^2$. Moreover, we assume that the masses of neutrino and L lepton are large (as numerical calculations show, the residual concentration of very heavy relic neutrinos and the output of cosmic rays from neutrino annihilation depend weakly on M_L), and in addition $M_L^2 \gg m_\nu^2 \gg M_Z^2$. Then

$$\beta\sigma_H \cong \frac{G^2 M_W^2}{8\pi} D_H' \hat{m}_\nu^4 \hat{\Delta}(6 - m_H^2/m_\nu^2) , \tag{37}$$

where $D_H' = D_H M_W^4$. We note that this expression can be negative as well as positive since σ_H is only a part of the total cross section $\nu\bar{\nu} \to W^+W^-$.

We must compare expression (37) with the formula for the cross section of process (5) in the approximation under consideration,

$$\beta\sigma_W \cong \frac{G^2 M_W^2}{16\pi}\hat{m}_\nu^2 , \quad (38)$$

and thus

$$\sigma_H/\sigma_W \cong \frac{2m_\nu^2 M_W^2 \hat{\Delta}}{(s-m_H^2)^2+\Gamma_H^2 m_H^2}(6-m_H^2/m_\nu^2) . \quad (39)$$

We see that at $m_H \gg m_\nu$ or $m_H \ll m_\nu$ the ratio σ_H/σ_W is small and only in the region of resonance $m_H \cong 2m_\nu$ can the ratio $\sigma_H/\sigma_W \cong M_W^2\hat{\Delta}/\Gamma_H^2$ be large. Indeed, calculating the residual concentration of the relic neutrinos, we evaluate the annihilation cross section at $s = \langle s_f\rangle$ [according to expression (10) this corresponds to $\hat{\Delta} = 6m_\nu T_f/M_W^2$] and at the width value $\Gamma_H \cong 61$ $(m_H/500\text{ GeV})^3$ GeV in the case $m_H \gg M_Z$ [14] so that $\sigma_H/\sigma_W \cong 0.4m_\nu^2/\Gamma_H^2 > 1$ at least for 250 GeV $< m_H <$ 500 GeV.

In the region of neutrino masses below the mass M_W, the dominant annihilation channels of heavy neutrinos are reaction (8) and

$$\nu\bar{\nu} \to H \to f\bar{f} , \quad (40)$$

where f denotes a fermion. The formula for the cross section of process (40) is given in Ref. [4], and using also the result (9), we obtain (at $m_\nu < M_W$)

$$\sigma_H(f\bar{f})/\sigma_W(f\bar{f}) \cong 3\frac{\hat{\Delta}m_\nu^2 m_f^2}{\hat{s}M_W^4}\frac{D_H}{D_Z} . \quad (41)$$

Far from H resonance, expression (41) is small. However, at H resonance ($m_H \cong 2m_\nu$) and far from Z resonance, this expression takes the form $\sigma_H/\sigma_W \sim 0.1(m_b/\Gamma_H)^2 \gg 1$ as a result of the smallness of the width of the Higgs meson $\Gamma_H \cong 0.45$ $(m_H/10\text{ GeV})$ MeV in this region (here we put the value m_f equal to the mass of the b quark, which is the heaviest fermion in this region since the t-quark mass $m_t > M_W$).

Thus, if the neutrino mass is around the resonance one for Higgs meson exchange, the annihilation cross section of very heavy neutrinos sharply increases, and this leads to a sharp decreasing in the residual (relic) concentration of such neutrinos in the Universe. One the other side, in the calculations of the annihilation cross section of very heavy neutrinos in the halo of the Galaxy, the effect of the Higgs meson is not important. This is due to the fact that such neutrinos in the Galaxy are strongly nonrelativistic [see Eq. (25)], and therefore the value $\hat{\Delta}$ in this case is very small ($\hat{\Delta} \cong 2\times 10^{-6}\hat{m}_\nu^2$). Also an additional suppression factor in Eq. (41) will appear, as a result of the smallness of the electron mass, when we consider the most interesting channel of the annihilation in the Galaxy, $\nu\bar{\nu} \to e^+e^-$, which produces a practically monochromatic line in the spectrum of cosmic electrons at energies $E_e \cong m_\nu$.

Results of numerical calculations of the electron flux from the very heavy neutrino annihilation in the galactic halo are presented in Fig. 4. In these calculations the residual concentration of relic neutrinos was evaluated by formula (4) and the flux of relativistic electrons was estimated according to expression (22). For the averaging of the cross section over the temperature, we used the formula of Ref. [15]:

$$\langle\sigma v\rangle = [8m_\nu^4 T K_2^2(m_\nu/T)]^{-1} \times \int_{4m_\nu^2}^{\infty} ds\, s^{1/2}(s-4m_\nu^2)K_1(s^{1/2}/T)\sigma(s) , \quad (42)$$

where $K_n(z)$ is the MacDonald function.

The experimental spectrum of cosmic electrons integrated over the energy resolution of the detector, $\Delta E \cong$ 1 GeV, is also shown in Fig. 4. As follows from Fig. 4, Higgs meson exchange modifies significantly in the resonance region the constraint (29) (which was obtained assuming $m_H \gg m_\nu$) to the form 60 GeV $< m_\nu <$ 90 GeV at $m_H = 200$ GeV.

In this section we have shown that Higgs meson exchange affects significantly the constraint on the mass of very heavy neutrinos in the case when the neutrino mass is around the resonance value $m_H \cong 2m_\nu$. The numerical calculations and comparison with the experimental data show that, in this case, the constraint on the neutrino mass is absent because of the smallness of the residual (relic) concentration of neutrinos in the neutrino mass range $\Delta m_\nu \sim 10\Gamma_H$. If, below the threshold of W-boson pair production, this value is not large ($\Delta m_\nu < 0.1$ GeV), then above the threshold the width of the Higgs meson is large and the constraint on the neutrino mass is eliminated in the region Δm_ν of the order of 10 GeV.

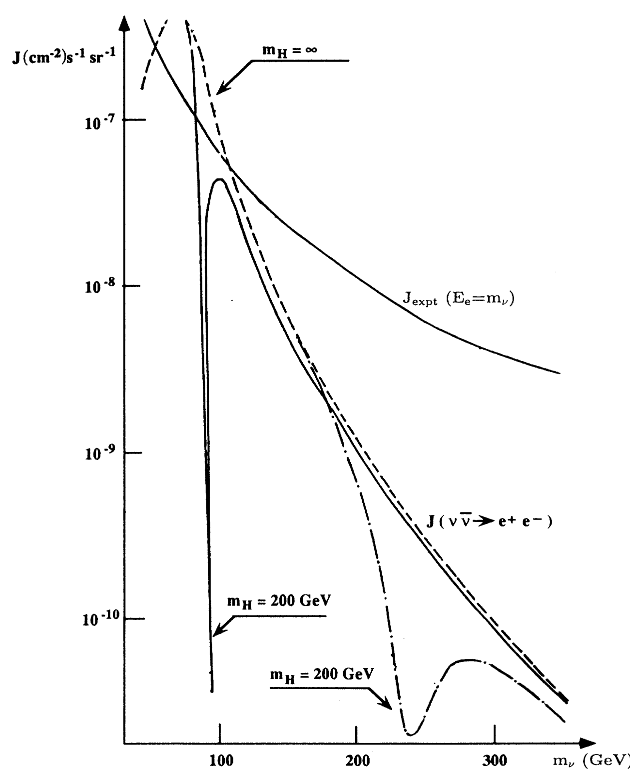

FIG. 4. Flux of cosmic electrons from the neutrino annihilation in the Galaxy as function of the neutrino mass for three values of the Higgs meson mass.

VI. CONCLUSION

In this paper, using the idea of Ref. [6] about neutrino condensation in the gravitational field of collapsing matter at the stage of Galaxy formation and analyzing the processes of cosmic ray production due to the relic neutrino annihilation in the Galaxy, we excluded the possibility of the existence of heavy stable neutrinos in the mass range 60–115 GeV. These bounds on neutrino masses imply also a corresponding bound on the total energy density of the Universe. We have shown that in the allowed neutrino mass range 44 GeV $< m_\nu <$ 60 GeV the neutrino energy density is less than the luminous baryonic density and only in the mass range 115 GeV $< m_\nu <$ 300 GeV can the relic neutrino density be large enough to be of observational interest (comparable to that of baryons), i.e., $3 \times 10^{-3} < \rho_\nu/\rho_c < 10^{-2}$, where $\rho_c = 10^{-29}$ g/cm^3. These constraints could be considerably extended by improving the precision of measurements of electron and positron spectra and their energy resolution, and by separate measurements of electron and positron spectra at high energies. It seems that the study of photon production in the Galaxy could reduce the relic neutrino density to a negligible value even in comparison with the luminous baryonic one, thus leading at least to a severe restriction on the role of such cold dark matter particles in solving the dark matter puzzle in the galactic halo.

Our treatment can be easily extended to any other weak interacting stable (neutral) particles, and, what is more, the weaker is their interaction the larger is their residual relic concentration in the Universe and the more stringent constraints could be obtained on the parameters of these particles from astrophysical experiments. It is also applicable to the analysis of the expected distribution and effects in the Galaxy of hypothetical strongly interacting massive particles (SIMP's), assumed to form collisional, but nonradiating gas [16], which may also condense via the considered mechanism [6]. It is also important to emphasize that our approach is valid in the case of a multicomponent dark matter [9] and may be used for species that give even a negligible contribution to the total energy density of the Universe.

The detection of an anomalous output of positrons with energy above 50 GeV would be a clear signature of the annihilation of Dirac neutrinos in the galaxy halo because the annihilation of Majorana fermions in the Galaxy is severely suppressed. This could be strong confirmation of a multicomponent dark matter content in the galaxy halo.

The search for heavy neutrinos at accelerators in the reaction $e^+e^- \to \nu\nu\gamma$ could give a possibility of analyzing the mass region $m_\nu \sim M_Z/2$, which is difficult for an astrophysical investigation. This problem will be discussed elsewhere.

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