Nonequilibrium shells of neutron stars and their role in sustaining x-ray emission and nucleosynthesis

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A review is presented of theoretical concepts of properties of matter and of processes in neutron star shells and of their relation to observations. The formation of a hot neutron star and its subsequent cooling leads to the appearance of a nonequilibrium layer in which energy of up to 10^{48} ergs is stored. The lack of equilibrium consists of the presence of superheavy nuclei with a large neutron excess near the limit of neutron evaporation $Q_n = 0$ and of free neutrons. The slow diffusion of neutrons into the interior of the star maintains the x-ray luminosity $L_X \gtrsim 10^{34}$ ergs/sec during $t \approx 10^4$ years. Nonstationary processes are discussed which are associated with the existence of the nonequilibrium layer. Outward transport of nonequilibrium matter and nuclear explosions may be associated with the observed x-ray bursts. Breaking up of the crust accompanying the diffusion of neutrons into the interior of the star and the expansion of the shell may explain abrupt changes in the period of the Crab pulsar. Expulsion of neutrons and superheavy nuclei into interstellar space accompanying explosions and subsequent electrodynamic acceleration of particles may explain the origin of heavy elements with $A > 150$ and also of deuterium.

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CONTENTS

1. INTRODUCTION

Neutron stars were discovered in 1968 as sources of pulsed radio emission (pulsars).[1] This occurred more than 30 years after the existence of such stars had been predicted by astronomers who had studied supernovae[2] and theoreticians concerned with the internal structure of stars.[3,4]

Interest in the physics of neutron stars increased rapidly after this outstanding discovery. This discovery of x-ray pulsars[5] confirmed the notion that x-ray sources may be associated with neutron stars and generated a new and still continuing wave of activity in neutron-star-research. One of the basic problems, and one to which many papers have been devoted, is that of explaining the luminosity of neutron stars. The accretion model, which can be applied to neutron stars in binary systems,[6] has been developed farthest toward explaining their luminosity in the x-ray region.[1)] The radio emission of pulsars is explained by processes in the magnetospheres of rotating stars.[7]

Another group of papers has been devoted to the internal structure of neutron stars and determination of their mass limit.

In the present paper we shall consider problems with a bearing on the properties of matter in the crust of a neutron star at densities of $10^{10} < \rho < 10^{12}$ g/cm^3 and show that physical phenomena that are important for

[1)]The conclusion that neutron stars must be x-ray sources was drawn in Ref. 141; accretion on neutron stars in binary systems was first considered in Refs. 142 and 143.

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various astrophysical applications unfold there. We shall also touch briefly upon problems of the internal structure of an equilibrium neutron star, its upper mass limit, and the formation and observational manifestations of neutron stars.

During cooling, a nonequilibrium layer in which a stored energy of up to 10^{48} erg accumulates is formed in the crust of a neutron star.[33,34] The lack of equilibrium consists of a large excess of neutrons, and this leads to the appearance of superheavy nuclei near the stability limit $Q_n=0$ and to the presence of free neutrons. The slow evolution of the nonequilibrium layer by diffusion of neutrons into the interior of the star[67] is responsible for the long persistence of the x-ray luminosity, $\sim 10^4$ yr. Nonstationary processes brought about in the neutron star by "starquakes" and abrupt changes in period[77-82] result in outward expulsion of nonequilibrium matter and explosive energy release, and may explain the observed γ-ray bursts.[91-96] The interstellar medium is enriched in heavy elements as the nonequilibrium matter escapes, and in deuterium as a result of combination of the escaping neutrons with surrounding protons.[67]

2. THE THEORETICAL MASS LIMIT OF NEUTRON STARS,OBSERVATIONAL DATA AND MECHANISMS OF FORMATION

Since the bulk of matter in neutron stars is in its lowest energy state, study of the physical properties of neutron-star matter reduces to calculation of the thermodynamic functions of matter at atomic-nucleus densities $\rho_n \sim 2\times 10^{14}$ g/cm^3 and higher. The methods and results of study of the properties of nuclear matter within neutron stars and calculations of neutron-star structure on the basis of the resulting equation of state are described in the reviews of Refs. 8–11, 167. The theory of nuclear matter, which is the basis of derivation of the equation of state and calculation of other thermodynamic functions, is still far from perfection. The masses of neutron stars have an upper limit $M_{\rm lim}$, which depends in the final analysis on the equation of state $P(\rho)$. The most complete equation of state calculations based on the theory of Bethe-Bruckner-Goldstone[12] and Pandaripande[13] give values of no more than $2M_\odot$ for the upper mass limit of neutron stars.[10] Recently published papers give higher mass limits for neutron stars. Their authors, appealing to the imperfection of methods for calculating the properties of nuclear matter, find various upper limits for $M_{\rm lim}$. The hardest equation of state,[14] which retains the causality principle $v_s \leqslant c$ (v_s is the velocity of sound), was used in Ref. 15. Here it was assumed that all calculations are unreliable at $\rho > \rho_* = 5.09\times 10^{14}$ g/cm^3,[2)] and the hardest equation of state $p=\varepsilon$ was used. Under these assumptions, $M_{\rm lim}$ was found equal to be $3.2M_\odot$. The way to increase the neutron-star mass limit is to abandon the causality principle and consider incompressible matter. Abandonment of causality at densities $\rho > \rho_*$ gives[16] $M_{\rm lim}=5M_\odot$, and if it is assumed that causality is violated at $\rho > 2\times 10^{14}$ g/cm^3, which corresponds to nuclear density, $M_{\rm lim}$ reaches $8M_\odot$.[17] The value of $3.2M_\odot$ can apparently be regarded as the most realistic upper limit for $M_{\rm lim}$. A calculation based on relativistic many-body theory results in an equation of state that puts the mass maximum at $2.39M_\odot$.[18] At $M > M_{\rm lim}$, a cold star should become a relativistic object—a black hole. In the x-ray source Cyg X-1, the mass of the compact object is no less than $6M_\odot$, and for this reason a black hole is assumed to exist in this source.[19] Radiopulsars, which now number 147,[20] and x-ray pulsars, of which, counting the recently discovered long-period pulsars, 14 have already been identified,[21,157] are, in all probability, neutron stars.[3)]

The discovery of x-ray sources in binary systems[6] and that of a radiopulsar in a binary system[158] have made it possible to estimate the mass of a neutron star from observations. A detailed investigation of the masses of observed neutron stars in binary systems[148] has shown that they may all lie in the range $(1.2-1.8)M_\odot$, although the upper limit may be higher for each star taken alone and reach $4M_\odot$ for SMC X-1, while the lower limit is $0.6M_\odot$ for Cen X-3. These results are in good agreement with theoretical estimates of the upper mass limit: $M_{\rm lim}=(2-3)M_\odot$.

According to present-day theoretical concepts, neutron stars are formed in nonstationary processes associated with supernovae. The attempt to explain the observed energies of supernovae was one of the principal factors in prediction of the existence of neutron stars.[2] Investigation of the last stages in the evolution of stars inevitably leads to the conclusion that a hydrodynamic-compression stage leading to the formation of a neutron star is necessary.[22]

One possible picture of the formation of neutron stars and black holes can be sketched briefly on the basis of today's concepts of the evolution of "ordinary" stars. A degenerate carbon core forms in stars with masses $3M_\odot < M < 10M_\odot$ at the red giant stage as a result of the reaction $3\mathrm{He}^4 \to \mathrm{C}^{12}$. Enlargement of the core proceeds smoothly until its mass approaches the Chandrasekhar limit. For a molecular weight $\mu_e = A/Z = 2$, which corresponds to the chemical composition of carbon, the Chandrasekhar limit is $M_L = (5.75M_\odot/\mu_e^2)\sim 1.44M_\odot$. Density and temperature increases inexorably in this core (the pressure of the degenerate electrons is insufficient), and explosive combustion of the carbon develops, with the resulting ejection of outer layers and formation of a neutron star with $M_{ns}=1.4M_\odot$ or, if combustion develops at lower densities, total explosion with supernovae effects.[39,145,166] In the evolution of massive stars ($M > 10M_\odot$), the temperatures at which carbon is formed and burns are reached at lower densities (the electrons are nondegenerate): quiet combustion continues until the star has formed a core of

[2)]An equation of state for $\rho < 5.09\cdot 10^{14}$ g/cm^3 was calculated in Ref. 25 on the basis of the Bethe-Bruckner-Goldstone theory. According to Ref. 25, there are not reliable calculations for higher densities.

[3)]As of the end of 1977.

iron-peak elements. During subsequent contraction of the star in the process of energy loss by neutrino emission and neutronization of its matter, it loses hydrodynamic stability. The cause of hydrodynamic-stability loss is a decrease in the adiabatic exponent to $\gamma < 4/3$ as a result of neutronization and breaking up of nuclei of the iron-peak.[46,47,51,162] Models of star cores at the stability limit have been developed.[144] An extended evolution of stars with formation of "iron" cores with masses of $(1.2-1.6)M_\odot$ has been developed to the point of stability loss.[160,161] Collapse of these cores results in the formation of hot neutron stars surrounded by dense shells. The mass of the neutron star is equal to that of the separated core. Matter is not ejected[36-38]; the subsequent fate of the core and the shell is not discussed. Therefore the question as to whether neutron stars or black holes are formed in this model cannot be resolved without analysis of the fate of the massive ($M_{sh} \gtrsim 10M_\odot$) shell. Subsequent ejection of the shell and an observable spectacle of a supernova in combination with a neutron star in the form of a compact residue can apparently occur as a result of operation of a magnetorotation mechanism[146,147,163] in which the energy source is rotation and the energy of the rotation is converted into energy of ejection by a magnetic field. The fact that at least 50% of all stars are components of binary systems makes calculations of neutron-star and black-hole formation even more uncertain; phenomena of gas transfer from one star to the other greatly complicate evolution as compared to single stars. The discovery of neutron-star pulsars provided observational confirmation that supernovae are linked to the birth of neutron stars. Two of the youngest pulsars, PSR 0531 +21 and PSR 0833 − 45, are found in the Crab and Vela supernova remnants. Nebulosity associated with supernova remnants has also been detected around the pulsars PSR 0611 +22, 2021 +51, and 1154 − 62.[21] Thus, there is both theoretical and observational evidence favoring the formation of neutron stars as a result of a nonstationary process leading to a supernova. The ages of other neutron stars and the parameters of the interstellar medium around them are apparently such that the supernova remnants have had time to expand considerably and therefore become invisible.

Below we shall discuss the chemical composition of neutron-star shells during the cooling process, i.e., on the basis of an evolutionary approach. For comparison, we shall first examine the equilibrium neutron-star shells that are usually discussed.

3. EQUILIBRIUM NEUTRON-STAR SHELLS

For neutron-star shells at densities no greater than nuclear ($\rho < 2\times10^{14}$ g/cm³), the equation of state at full thermodynamic equilibrium can be calculated with high reliability, although even here the results obtained by different authors differ sharply. This is due basically to the differences in the computing methods and especially to the different treatments given the surface energy of the core. In the early papers,[24,25] the equilibrium composition of the core was characterized by a sharp increase of A and Z with increasing density. The expression for the surface energy was later improved,[26-28]

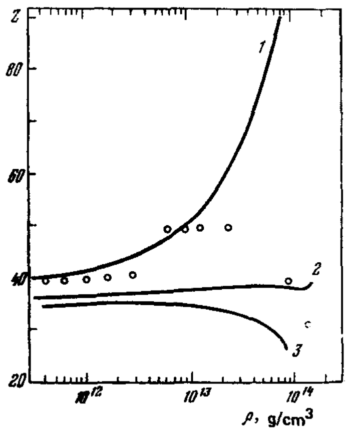

FIG. 1. Dependence of nuclear charge z in crust of a neutron star on density as constructed from calculations of various authors. 1) curve from Ref. 25; 2) from Ref. 28; 3) from Ref. 27; the circles represent the results of Ref. 29.

and this limited the growth of Z with increasing ρ at the level $Z=20-50$. $Z(\rho)$ curves from Refs. 25 and 27–29 appear in Fig. 1. Thus the nuclear composition of the equilibrium neutron-star shell has the following character (see, for example, Ref. 10). At densities $\rho < 8.1\times10^6$ g/cm³, matter consisting of ^{56}Fe—the element with the highest binding energy per nucleon B_N, has the lowest energy.[4)] As the density of matter increases, the Fermi energy of the electrons $\varepsilon_{f,e}$ increases to the point at which the appearance of heavier nuclei that have smaller B_N but an excess numbers of neutrons is energetically favored. The equilibrium shifts in the direction of larger relative numbers of neutrons and fewer electrons and protons. First to appear is ^{62}Ni, then ^{58}Fe and so forth all the way to $\rho = 4.3\times10^{11}$ g/cm³ (see also Ref. 154), when matter consists of ^{118}Kr. Then the electron Fermi energy $\varepsilon_{f,e} = 26.2$ MeV, and the ratio of the number of protons to the atomic weight A of the nucleus is $Z/A = 36/118 \approx 0.3051$. Free neutrons appear on a further increase in density under equilibrium conditions. We note that this change in composition corresponds to thermodynamic equilibrium states. In reality, the composition will be nonequilibrium as the cold matter contracts. At first, the change will take place with A constant and Z decreasing, until the stability limit $Q_n = 0$ is reached; at that point, free neutrons will appear and the nuclei will move along this boundary to $Z\approx 6$ at $\rho\approx 3\times10^{12}$ g/cm³, when heating becomes inevitable and, apparently, equilibrium is established.[34]

Below the threshold at which neutrons appear, the

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4)The isotope ^{62}Ni has the largest mass defect per nucleon, but the difference between Q_N for ^{56}Fe and ^{62}Ni is small[48]:

$$Q_N = \frac{Zm_p + (A-Z)\,m_n - m_{A,Z}}{A}.$$

$Q_N(^{62}\mathrm{Ni}) = 8.7950$ MeV, $Q_N(^{56}\mathrm{Fe}) = 8.7907$ MeV, and $\Delta Q_N/Q(^{56}\mathrm{Fe}) = 4.9\cdot10^{-4}$. It is assumed that the binding energy of the nucleus is reckoned from a fixed state of matter, for example the protonic state, so that the nuclear binding energy per nucleon is $B_N \neq Q_N c^2$. In particular, $Q_N(^{62}\mathrm{Ni}) > Q_N(^{56}\mathrm{Fe})$ while $B_N(^{62}\mathrm{Ni}) < B_N(^{56}\mathrm{Fe})$.

pressure is determined chiefly by the degenerate relativistic electrons, for which we have[30]

$$P_e = \frac{3^{1/3}\pi^{2/3}}{4}\frac{\hbar c}{m_p^{4/3}}\left(\frac{\rho}{\mu_e}\right)^{4/3} \approx 1.2\cdot10^{15}\left(\frac{\rho}{\mu_e}\right)^{4/3}\ (\text{dyn/cm}^2); \quad (3.1)$$

μ_e is the number of nucleons per electron.

The pressure of the nuclei is negligibly small at all densities in a cold neutron star, and the contribution made by the pressure of the degenerate nonrelativistic neutrons

$$P_n = \frac{(9\pi)^{1/3}}{5}\frac{\hbar^2}{m_p^{8/3}}\rho_n^{5/3} \approx 5.3\cdot10^9\rho_n^{5/3}, \quad m_n \approx m_p, \quad (3.2)$$

increases progressively with increasing density. It reaches 20% at $\rho = 1.5\times10^{12}$ g/cm^3 and as much as 80% for the minimal-energy state[5)] at $\rho = 1.5\times10^{13}$ g/cm^3.

The interaction between neutrons becomes significant as the density increases. It influences both the chemical composition of matter and the equation of state. As we have noted, calculations of the thermodynamic properties of matter in a state in which nuclei, electrons, and neutrons are present (the Aen phase) are highly sensitive to the description of core surface energy. As the density increases to $\rho = \rho_0 \sim 10^{14}$ g/cm^3, at which there is a transition from the Aen phase to homogeneous nuclear matter, the surface energy E_S vanishes, and its behavior near ρ_0 is highly important. An improvement in Refs. 26–29 as compared to Ref. 25 consists of the fact that E_S tends to zero more rapidly in a more exact calculation than had been assumed empirically in Ref. 25. The interactions of free neutrons and nuclei in the core must be described consistently to obtain noncontradictory results, and this was done in Refs. 26–29, in contrast to Ref. 25, where the empirical formula of Myers and Swiatecki[31] was used for the nuclei. Calculations of the interaction energy of nucleons in the theory of nuclear matter by the method of hyperspherical functions are compared in Ref. 32 with the result from the empirical Weizsäcker formula. It is shown that the Weizsäcker formula can be used for rough estimates of the binding energy because the use of various nuclear potentials results in scatter comparable to the difference from the Weizsäcker formula. The equation of state of matter at subnuclear and nuclear densities and calculations of neutron-star models are presented in greater detail in the review.[167]

With this we conclude our review of results of calculations of the equilibrium state of matter in neutron stars and proceed to a more detailed evolutionary treatment of the properties of matter at densities $10^9 < \rho < 10^{12}$ g/cm^3, following Refs. 33 and 34.

4. PHYSICAL PROCESSES IN THE SHELLS OF HOT NEUTRON STARS

Above we briefly set forth a scenario of neutron-star formation in accordance with present-day concepts. One of the important stages in its evolution is that of the hot neutron star.[35] In present calculations of the dynamics of the collapse,[36–39] temperatures of the order of 10^{10}–10^{11} °K are attained in matter when it reaches neutron-star densities. Under these conditions, various reactions that due to high Coulomb barriers proceed extremely slowly in matter at ordinary temperatures may now proceed rapidly. Let us explain this in greater detail.

The Coulomb barrier (expressed in megaelectron-volts) is

$$B = \frac{Z_1 Z_2}{A^{1/3}}\ (\text{MeV}), \quad (4.1)$$

where A is the baryon number and Z is the charge of the nucleus. $B \approx 7$ MeV for the interaction of protons with iron-group nuclei. The average thermal energy of the protons is

$$E_T = \frac{3}{2}kT = 0.133T_9\ (\text{MeV}) \quad \left(T_9 = \frac{T}{10^9\ ^\circ\text{K}}\right). \quad (4.2)$$

As a rule, the temperatures do not exceed 10^9 °K in stationary stars, and $kT \ll B/10$; under stellar conditions, therefore, the Coulomb barrier completely determines the rates of nuclear reactions with charged particles.[6)] If the temperatures are high enough so that the rates of nuclear reactions are high compared to the rates of other processes that change conditions in matter (in our case density and temperature), it may be assumed that all nuclear-reaction channels are open. In this case, matter is under detailed equilibrium conditions with respect to nuclear reactions. For characteristic processes in stars, the conditions of detailed equilibrium for nuclear reactions are reached when the temperature exceeds $T > (3-5)\cdot 10^9$ °K. Nuclear reactions will not change the ratio of the total number of neutrons N_n (in free and bound states) to the total number of protons N_p.

We may therefore introduce the parameter $R_N = N_n/N_p$ and consider the equilibrium at a given R_N.[40] In real physical situations, however, R_N varies as a result of β processes and is, generally speaking, time-dependent. Under certain conditions, it is possible to eliminate the time dependence, for example, under conditions of full thermodynamic equilibrium, when detailed equilibrium is realized for β processes. Situations in which it is difficult to satisfy conditions for complete thermodynamic equilibrium are most often encountered in nature. This is because of the enormous differences in the times of the processes (for example, $t_{\text{nuc}} \ll t_\beta$ at $T \gtrsim 5\times10^9$ °K for nuclear and weak interactions) or because of the absence of equilibrium for some type of particle (neutrino) involved in the processes. Under these conditions, we may separate fast and slow processes and consider the detailed equilibrium for fast processes and the kinetic equations for slow processes.

Such a situation occurs in the collapse stage of the star at densities $\rho < 10^{12}$ g/cm^3 and temperatures $5\cdot10^9$

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5) The electrostatic and magnetic interactions of electrons and nuclei lower the pressure P in comparison with the pressure P_e of the ideal Fermi electron gas. In the case of ultrarelativistic electrons, the correction depends only on the nuclear charge Z (Refs. 149, 22): $P/P_e = 1.00116 - 4.56\cdot10^{-3}Z^{2/3} - 1.78\cdot10^{-5}Z^{4/3}$.

6) We note that all nuclei are assumed to have Maxwellian distributions; then the reaction rate and the effective energies of the particles participating in the reaction are determined by the method of steepest descent and $E_{\text{eff}} \approx 10kT$.

$< T < 10^{10}$°K. Matter is transparent for the neutrino, and there is no equilibrium for it, while detailed equilibrium exists for nuclear reactions. An important particular case is kinetic β-process equilibrium.[42-45]

Let us explain the concept of kinetic β-process equilibrium in greater detail. When there are several reaction channels leading to both increases and decreases in R_N, a situation may arise in which the resultant rate of decrease is equal to the resultant rate of increase of R_N, i.e., R_N remains constant. Then the time of the temperature decrease due to free escape of ν, $\bar{\nu}$ may prove to be quite large.[46]

We stress that detailed equilibrium does not come about in each of the channels. Such equilibrium is known as kinetic equilibrium with respect to β-processes. All the above points up the great importance for the analysis of physical models of the relation between the rates of the processes or the ratio of the characteristic times for processes taking place in matter.

Let us examine the basic processes that characterize conditions in stars at the later stages of evolution, and in the process of hydrodynamic collapse or disintegration of the star's dense core. The basic thermodynamic characteristics of the state of matter under stellar conditions are density and temperature. Chemical composition also depends on the thermodynamic variables, which are determined in turn from stellar evolution and hydrodynamic calculations. If the star is not in equilibrium, the physical conditions in its matter vary with a hydrodynamic time that can be determined, according to Ref. 22, from the expression

$$t_H = \frac{1}{\sqrt{4\pi G\rho}} \approx \frac{10^3}{\sqrt{\rho}}\,(\mathrm{sec}), \tag{4.3}$$

where G is the gravitational constant and ρ is the density (in g/cm^3). When the star is in mechanical equilibrium, its evolution is determined by energy losses in photon or neutrino emission. At high temperatures, the basic energy losses occur in the form of neutrino emission. Then the characteristic time of evolution if matter is transparent to neutrinos is

$$t_\nu = \frac{E_Q}{q_{\nu\bar{\nu}}}, \tag{4.4}$$

where E_Q is the thermal energy of the star and $q_{\nu\bar{\nu}}$ is the energy flux in the neutrino emission. If the size of the star exceeds the neutrino mean free path for absorption, $l = 1/\sigma_\nu n_n$, the time for cooling due to neutrino emission increases by a factor $\sim R/l$ when $R/l \gg 1$, and the neutrino emission flux decreases by a factor $e^{\sigma_\nu n_n R}$ when $R/l < 1$.[35] The quantity $\tau_\nu = \sigma_\nu n_n R = (\sigma_\nu / m_p)\rho R x_n$ is naturally referred to as the optical thickness for neutrinos in the nucleon gas, while $\varkappa_\nu = \sigma_\nu x_n / m_p$ is the neutrino opacity and x_n is the fraction of neutrinos by weight. At $\tau_\nu \gg 1$, a neutrino in the interior of the star is in thermodynamic equilibrium with matter, and the neutrino luminosity is determined approximately (assuming $\mu_\nu = 0$) by the temperature T_ν and the size R_ν of the neutrino photosphere, on which $\tau_\nu = 1$ (Ref. 41):

$$L_\nu = \frac{7}{8}\,\sigma T_\nu^4 \cdot 4\pi R_\nu^2, \tag{4.5}$$

where σ is the Stefan-Boltzmann constant and the multiplier 7/8, by which relation (4.5) differs from the corresponding relation for photons, appears because the energy density of the equilibrium neutrinos and antineutrinos for $\mu_\nu = \mu_{\bar{\nu}} = 0$ equals[30]

$$E_{\nu\bar{\nu}} = \frac{7}{8} E_\gamma. \tag{4.6}$$

TABLE I. Values of $\lg t_\nu$ (sec) (t_ν is the time of neutrino energy loss in the URCA process as calculated from (4.4))

$\lg \rho/T_9$	5	8	10	12	16
9	7	4	2.3	1.3	0.3
10	12	6.1	3.9	2.8	0.9
11	—	8.5	5.5	4.6	2

Another important parameter for determination of the properties of matter at high temperatures is t_{nuc}—the characteristic time of establishment of detailed equilibrium for nuclear reactions. As a characteristic time we shall take the proton photodetachment time, since the rates of nuclear reactions involving neutrons and the rates of radiative capture of neutrons and protons are much higher than the proton photodetachment rate under conditions of equilibrium between matter and radiation. The proton photodetachment time has been determined for equilibrium conditions.[34,46]

The next important process that describes the state of the matter is the β process. As we noted above, it can be assumed in view of the small neutrino-antineutrino interaction cross section at densities $\rho < 10^{12}$ g/cm^3 and temperatures $T_9 < 100$ that only bulk energy losses due to neutrinos occur, i.e., neutrinos and antineutrinos escape freely from matter.[47] Therefore thermodynamic β-process equilibrium is not established under these conditions. The kinetics of the β processes must be taken into account.

The characteristic times of the porcesses discussed above are given in Tables I–III taken from Ref. 34. Comparing the characteristic times, we may conclude that the chemical composition at the stage of rapid hydrodynamic compression of hot matter can be regarded as being in equilibrium with respect to nuclear reactions and frozen with respect to β processes, since $t_{\mathrm{nuc}} < t_H < t_\beta$. During the existence of the hot neutron star, which is being cooled by neutrnio energy losses, the validity of $t_\nu \gg t_\beta$ means that the chemical composition can be calculated from the condition of β-process kinetic equilibrium where neutrinos escape freely from the matter. Calculations of chemical composition under kinetic equilibrium conditions indicate[44,45] that neutrons compose considerably more than half of the mass at

TABLE II. Values of $\lg t_\beta$ (sec) (t_β is the characteristic β-process time calculated in Ref. 34, and t_H was calculated from (4.3))

$\lg \rho/T_9$	5	8	10	12	16	$\lg t_H$
9	4.2	2.7	1.3	0.69	−0.6	−1.5
10	7	5.2	3.4	1.96	−0.2	−2
11	—	7.9	5.2	3.6	1.2	−2.5

TABLE III. Values of lg t_{nuc} (sec) as a function of Q_{p6}—the proton detachment energy in MeV (t_{nuc} is the nuclear reaction time calculated in Ref. 34)

T_9/Q_{p6}	10	16	22	28
3	3.185	13.205	—	—
5	−4.428	0.72	5.96	11.28
10	−12.7	−9.2	−5.7	—

densities $\rho > 10^{10}$ g/cm³ and temperatures $T_9 \sim 5$.

The above leads us to the following physical model of the evolution of the neutron-star shell. As a result of loss of mechanical stability, the star is compressed from the presupernova state with $\rho_s \lesssim 10^9$ g/cm³ (Ref. 144) to high densities of the order of or greater than nuclear. High temperatures are reached in matter during the compression process. The result is formation of a hot neutron star (NS). We shall consider only the layers of NS matter that are in the density range $10^{10} < \rho < 10^{12}$ g/cm³. At the moment of formation, as we indicated above, the condition of detailed equilibrium for nuclear reactions and kinetic equilibrium for β processes is satisfied in this layer of matter. The temperature drops rapidly as a result of neutrino-emission cooling. However, the nuclear equilibrium conditions remain satisfied as long as the temperature in matter remains above $T_9 > 5$. At this temperature, the time to establish detailed equilibrium with respect to nuclear reactions becomes larger than the neutrino-cooling time—the thermal time. We may therefore conclude that a phase of incomplete equilibrium begins at this moment and untimately results in formation of a nonequilibrium layer.

5. FORMATION OF THE NONEQUILIBRIUM LAYER AND ITS STRUCTURE

When the temperature drops below $T_9 < (4-5) \cdot 10^9$ °K, reactions with charged particles slow down. At densities $\rho > 10^7$ g/cm³, matter consists of nuclei (the concentration of which is determined by the prior history of matter at high temperatures), neutrons, and relativistically degenerate electrons. Under these conditions, reactions with neutrons, photodetachment and neutron capture, β decays at $\varepsilon_\beta > \varepsilon_{f,e}$, and e^- captures at $\varepsilon_\beta < \varepsilon_{f,e}$ are possible; here ε_β is the energy of β decay of the nucleus and $\varepsilon_{f,e}$ is the electron Fermi energy, which is determined as follows:

$$\varepsilon_{f,e} = m_e c^2 \left[\frac{\rho \text{ g/cm}^3}{\mu_e \cdot 10^6}\right]^{1/3} \text{(erg)}, \tag{5.1}$$

where $\mu_e = 1/\sum_i (Z_i/A_i) x_i$ is the number of nucleons per electron and x_i is the concentration by weight of the ith element. Let us estimate the characteristic time of neutron photodetachment as a function neutron binding energy in the nucleus. We have the following expression for the characteristic time $\tau_{\gamma n}$ of neutron photodetachment with averaging over the Planck photon distribution:

$$\frac{1}{\tau_{\gamma n}} = 6.95 \cdot 10^{33} T_9^{3/2} \langle \bar{\sigma}_{th} \rangle \exp\left(-\frac{11.605 Q_{nb}}{T_9}\right), \tag{5.2}$$

where Q_{nb} is the binding energy of the neutron in MeV and $\langle \bar{\sigma}_{th} \rangle$ is the averaged cross section of the reaction $A(n,\gamma)(A+1)$. Then, taking $\sigma_{th} \sim 1b$ and $T_9 < 0.04 Q_{nb}$, we obtain $\tau_{\gamma n} > 10^7$ yr, i.e., practically no neutrons with binding energies $Q_n > Q_{nb} = 10$ MeV are detached from the nucleus at $T_9 < 4$. Thus, equilibrium with respect to reactions with neutrons prevails in the system for nuclei with last-neutron binding energies below 10 MeV.

Let us explain the formation of nonequilibrium chemical composition with reference to the scheme in Fig. 2. The diagram represents the (A, Z) plane, on which the fission and stability limits of the nuclei have been plotted. Let us trace the track of variation of (A, Z) for a nucleus situated in a surrounding of neutrons that account for one-half or more of the weight of matter. The (A, Z) plane is divided into three regions:

Region I : $Q_n > Q_{nb}$,
Region II: $0 < Q_n < Q_{nb}$ and $\varepsilon_\beta > \varepsilon_{f,e}$,
Region III: $0 < Q_n < Q_{nb}$ and $\varepsilon_\beta < \varepsilon_{f,e}$.

At high neutron concentrations, nuclei corresponding to the above conditions will quickly capture neutrons and go from region I to region II or III irrespective of ε_β. This forms nuclei that have large excesses of neutrons and are situated far from the "stability valley" of the nuclei. Equilibrium with respect to neutron reactions prevails in region III; β processes, namely capture of electrons with high Fermi energies, always result in a decrease in Z because of the inequality $\varepsilon_\beta < \varepsilon_{f,e}$, with the result that the Z of all nuclei in this region decreases and they enter the region below level ab. In region II, on the other hand, the equilibrium for reactions with neutrons is accompanied by β^- decays with increasing Z until all nuclei have entered the region above level ed. We see that under limited-equilibrium conditions, the chemical composition of matter is determined by the rather narrow region in A and Z bounded by $abcd$. Nothing leaves this region because of the absence of admissible β processes and neutron photodetachment. At $Q_{nb} < 1$ MeV and $T_9 < 0.4$, there remains only one nucleus on the stability boundary, for

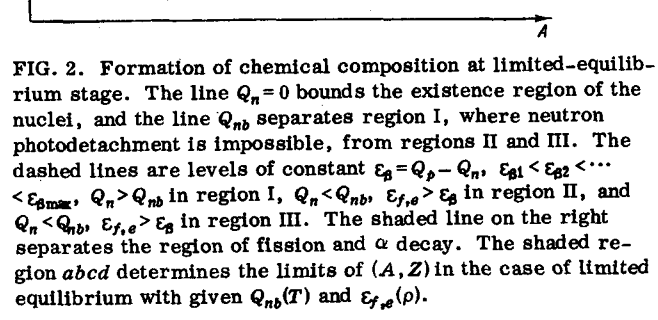

FIG. 2. Formation of chemical composition at limited-equilibrium stage. The line $Q_n = 0$ bounds the existence region of the nuclei, and the line Q_{nb} separates region I, where neutron photodetachment is impossible, from regions II and III. The dashed lines are levels of constant $\varepsilon_\beta = Q_p - Q_n$, $\varepsilon_{\beta 1} < \varepsilon_{\beta 2} < \cdots < \varepsilon_{\beta max}$, $Q_n > Q_{nb}$ in region I, $Q_n < Q_{nb}$, $\varepsilon_{f,e} > \varepsilon_\beta$ in region II, and $Q_n < Q_{nb}$, $\varepsilon_{f,e} > \varepsilon_\beta$ in region III. The shaded line on the right separates the region of fission and α decay. The shaded region $abcd$ determines the limits of (A, Z) in the case of limited equilibrium with given $Q_{nb}(T)$ and $\varepsilon_{f,e}(\rho)$.

which

$$\varepsilon_\beta = Q_p - Q_n \approx Q_p = \varepsilon_{f,e}. \tag{5.3}$$

Parasitic effects for the above process of formation of excess-neutron nuclei are nuclear fission, α decay, and reactions of the type $2(A, Z) \to A_1Z_1 + A_2Z_2 + A_3Z_3 + \ldots$, which lead to the formation of new embryonic nuclei capable of rapid absorption of excess nonequilibrium neutrons.

Estimation of the characteristic time τ_Z of the reaction of two nuclei with large Z indicates[34] that $\lg \tau_Z \gg 100$ when $\rho \sim 10^{10}$ g/cm^3, $T \sim 10^9$°K, and $Z = 26$. This estimate took account of the strong electron screening with pycnonuclear corrections in accordance with Ref. 49.

To estimate the influence of α decay and nuclear fission, it is necessary to make certain assumptions as to the properties of nuclei far from the "stability valley." We do not now have a rigorous theory and calculation of the parameters of such nuclei, let alone experimental material on heavy nuclei far out in the instability region. We can only make certain semiquantitative estimates. The uncertainty of these premises does not influence the qualitative aspect of the problem, although certain quantitative estimates may be affected. We shall adopt the condition that the stability boundary of the nuclei with respect to neutron evaporation occurs at $A = 4Z_0$ for $Z > 6$ and that the proton binding energy is then[7)]

$$Q_{pb} = 33 - \frac{Z}{7}\ (\text{MeV}). \tag{5.4}$$

For $Z < 6$, the value of Q_p drops off sharply at the limit of existence of nuclei. (A, Z) was calculated in Ref. 50 for the most abundant element under these conditions. The accuracy of formula (5.4) is sufficient for qualitative estimates of chemical composition. For cold matter, we have taking (5.3) and (5.4) into account

$$Z_0 = 7(33 - \varepsilon_{f,e}) = 7\left[33 - 0.511\left(\frac{\rho_6}{\mu_e}\right)^{1/3}\right], \quad A_0 = 4Z_0; \tag{5.5}$$

where $\varepsilon_{f,e}$ is in MeV and $\rho_6 = \rho/10^6$ g/cm^3.

The change in the stability-boundary parameters of the last nucleon does not affect the basic conclusions of this study concerning the formation of matter with excess neutrons and excess-neutron nuclei. The change in the quantitative results affecting the values obtained for the mass number A and the nuclear charge Z does not alter the qualitative picture of the processes unfolding in matter, but it does indicate the importance of studying the properties of nuclei with large neutron excesses.

It will be assumed from now on that the neutron concentration is high, so that even when all nuclei are on the stability boundary with $Q_n \approx 0$ matter still contains many free neutrons that can neither decay with $\varepsilon_{f,e} > 0.78$ MeV nor be captured by nuclei. At $\varepsilon_{f,e} > 33$ MeV, stable nuclei no longer exist in our approximation, and

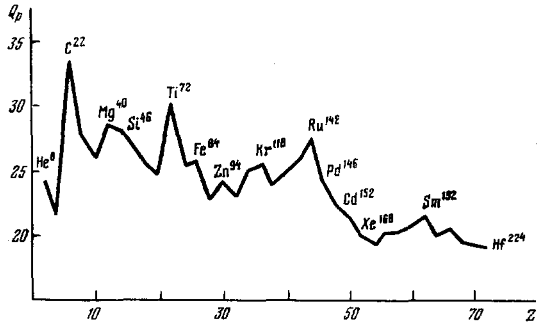

FIG. 3. Proton detachment energy Q_p vs. z for nuclei on the stability boundary with $Q_n = 0$, constructed according to P. É. Nemirovskiĭ's quantitative estimates.

electron capture begins to proceed under nonequilibrium conditions, with resulting rapid heating,[51-54] formation of new embryonic nuclei, and arrival of matter at a state close to thermodynamic equilibrium.[34] Figure 3 presents a plot of the proton detachment energy Q_p at the boundary $Q_n = 0$ according to P. É. Nemirovskii's calculations. We note that $\varepsilon_\beta = Q_p - Q_n$ and that $\varepsilon_\beta \approx Q_p$ for $Q_n = 0$ at the stability boundary. The inner boundary of the nonequilibrium layer occurs at the density given by

$$\rho_2 = \mu_e \cdot 10^5 \left(\frac{33}{0.511}\right)^3 \approx 2.7 \cdot 10^{11}\mu_e\ (\text{g/cm}^3), \qquad P_2 \approx 1.2 \cdot 10^{15}\left(\frac{\rho_2}{\mu_e}\right)^{4/3} \approx 2 \cdot 10^{30}\ \text{dyn/cm}^2. \tag{5.6}$$

Since $\mu_e > 4$, we have $\rho_2 > 1.1 \cdot 10^{12}$ g/cm^3.

We proceed directly to evaluation of the effect on the picture presented of α decay and spontaneous nuclear fission.

Estimation of the α-decay energy E_α with the Weizsäcker formula indicates that $E_\alpha < 0$ for nuclei with large Z and $A = 4Z$, and that decay does not occur. To estimate the time $T_{1/2}$ of spontaneous nuclear fission, we use the semiempirical formula

$$\lg T_{1/2} = 1.57 - 3.75\,\frac{Z^2}{N}\quad (T\ \text{in sec}). \tag{5.7}$$

At the stability boundary $A = 4Z$, we have $\lg T_{1/2} = 1.57 - 0.93Z$, i.e., $T_{1/2} > 3 \cdot 10^7$ yr for $Z < 153$. Induced fission is not considered because in the absence of an exact theory and experimental data, the correction would be within the error of the qualitative picture of the processes.

The upper boundary of the nonequilibrium layer at large Z is determined from the condition for no nuclear fission: Z increases with decreasing ρ, $Z_{max} = 153$, and this gives

$$\rho_1 \approx \mu_e \cdot 10^6 \left(\frac{11}{0.511}\right)^3 = 10^{10}\mu_e\ (\text{g/cm}^3), \qquad P_1 = 1.2 \cdot 10^{15}\left(\frac{\rho_1}{\mu_e}\right)^{4/3} = 2.5 \cdot 10^{28}\ \text{dyn/cm}^2 \tag{5.8}$$

with $\mu_e => 4$ and $\rho_1 > 4 \cdot 10^{10}$ g/cm^3.

Thus, a nonequilibrium layer with a large excess of energy stored in free neutrons forms in the shell of the neutron star. The mass of the layer can be estimated from a relation that follows from the thin-layer equilibrium condition[34] and application of (3.1) and (5.6):

[7)]Formula (5.4) interpolates approximately the quantitative estimates of P. É. Nemirovskiĭ, the results of which are shown in Fig. 3.

$$M=\frac{4\pi R^4}{GM}(P_2-P_1)\approx\frac{4\pi R^4}{GM}P_2\approx 0.1P_2\approx 2\cdot 10^{29}\ \mathrm{g}=10^{-4}M_\odot. \quad (5.9)$$

All the above indicates that the neutron star has a stored energy that makes manifestation of its activity possible even without external sources.

6. QUASISTATIONARY PROCESSES IN THE SHELLS OF NEUTRON STARS

a) Stored energy of a young neutron star

In the collapse process that preceeds the formation of a neutron star, its matter is heated by adiabatic contraction, nonequilibrium β-processes,[51-54] and by the possible detonation of nuclear fuel. The newly formed neutron star is hot as a result. Nor may we expect that it will be formed immediately in a state of static equilibrium. Oscillatory motions about the point of stable equilibrium are inevitable. These two forms of energy—thermal and oscillatory—have been assumed to be present in the young neutron star. The total stored energy at the instant of formation may be quite large. Models of hot neutron stars were constructed in Ref. 35. Hydrodynamic calculations[36-38] showed that the characteristic parameters of forming neutron stars agree well with statistical calculations.[35] Assuming, in accordance with Refs. 35–38, that the temperature of the hot neutron star averages $T_n\approx 10^{12}$ °K we obtain the second thermal energy

$$Q_q\approx\frac{3}{2}kT_n\frac{M}{m_n}\approx 2\cdot 10^{53}\ \mathrm{erg}. \quad (6.1)$$

To estimate the energy of the oscillations, we put the amplitude at the surface equal to 10% of the Keplerian velocity:

$$v_{0s}\approx 0.1\sqrt{\frac{GM}{R}}\approx 0.045c. \quad (6.2)$$

The corresponding oscillatory energy at $v_0=v_{0s}r/R$ equals

$$Q_0\approx\frac{4\pi}{2}\int_0^R v_0^2(r)\,\rho r^2\,dr=\frac{2\pi\cdot 0.002c^2}{100R^2}\int_0^R\rho r^4\,dr\approx 10^{51}\ \mathrm{erg}. \quad (6.3)$$

Here we have used the value $I_c=4\pi\int_0^R\rho r^4 dr\approx 10^{45}$ cgs units for the central moment of inertia of the neutron star in accordance with Refs. 55 and 56, with $M\approx M_\odot$ and $R=10^6$ cm.

In spite of this impressive total store of energy, it would not long maintain luminosity at the level at which observations of neutron stars could sensibly be expected: $L\approx L_\odot\approx 4\cdot 10^{33}$ erg/sec. The temperature of the star drops very quickly as the result of neutrino emission. Matter can be treated as transparent for neutrinos only outside the dense hot core.[35-38] The temperature of the neutrino photosphere[38] $T_{\nu_p}\approx 6\cdot 10^{10}$ °K at $T_n\approx 10^{12}$ °K, and its radius is $R_\nu=10^6$ cm. Estimating using relation (4.5), we obtain $L_\nu\approx 8\cdot 10^{51}$ erg/sec. Then the characteristic cooling time is

$$\tau_\nu=\frac{Q_q}{L_\nu}\approx 25\ \mathrm{sec}. \quad (6.4)$$

At $T_n\approx T_\nu\approx 10^{10}$ °K, the interaction cross section of neutrinos with matter decreases $\propto T_\nu^2$, and the matter becomes transparent. The cooling rate decreases as the temperature drops further; neutrino luminosity gives way to photon luminosity. The magnetic field and the decrease in heat capacity due to the superfluidity of the neutrons are important factors in the rate of photon cooling. The interaction cross section for quanta corresponding to an extraordinary wave with frequency ν and moving along the field decreases sharply, $\sigma=\sigma_0(\omega/\omega_B)^2$ with $\omega<\omega_B$ and $\omega_B=eB/m_ec$,[57] which shortens the cooling time appreciably. According to calculations made in Ref. 58, the luminosity of the neutron star drops below the solar luminosity after a time $\tau_{ph}\approx 400$ years in a uniform field $B\approx 4\cdot 10^{13}$ G with allowance for superfluidity, which is much smaller than the cooling time $\tau_{\nu_p}\approx 10^3$ yr obtained[59] without consideration of the field or superfluidity.

The oscillatory energy of the star could also become an energy source. Generation of waves in the hot atmosphere of the neutron star would result in conversion of oscillatory to thermal energy and give the star a relatively high x-ray luminosity. But even this energy source dries up quickly as a result of nonequilibrium weak-interaction processes unfolding within the star.[60] The oscillations are damped because their energy is carried away by neutrinos and transformed into heat. The damping time of the oscillations due to reactions in which electrons, protons, and neutrons participate is estimated at about 100 yr,[60,61] and if there are hyperons at the center of the star, dissipation of the oscillatory energy into heat would occur during a vanishingly small time, ~1 sec.[62]

Thus, the neutron star quickly loses both its thermal and its oscillatory energy.

Rotation, which gives pulsars radio luminosity for up to 10^9 yr,[63] is an important source of neutron-star energy. A neutron star in a close pair is a powerful x-ray source that is luminous by virtue of accretion.

Observations indicate that single neutron stars may also be x-ray sources. The pulsars in the Crab and Vela nebulae emit in the x-ray as well as the radio bands, and much of the x-ray flux is constant and may be of thermal nature. The age of the Vela pulsar, ~10^4 yr, is such that its intrinsic thermal energy and oscillatory energy could no longer explain the observations. Many unidentified x-ray sources[64] are faint objects whose luminosities can also be explained by emission from the surfaces of single neutron stars. Arguments in favor of a galactic nature for the faint x-ray sources were presented in Ref. 65. Evolution of the nonequilibrium layer considered above would result in a release of energy the stored amounts and release rates of which could explain the thermal luminosity of single neutron stars.

b) Neutron diffusion

The deviation from equilibrium in a layer of the crust of a neutron star consists firstly of the presence of a significant number of free neutrons and, secondly, of the nuclear composition of matter, which is far off equilibrium. The nuclei apparently form a crystal lattice[66] and can therefore be regarded as immobile. In contrast, the neutrons are free and can diffuse under the action of the concentration gradient and the force of gravity. The neutron concentration gradient may change sign in the nonequilibrium layer, and, specifically, is always directed outward (or equal to zero) at the outer boundary, since there are no free neutrons on the outside at $\rho < 10^{10}\ \mu_e$ g/cm³. The gravitational force acting on the neutrons is always directed toward the interior of the star. The pressure in the nonequilibrium layer of the crust is determined basically by electrons. The ratio of the pressures of the relativistically degenerate electrons P_e (3.1) and the neutrons in the nonrelativistic-degeneracy approximation P_n (3.2) is

$$\frac{P_n}{P_e} = \frac{5.3 \cdot 10^9 \rho_n^{5/3}}{1.2 \cdot 10^{15} (\rho/\mu_e)^{4/3}} = 2.8 \cdot 10^{-5} \rho^{1/3} \frac{x_n^{5/3}}{(1 - x_n)^{4/3}} \tag{6.5}$$

(where we have assumed $Z = A/4$ for the nuclei).

For the highest accepted value for the neutron weight concentration $x_n = 0.5$ and $\mu_e = 4(1 - x_n)$ for nuclei with $Q_n = 0$ at the inner boundary of the nonequilibrium layer at a density $\rho_2 \approx 3 \cdot 10^{11} \mu_e$ g/cm³, we have $(P_n/P_e)_{\max} \approx 0.3$. On the average, the role of gravity in neutron diffusion is dominant, especially at the outer boundary, where the neutron pressure is negligibly small. At the same time, large neutron density gradients may appear in dense regions of the nonequilibrium layer, and they may have an important influence on diffusion.

From the computational viewpoint, allowance for the term with $\partial n_n/\partial r$ in the diffusion equation increases its order and greatly complicates the solution. There are also fundamental difficulties associated with the need to specify additional boundary conditions at the layer boundary. This requires knowledge of the physical conditions outside the nonequilibrium layer and investigation of reactions with neutrons that take place outside it. All this would greatly complicate the problem, whereas semiquantitative results on the evolution of the layer and the luminosity of a neutron star can be obtained by retaining only the principal gravity-determined term in the diffusion equation. A solution of this kind was obtained in Ref. 67. The problem was solved in the plane approximation under static-equilibirum conditions defined by

$$P = \frac{GM_0 M}{4\pi R_0^4}, \quad r = \int_P^{P_2} \frac{R_0^4\, dp}{\rho G M_0}, \quad r \ll R_0; \tag{6.6}$$

here M is the mass of the nonequilibrium layer above a given radius r reckoned from the lower boundary of the layer, where $P = P_2$. The atomic weight and charge of the nuclei are determined from (5.5) and depend on the electron Fermi energy and hence also on density. Diffusion is accompanied by a redistribution of matter in the layer, and by loss of neutrons into the core of the star, as a result of which the density and chemical composition of matter in the layer are altered. As the density and electron Fermi energy decrease, a nucleus becomes capable of emitting an electron, and then, owing to the presence of free neutrons, of capturing approximately four neutrons to arrive at a state with $Q_n = 0$. Here part of the internal energy of the nucleus is converted into heat. Thus, nuclear-chemical energy of nuclei and neutrons is released when the density decreases as a result of diffusion of neutrons into the star. The transformation of nuclei must be taken into account in the diffusion equation, since some of the neutrons are attached to or detached from nuclei, and in the equation determining the energy balance in the nonequilibrium layer. A diffusion equation allowing for nuclear transformation is easily obtained in an approximation in which the reactions of neutrons with nuclei take place much more rapidly than the diffusion process. This condition is always satisfied with accuracy of a very high order.

Although the atomic weights of the nuclei may change, their number remains constant. This enables us to introduce the computationally convenient Lagrangian coordinate N_A—the number of nuclei in the layer measured from the lower boundary with $P = P_2$ to the local pressure P:

$$N_A = \int_P^{P_2} \frac{4\pi R_0^4 (1 - x_n)}{G M_0 m_A}\, dP, \quad N_A(P_1) = N \text{ — the total number of nuclei.} \tag{6.7}$$

The mass-related Lagrangian coordinates normally used cannot be used in this case, since diffusion is accompanied by mass redistribution. The masses of nuclei $M_A(P)$ and of neutrons M_n in the layer with pressures from P to P_2 are

$$M_A(P) = \int_P^{P_2} \frac{4\pi R_0^4 (1 - x_n)}{G M_0}\, dP, \quad M_n(P) = M_2 - M(P) - M_A(P); \tag{6.8}$$

here $M_2 = 4\pi R_0^4 P_2 / GM_0$ is the total mass of the shell with $P < P_2$, and $M(P)$ is defined in (6.6). The mass of the shell above the nonequilibrium layer is $M_1 = 4\pi R_0^4 P_1 / GM_0$. Relations (3.1), (5.5), and (6.6)–(6.8) fully define the static equilibrium of the shell at a given $x_n(P)$ or $x_n(N_A)$.

Let us derive the diffusion equation. In a layer of unit density and thickness dr there are

$$\xi = n_n\, dr + \left(\frac{A\, dN_A}{4\pi R_0^2}\right) \tag{6.9}$$

nucleons, where dN_A is the number of nuclei in the layer dr.

Both n_n and A change in the diffusion process, but N_A does not depend on time. The change in the number of nucleons in the layer during time dt equals the difference between the diffusion neutron fluxes J_n (cm^{-2} $\times$sec^{-1}) at the boundaries of layer dr during time dt. Thus,

$$\xi(t + dt) - \xi(t) = [J_n(N_A + dN_A) - J_n(N_A)]\, dt. \tag{6.10}$$

Expanding in dt and dN_A in (6.9) and (6.10) and dividing by $dN_A dt$, we obtain the diffusion equation

$$\frac{\partial}{\partial t}\left(n_n \frac{dr}{dN_A} + \frac{A}{4\pi R_0^2}\right) = \frac{\partial J_n}{\partial N_A}$$

with the boundary condition

$$J_n(t, N) = 0. \tag{6.11}$$

The flux J_n is defined in a system bound to the nuclei and is positive if the neutrons are diffusing into the star. It remains to express the flux J_n in terms of matter parameters. We use the first approximation of the kinetic theory of diffusion in a binary mixture for an elastic-sphere model[68]:

$$J_n = \frac{3}{4}\sqrt{\frac{6}{\pi}}\,\frac{n_n}{n_A\sigma_{nA}v_n}\,\frac{GM_0}{R_0^2},\quad n_A=\frac{\rho_A}{Am_p},\quad \rho_A=\rho-\rho_n; \tag{6.12}$$

here σ_{nA} is the total neutron-nucleus interaction cross section. The velocity v_n was assumed equal to the velocity of a neutron at the Fermi limit:

$$v_n = 1.6\cdot 10^5\rho_n^{1/3}\ \text{(cm/sec)}. \tag{6.13}$$

The computed results given below indicate that the luminosity, according to Refs. 58 and 59, the temperature is such that the neutrons are partially degenerate and the neutrino luminosity does not exceed the photon luminosity. The validity of formulas (6.12) and (6.13) is confirmed by calculations of the neutron diffusion coefficient with allowance for degeneracy on the basis of solution of the kinetic equation by the Chapman-Enskog method.[159]

Heat is released in the process of neutron diffusion and is transferred outward by conduction and radiated, while some of it is transferred inward and heats the star. Generally, the heat conduction equation must be solved with allowance for energy sources in the star. In our approximate formulation, it is sufficient to confine the discussion to the stationary case in which all the energy released is radiated. In this case, it is sufficient to calculate the energy of the nonequilibrium layer in order to calculate the luminosity L.

An expression for the luminosity of a neutron star that takes into account the changes in the gravitational, chemical, and kinetic energies of the nuclei and electrons is given in Ref. 67.

c) Results of calculations

A numerical solution of Eq. (6.11) taking into account (3.1), (5.5), (6.6)–(6.8), (6.12), and (6.13) was obtained in Ref. 67. The calculations were made for a neutron-star mass $M = 1M_\odot$ and the cross-section value $\sigma_{nA} = \sigma_0(1 + A^{1/3})^2$, $\sigma_0 = 10^{-23}$ cm^2 = 10 b. The cross-section chosen is characteristic for elastic scattering of neutrons by heavy nuclei in the presence of a large number of resonances.[69] A more definite choice of the cross section σ_{nA} is as yet impossible because of the lack of experimental data and calculations for nuclei with large neutron excesses. The basic results of the calculations appear in Figs. 4–7.

The calculations were made for the initial distribution $x_{n_0} = 0.5$ and $x_{n_0} = 0.1$. Figure 4 shows distributions of

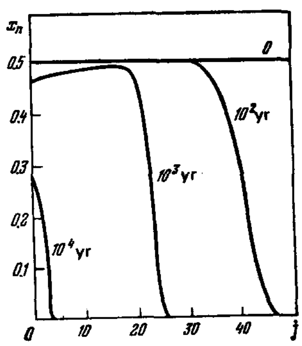

FIG. 4. Concentration of neutrons by weight x_n vs. number j of difference-grid node and at various times.

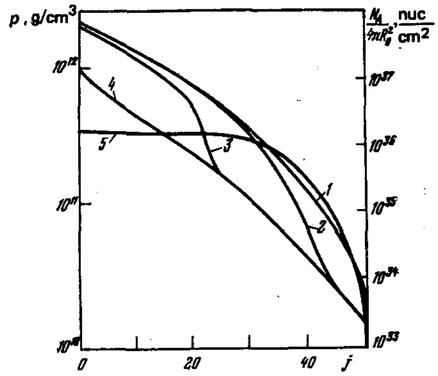

FIG. 5. Distribution of mass density ρ (g/cm^3) in the nonequilibrium layer at various times. t (yr) = 0 (1), 10^2 (2), 10^3 (3), 10^4 (4). Curve 5 represents the number of nuclei per square centimeter reckoned from the upper boundary of the layer as a function of difference-grid node number.

x_n over the layer at different times for $x_{n_0} = 0.5$, and Fig. 5 the analogous density distribution. The relation between the Lagrangian coordinate N_A and the node number j of the difference grid also appears in Fig. 5. At the starting time, neutron diffusion is strongest in the outer layers, and x_n and the density decrease rapidly there. This is because both the average random neutron velocity (6.13) and the ratio $n_A\sigma_{nA}/n_n \sim (x_A/x_n)A^{-1/3}$ are smaller in the outer layers, since A increases with decreasing density in accordance with (5.5). According to (6.12), therefore, the neutrons first leave the upper layers of the nonequilibrium region. The density change ceases when there are negligibly few neutrons. Figure 6 shows plots of the radii corresponding to a given coordinate N_A against time. The decrease in the number of neutrons due to diffusion results in a decrease in the mass above a given N_A and, consequently, the pressure P should decrease according to (6.6). Since P is determined by the electrons, the number of which increases only as a result of β decays, the layer must expand to maintain equilibrium. Then the internal stresses in the solid crust increase. Figure 7 shows the luminosity (L) curve of a neutron star during neutron diffusion in two variants: $x_{n_0} = 0.5$ and $x_{n_0} = 0.1$. The luminosity and the duration of the active phase are much smaller in the latter case.

We note that the duration of the active phase increases in proportion to the cross section σ_{NA}, but that there is a corresponding decrease in luminosity, in order that

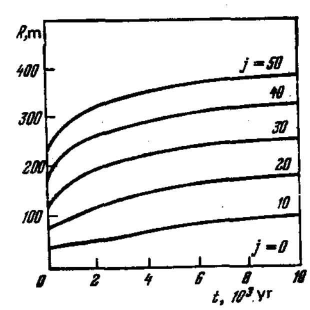

FIG. 6. Variation with time t (in years) of height R (in meters) of nonequilibrium-layer levels corresponding to various values of j.

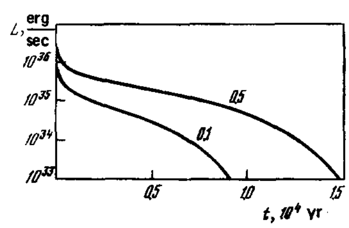

FIG. 7. Variation of neutron-star luminosity L with time t for initial distributions of neutron weight concentration. 1) $x_{n0} = 0.5$; 2) $x_{n0} = 0.1$; luminosity in erg/sec and time in years.

the total energy release should remain unchanged. The sharp drop in luminosity at the end of the active phase is apparently to be explained by our failure to take into account the neutron concentration gradient, which lowers J_n. Inclusion of this factor should result in slower energy release at the end of the process and smoother extinction of the neutron star.

d) Comparison with observations

The emission spectrum of a neutron star that shines by releasing energy in a nonequilibrium layer will apparently be near-thermal, and the surface temperature will be

$$T_S = \left(\frac{L}{ac\pi R_0^2}\right)^{1/4} \approx 3.5 L_{35}^{1/4} \cdot 10^6 (^\circ\mathrm{K}), \quad L_{35} = \frac{L}{10^{35}} \text{ (erg/sec)}. \quad (6.14)$$

The thermal-emission maximum occurs at the energy

$$\varepsilon_{\max} = 2.7kT \approx 0.84 L_{35}^{1/4} \text{ (keV)}.$$

Thus, sources that emit by nonequilibrium-layer decay and have luminosities $< 10^{35}$ erg/sec during most of their lifetimes (see Fig. 7) emit in the soft x-ray range, < 1 keV.

ANS satellite observations[70] detected nonpulsating luminosity of the Vela pulsar in the 0.2–0.28 keV range at a level of $8 \cdot 10^{-11}$ erg/cm^2 · sec · keV. Assuming a distance $R = 500$ pc, absorption corresponding to $N = 2 \times 10^{20}$ cm^{-2},[70,71] and an effective energy range of 0.25 keV, we estimate the luminosity at $2 \cdot 10^{34}$ erg/sec. Observations of the Vela pulsar[72] in the harder 2–60 keV range also indicated the presence of a nonpulsating component only. It is difficult to state a definite value for the luminosity of the pulsar in this range due to the strong nebular background. However, even if we assume that all the emission is determined by the pulsar, we obtain $8 \cdot 10^{33}$ erg/sec under the same assumptions. But the spectral distribution[72] indicates a nonthermal nature for the emission in the 2–60 keV range.

Comparison of the estimate $2 \cdot 10^{34}$ erg/sec with Fig. 7 ($x_{n0} < 0.5$) indicates that the luminosity of a star shining because of a nonequilibrium layer reaches the observed value after $\sim 10^4$ yr. Estimation of the age of the Vela pulsar on the basis of the rotation slowdown gives the very same value $\sim 10^4$ yr.[20]

The nonpulsating component of the emission from the region of the Crab Nebula pulsar is much stronger than the pulsating component because of the strong emission from the nebula. The pulsating flux in the 1–7 keV range[70] is $1.3 \cdot 10^{-9}$ erg/cm^2 sec, which, at $R = 2$ kpc, corresponds to $L \approx 10^{36}$ erg/sec over the entire frequency range, which is approximately twice the flux at 1–7 keV. Analysis of the nonpulsating thermal emission L_{th} of the neutron star in the Crab Nebula during a lunar occultation yielded the upper limit $L_{th} < 6 \cdot 10^{34}$ erg/sec, which corresponds to a temperature $T \leq 3 \times 10^6$ °K.[155] However, it must be remembered[155] that because of the strong absorption ($N_H = 2 \cdot 10^{21}$ at/cm^2), improvement of this upper limit will be very difficult. Comparison with Fig. 7 at $t = 10^3$ yr indicates that the initial neutron concentration in the nonequilibrium layer did not exceed $x_{n0} < 0.1$ for the σ_{nA} used in Ref. 67. The uncertainty as to σ_{nA} prohibits more definite conclusions as to the properties of the nonequilibrium layer.[8)]

If the dipolar magnetic field of a neutron star is large, it may produce differences in the emission at the pole and at the equator. The magnetic fields at the surfaces of the neutron stars in the Crab and Vela pulsars have been estimated at 10^{12} G.[20] Photons with frequencies below the Larmor frequency interact differently with matter as they propagate along and across the field. As we noted in Sec. 6a radiative heat conduction is considerably greater along than across the magnetic field because the cross section for interaction with electrons decreases $\propto B^{-2}$ for the extraordinary wave. The quantity $\hbar\omega_B \approx 12$ ($B/10^{12}$ G) keV also exceeds the characteristic energy (6.14) by a considerable margin. When energy is released in a thin nonequilibrium layer, the distribution of luminosity over the star in the magnetic-field case differs from the lumonisity distribution of a cooling star. In the latter case, the heat flux is at maximum in the direction along the magnetic field, and modulation of the luminosity is possible with rotation. If energy release occurs uniformly in the nonequilibrium layer, the heat flux across the surface will be uniform, but the presence of a magnetic field will result in modulation of the spectrum and polarization, which will be much more difficult to detect. The circular polarization of the radiation should increase sharply along the field, since it is basically the extraordinary wave that escapes and the cross section for this wave is small. The nonequilibrium layer may give rise to various nonstationary phenomena, which will be considered in the next chapter.

7. NONSTATIONARY PROCESSES INVOLVING THE NONEQUILIBRIUM LAYER

a) Stability of layer to buildup of oscillations

There are two factors that make for instability of the nonequilibrium layer. Firstly, the outer boundary of the layer is oscillatory-unstable. Expansion of the layer, during which the density of the outer layer drops below $\rho_1 = 10^{10}\mu_e$ g/cm^3, results in rapid spontaneous nuclear fission and an increase in thermal energy, slows down the pressure decrease, and weakens the restoring force.

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8) Trümper[150] (see also Ref. 70) observed that the x-ray band from 1 to 7 keV contains the minimum of the pulsating-radiation fraction of the central Crab Nebula source, apparently because the extended and point sources have spectra of different kinds.

Thus, the release of energy on expansion promotes buildup of oscillations at the outer boundary of the layer.

The other factor tending to develop unstable oscillations is the small value of the adiabatic exponent for oscillations in which partial equilibrium is maintained among nuclei, neutrons, and electrons and formula (5.5) is satisfied. In this case, the approximate equilibrium adiabatic exponent that takes account only of electron pressure was calculated as follows.

We obtain from expressions (3.1) and (5.5), the expression for $\rho = n_A A m_p/(1-x_n)$, $\mu_e = A/Z(1-x_n)$, and the condition $d(\rho/n_A)=0$

$$\left.\begin{aligned} d\ln P &= \tfrac{4}{3}(d\ln\rho + d\ln A),\\ d\ln A &= \tfrac{1}{4}\tfrac{28}{A}\left(33-\tfrac{A}{28}\right)d\ln P. \end{aligned}\right\}$$

From this we obtain at once

$$\gamma = \left(\frac{d\ln P}{d\ln\rho}\right)_{S=0} = \frac{4/3}{(2/3)+(308/A)}. \tag{7.1}$$

The composition was assumed to be at equilibrium in the calculation of (7.1). The exponent γ then reaches 0.091 at $A=22$ on the inner boundary of the layer. This small value of γ may result in building up oscillations of the inner layer boundary with the corresponding characteristic period

$$T \approx \frac{2\pi}{\sqrt{4\pi G\rho}} \approx 10^{-2}\ \text{sec for}\quad \rho = 10^{12}\ \text{g/cm}^3. \tag{7.2}$$

Owing to the finite rate of the β processes maintaining the equilibrium A of (5.5), deviations from the equilibrium A arise during the oscillations. The latter are damped by neutrino emission and conversion of oscillation energy into heat in nonequilibrium β processes.[51–54,74] Quasistationary oscillations with the period (7.2) may be generated as a result.

We note that instability at the outer boundary results in decay of nuclei in the outer regions and in an approach to complete thermodynamic equilibrium. This instability ceases to operate when the equilibrium layer becomes sufficiently thick and when strong density perturbations that do not in fact occur would be needed to cause instability.

b) Abrupt changes in period accompanied by a decrease in speed of rotation

Diffusion of neutrons into the interior of a star lowers the mass of the nonequilibrium layer and increases its thickness. The nuclei form a crystal lattice; therefore, stretching of the layer will deform the lattice and give rise to internal stresses that resist stretching. When the internal stresses exceed the strength limit, the crystal breaks, the internal stresses are relieved abruptly, and the layer thickness increases suddenly to the equilibrium value. This results in an abrupt increase in the moment of inertia and a decrease in the speed of rotation. Abrupt increases of period have been observed in the Crab pulsar.[75] The maximum relative slowdown $\eta_- = \Delta\omega/\omega \approx -10^{-9}$ in smaller than the relative acceleration $\eta_+ = \Delta\omega/\omega = 2.5\cdot 10^{-9}$ observed in the same pulsar.[76 9)]

Even sharper period jumps accompanied by increases in angular velocity have been observed in the Vela pulsar: $\Delta\omega/\omega \approx 2\cdot 10^{-6}$.[77–80] Jumps that increase angular velocity are usually assoicated with the accumulation of internal stresses and rupture of the crust or of the crystalline neutron core owing to deformation of the star during a process of slow angular-momentum loss.[81,82] A sharp increase in angular velocity might also result from a rapid decrease in the thickness of the shell due to neutronization.[73,83] Neutronization also occurs as a result of a density increase in the shell as angular momentum is lost. Thus far, the only possibility for the appearance of period jumps that result in both acceleration and deceleration of rotation has been the mechanism proposed by Scargle and Pacini[84] of particle accumulation in the magnetosphere of the pulsar and their sudden ejection as a result of development of instability. However, this mechanism does not explain the observed difference between jumps of different signs: for the Crab, η_+ is larger than η_-, and only accelerations of rotation have been observed for the Vela pulsar. An essential defect inherent in this mechanism was noted by Shklovskiĭ.[85] The density of hot particles accumulated in the magnetosphere is so high that their recombination emission reaches 10^{41} erg/sec with recombination times of $\sim 10^{-10}$ sec; this is clearly in contradiction to observations.

The expansion of the nonequilibrium layer associated with neutron diffusion takes place most rapidly during the first $\sim 10^4$ yr: therefore jumps with increasing period occur much more often in young pulsars than in old ones. This is evidently why jumps with an increase in period have not been observed in the Vela pulsar.

A distinctive feature of nonequilibrium-shell expansion is the fact that the outer layers are deformed first and the deformation moves inward with the passage of time (see Fig. 6). Since most of the mass is situated in the inner layers, fractures should occur more often and be less powerful in young neutron stars. With time, the frequency of breaks decreases, but η_- increases. If such jumps should be observed in the Vela pulsar, η_- should be larger than in the Crab pulsar. We therefore posit different origins for jumps associated with decreases and increases in period. In the former case, the star is quickly compressed as a result of the starquakes caused by rotation slowdown or as a result of neutronization due to the same cause. In the latter case, there are also starquakes, but they are accompanied by an abrupt expansion of the crust and caused by diffusion of nonequilibrium neutrons into the star.

If the star is rotating, expansion of the layer results in the appearance of shearing deformations $\Delta\varphi$, but they are very small for the observed pulsar rotation

9)Groth[164] presents the law of variation of the period in the form of stochastic variations of both signs that are superposed on a smooth period increase. Below we shall refer to all fast increases in period as jumps for the sake of brevity.

speeds.[20] Using the estimate of Ref. 81, we have $\Delta\varphi \approx (\Omega/\Omega_{max})^2\ \Delta R/R \ll 10^{-3}$ and $\Omega_{max} \approx 10^4$ sec^{-1}. The linear deformation is much larger than $\delta \approx \Delta r/h$, where h is the thickness of the layer. A linear deformation $\delta = 10^{-2}-10^{-3}$ is accumulated (see Fig. 6) in 2–0.2 yr during the first hundred years and 30–3 yr after a thousand years. If we take $10^{-2}-10^{-3}$ as the maximum relative deformation[81,82] for both the shear and stretching, the values quoted above give the time between crust fractures and starquarks. Let us estimate the size of the period jumps. If the thickness of the crust has increased abruptly by Δr, we have, given conservation of angular momentum,

$$\frac{\Delta\omega}{\omega} = -\frac{2\Delta m\,\Delta r\cdot R}{I_0}; \tag{7.3}$$

Δm is the mass fraction of the crust where the break occurs and Δr is the increase in the thickness of the layer with mass Δm. Let us assume $\Delta m = \varepsilon m$, where $m = 2\cdot 10^{29}$ g is the maximum mass of the nonequilibrium layer, $\Delta r = \varepsilon\delta r_0$, where $r_0 = 240$ m is the initial radius of the nonequilibrium layer, $I_0 = 10^{45}$ for the central moment of inertia of the neutron star, and $R = 10^6$ cm for the radius of the star. Then

$$\eta = \frac{\Delta\omega}{\omega} = -10^{-5}\varepsilon^2\delta. \tag{7.4}$$

With $\delta = 10^{-2}-10^{-3}$ and $\varepsilon = 0.1$, we have $\Delta\omega/\omega = -(10^{-9}-10^{-10})$ i.e., close to what is observed in the Crab. The upper limit of η_- for $\varepsilon = 1$ is $|\eta_-| < 10^{-7}-10^{-8}$ which is more than an order of magnitude smaller than the value $\eta_+ = 2\cdot 10^{-6}$ observed in the Vela pulsar.

c) Gamma-ray bursts from neutron stars

The gamma-ray bursts discovered in 1973[86–88] are among the most puzzling phenomena. Since the distances to these events are unknown, we have only relative observation data: total flux at the earth less than 2×10^{-4} erg/cm^2 in the spectral interval from 2 keV to 5 MeV; spectrum in this range with a maximum at ~150 keV; fine time structure in some bursts with details lasting ~0.1 sec. Detailed data on individual bursts will be found in Refs. 89 and 90, and there is also a catalog of ephemeral γ sources,[91] which gives the principal properties of known sources.

The distances to the sources and the total powers are definitely unknown. From the fact that their distribution is isotropic, we might conclude that they are either remote extragalactic ($\gtrsim$10 Mpc) or quite close ($\lesssim$300 pc) galactic sources. There are certain observational arguments in favor of a galactic nature for the bursts,[92] but they are not conclusive. Theoretical interpretation of the fine time structure of γ flares[89,90] would be very difficult if they are of extragalactic origin associated with supernova explosions.

Let us assume that the flares occur in nearby regions of the Galaxy. There are no optical objects in the directions of some γ flares[93]; we shall therefore suppose that the flares are associated with neutron stars, whose optical emission may be very weak. The appearance of a γ flare has been linked to a nuclear explosion occurring as a result of ejection of matter from the nonequilibrium layer to the surface at the time of a starquake.[94,95] There are certain other galactic-orgin hypotheses for γ flares, and they are reviewed in Ref. 96. Matter could be displaced from internal regions of the shell to the surface if the period jumps are related to starquakes (see Ref. 76) or to volcanic activity.[97] A recently proposed[98] starquake mechanism was based on the Pomeranchuck effect in a quantum neutron Fermi system in the presence of a liquid-crystal phase transition.

On the surface of a star, superheavy nuclei are unstable to β decays, which are the first and slowest processes; subsequent fission of nuclei and α decay take place much more rapidly and can be treated as instantaneous. The β-decay time can be written in the form $(\varepsilon_\beta - \varepsilon_{f,e} \gg m_e c^2)$

$$t_\beta \approx \frac{30ft}{\ln 2}\frac{(m_e c^2)^5}{(\varepsilon_p - \varepsilon_{f,e})^5} = 50-10^{-4}\ \text{sec} \tag{7.5}$$

for $ft = 10^3-10^5$; $\varepsilon_\beta - \varepsilon_{f,e} = 5-20$ MeV.

The time t_β determines the time of nuclear-energy release; the known time scales of γ bursts, 0.1–30 sec, lie in the range (7.5).

The formation of γ quanta in a nuclear explosion takes place basically as a result of γ emission by excited nuclei formed in β decays, nuclear fission, α decays, and neutron captures. Some γ- and x-ray emission arises when deceleration occurs, but basically it arises as a result of synchrotron radiation of electrons formed in β decay. The energy of the synchrotron quanta is given by $E_B \approx 3\cdot 10^{-4} B\gamma^2$ keV; with $B = 10^{11}$ G and $\gamma = E_e/m_e c^2 \approx 6$, we have $E_B \approx 100$ keV, while γ quanta that appear as a result of nuclear fission have energies of 0.002 to 3 MeV.[99] This is in close agreement with the observed spectral range.[89] The energy of the γ emission accounts for ~10% of the total energy of the nuclear explosion. If the average distance to the neutron star is 300 pc and the average total flux at the earth is ~10^{-4} erg/cm^2, the energy release in a single event is $\varepsilon_\gamma \approx 10^{39}$ erg. The energy yield Q from conversion of neutrons in the nucleus amounts to 10^{-2} of the rest energy M_0c^2, and equals ~$10^{-3}M_0c^2$ for superheavy nuclei. For the mixture we shall assume $\eta = Q/M_0c^2 \approx 3\cdot 10^{-3}$. The mass necessary for occurrence of a γ burst in a nuclear explosion is

$$M_\gamma \approx \frac{10}{3\cdot 10^{-3}}\frac{E_\gamma}{c^2} \approx 3\cdot 10^{21}\ \text{g}. \tag{7.6}$$

This mass is small compared to the total mass of the nonequilibrium layer, ~10^{29} g (see Sec. 5); thus, one neutron star may produce many γ bursts, and we may expect repeated bursts from the same part of the sky.

There are two possible variants of development of the process after the nonequilibrium layer appears on the surface of the star. The energy release may be sufficient to eject matter from the star to infinity. The gravitational potential at the surface of a neutron star is ~$0.2c^2$, so that either an additional energy source, for example rotation or a magnetic field, is required or the consumption of the nonequilibrium matter takes place ~100 times more rapidly. The γ radiation from the expanding matter depends on the optical thickness τ. If $\tau \gg 1$, the radiation is nearly black-body; when the fragment of matter is larger, its surface temperature T_S is low and γ emission does not occur. Thus,

γ emission can arise either in the initial stages of expansion, when the size is small and T_S is high, or in the last stages, when $\tau \lesssim 1$ and γ photons are not transformed to soft x-ray and optical radiation. An x-ray burst might be looked for in the intervening interval. Observed x-ray bursts[100] repeat quite often, and are therefore difficult to explain in terms of nuclear explosions.[10)] It is also unclear how γ flares are related to hard x-ray flares.[101] It is not impossible that we are dealing here with phenomena of the same nature and that the emission spectrum is determined by the dimensions and optical thickness of the exploding fragment of matter at the time of maximum energy release.

If the fragment of matter cannot be ejected out of the neutron star, it will always have a small dimension d and a high surface temperature, and γ-emission will occur. A hot spot on the surface of the star with $T = 5 \cdot 10^8$ °K and $d = 10^4$ cm must be formed to produce a γ flare similar to what is observed. In reality, we may expect the mean of these two cases.

The spectrum of the γ flare can be approximated either by a power-law or by an exponential curve.[89] The spectrum formed in a nuclear explosion is compatible with observations. It should be a mixture of a black-body spectrum with $T \lesssim 5 \cdot 10^8$ °K, a softened γ spectrum of a nuclear explosion, and the synchrotron emission of decay electrons in the x-ray range, $E \lesssim 100$ keV. Attempting to fit the observed spectrum exactly would make no sense because of the existing uncertainties and the broad range of options. We note that in this mechanism of γ-burst formation, polarization can be present only in the x-ray band and should be absent in the harder region. Absence of polarization in the x-ray region is also consistent with the model. Other models of γ-ray bursts[102,103] that are related to jumps in neutron stars assume a synchrotron nature for the γ radiation and require polarization over the entire spectral range. X-ray and γ-nuclear lines can exist in the present model in the hard range $E \gtrsim 200$ keV, but they have not yet been observed.[89]

The temporal fine structure and the presence of peaks of width ~0.1 sec[89] can also be explained in this model. The temporal fine structure of a γ-ray burst can be explained by several successive ejections of matter or by the appearance of several successive hot spots. Another cause of fine structure might be the presence of several successive β decays in the same fragment of matter, after each of which only part of the nuclear energy is released. This explanation fits observations, since the heights of the peaks and the distances between them are quite random. The mechanism of formation of the peaks by rotation of an electron cloud around the neutron star would give a more periodic peak pattern.[102,103]

Thus, the formation of γ-ray bursts in nuclear explosions on the surfaces on neutron stars can explain the energetics, statistics, and spectral and time features of this phenomenon.

d) Emission from the Crab Nebula

Matter of the nonequilibrium layer, including free neutrons, may be ejected in starquakes. Shklovskiĭ[85] showed that ejection of neutrons from the Crab pulsar could explain the emission from the Crab nebula and the observed correlation between the period jumps and the increase in the brightness of the nebular filaments. Estimates of the charged-particle flux out of the polar caps of a magnetized neutron star[85] have shown that the flux of these particles is approximately half of what would be required to make the nebulosity emit light. Ejection of charged particles is impossible away from the poles owint to the magnetic-field obstacle. Ejection of neutrons eliminates this difficulty. The neutrons pass freely through the region with the strong magnetic field, decay after $\tau \approx 13$ min at a distance $l \approx 0.4t \approx 10^{13}$ cm, are accelerated to high energies by the strong wave from the pulsar, and cause the nebulosity to emit.

This model requires a nearly continuous neutron flux from the pulsar to support emission from the nebula. It would hardly be possible for this flux to have the form of a continuous wind-type outflow. In all probability, small jumps are taking place continuously in the crust of the neutron star and causing quasicontinuous ejection of neutrons. The constant irregularities in the variation of the period of the Crab pulsar, observations of which are described in Refs. 84 and 164, may be associated with such small jumps. They can be interpreted as continuous small breaks in the outer part of the crust, where little mass is concentrated. In this case, we should observe a definite correlation between period irregularities and the emission from the nebula in much the same way as the "large" jump is related to an increase in nebular-filament activity.

8. ROLE OF NEUTRON STARS IN THE NUCLEOSYNTHESIS OF ELEMENTS

a) Processes of formation of elements

According to present-day concepts, most of the elements heavier than iron are formed by neutron capture with subsequent β decays.

A small number, about 30, of so-called bypassed elements with an excess of protons and concentrations of 10^{-2}–10^{-3} relative to their neighbors, is exceptional. It is believed that these elements may be formed when strongly excited nuclei lose neutrons with higher probability than protons. The excitation may be due to cosmic rays, energetic protons, and α particles.[104] Domogatskiĭ and Nadezhin recently developed an original model of the formation of bypassed elements as a result of interaction between ordinary nuclei and a neutrino pulse from the collapsing core of a star.[105]

Let us turn to the central problem of the synthesis of heavy elements beyond the iron peak. The following picture of heavy-element formation has existed since the classical paper of Ref. 106 was published. Two characteristic processes are distinguished: the s (slow)

[10)] X-ray bursts may occur on nuclear explosion of matter that has been delivered by accretion to the surface of the neutron star.[165]

process and the r (rapid) process, in which s and r elements, respectively, are formed. Most elements can be formed in either process. Neutron capture is the basic common feature of both processes. In the s process, the time between two successive neutron captures is much larger than the β-decay time. The characteristic β-decay times determined for Kr^{79} and Sm^{151} are 10^5 yr and 10 yr, respectively.

In the r process, successive neutron captures continue until the β-decay rate comes to exceed the rate of neutron capture. The characteristic conditions under which the r process occurs can be stated as follows[116]: a free-neutron concentration $n_1 \sim 10^{20}-10^{30}$ cm^{-3} at a temperature $\sim 10^9$ °K. Thus, the heaviest nuclei are formed in the r process. We note that elements heavier than Bi^{209} cannot form in the s process, since the isotope Bi^{210}, which is formed as a result of neutron capture by the Bi^{209} nucleus, is unstable to α decay and is transformed to Pb^{206}.

The paths of the s and r processes are shown in Fig. 8 in accordance with Ref. 129; also indicated is the formation pathway of the elements at high neutron concentrations, which we shall call the n process. The n process proceeds along the stability boundary of the nuclei, where $Q_n = 0$. We first proposed this process for the formation of heavy elements in neutron stars in Refs. 33, 34, and 50. The difference from the traditional r process is that the ordinary r process terminates at $A = 270-275$ and $Z = 93$ as a result of induced nuclear fission under the action of neutrons.[129]

In the case of high neutron concentrations, the n process can yield large values of $A \approx 300$, all the way up to the "stability island." We note that Berlovich and Novikov[107] obtained $Z = 107$ and $A = 291$ at $n_n \geqslant 10^{27}$ cm^{-3} when they allowed for an increase in binding energy for closed neutron shells.[107-109]

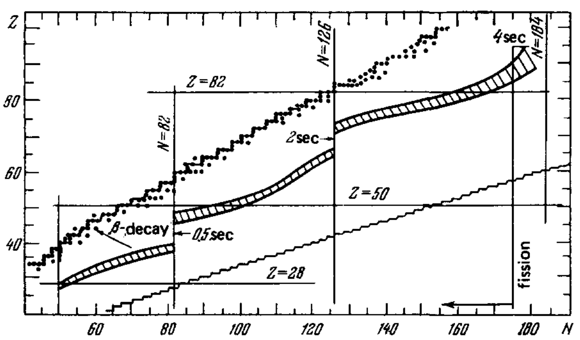

FIG. 8. Pathways of formation of elements on $N-Z$ diagram. The upper line indicates the path of the s process along the β-stability valley of the elements. Below it the shaded strip corresponds to the r process. The times indicated are those required to attain the respective concentration peaks. The numbers for the r process with the parameters $T_9 = 1.0$, $\lg n_n = 24$, and initial $Z = 26$ iron nuclei were taken from the calculations of Ref. 129. The plotted points represent stable nuclei formed in the r process. The lower stepped line describes the n process. It corresponds to the limiting r process at high neutron concentrations $\lg n_n > 30$, and its position does not change with increasing n_n. Nuclei formed in the n process fall into the stability valley as a result of β decay and fission with emission of neutrons.

The last element that has been synthesized on earth has $Z = 106$ and $N \sim 152$ and an extremely short lifetime, and is identified by spontaneous fission.[130] Closed-shell theory has yielded indications that elements with $Z = 114$ and $N = 184$ and $Z = 126$ and $N = 184$ may exist with rather long lifetimes with respect to α decay and fission.[126,127] In recent years, such elements have been looked for very actively in cosmic rays,[131,132] lunar and meteoritic rocks,[130,134,136] and in terrestrial materials.[133,135] Special interest attaches to recent results from a search for traces of superheavy elements in phenocrysts in Madagascar mica, which produce giant halos due to α decay. According to the results of Ref. 111, these traces of elements with $Z = 126$ and certain others have been found. However, no confirmation has as yet been obtained from experimental groups engaged in similar exploration. Further, explanation of the results of Ref. 111 that do not invoke the existence of superheavy elements have been offered.[112,114] The unfortunate result is that no hard, definite conclusions can as yet be drawn as to the existence of superheavy elements.

b) Possibilities of nucleosynthesis in supernovae

The relation of the processes described above to specific astrophysical objects is an important question. We shall not dwell upon the s-process, although certain difficulties are encountered with it. Our main emphasis will be on the r-process.

Figure 9 is a graph of the abundances of the elements that we have taken from Ref. 116. Figures 8 and 9 show that we have three peaks, with $N = 50$, 82, and 126, for typical r elements. It was shown already in 1957[106] that they are produced under different physical conditions. Thus, if we consider the formation of the three element peaks at the same neutron density $n = 10^{24}$ cm^{-2}, temperatures $T_9 = 1.45$, 1.0, and 0.8, respectively, are required for the peaks, whereas a single temperature $T_9 = 1$ requires different neutron concentrations: $n = 8 \cdot 10^{19}$, $3 \cdot 10^{23}$, and $1 \cdot 10^{26}$ cm^{-3}, respectively.

Figure 8 shows the characteristic cycle times necessary for formation of these peaks at a given neutron flux. We see that the characteristic time $\tau = 4$ sec for formation of the peak with $N = 126$ is too large compared to the time $\tau = 0.5$ sec for the $N = 50$ peak. This results

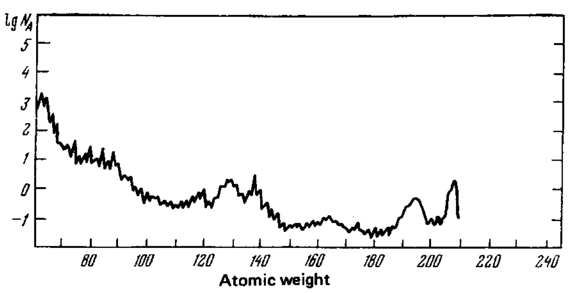

FIG. 9. Relative abundances of nuclides (nuclei with a given atomic weight) in the Solar System as functions of atomic weight A.[153] Normalized to $N_a(\mathrm{Si}) = 10^6$. The concentration peaks occur at $A = 138$ and 208 (s peaks) and at $A = 130$ and 198 (r peaks). The abundance peak at $N = 50$ corresponds to the plateau at $A = 80-90$.

in complete disappearance of the peak with $N=50$ at a characteristic time $\tau=4$ sec.

It is therefore difficult to obtain all three peaks in a single process. In nature, the r-process probably occurs in various astrophysical objects.

It is now believed that the conditions necessary for the r-process are realized in supernovae. It is difficult at this time to draw any definite conclusions as to the role of supernovae in creation of heavy elements, since there is no definitive model of a supernova. It is especially difficult to explain ejection of matter that has passed through high densities and in which many free neutrons have been formed. We shall not spend time on this problem because it requires a separate analysis. We note only that the few presently available rough calculations of supernovae[23,39,47] yield inadequate neutron concentrations in the exploded matter. Difficulties in explaining the origin of heavy elements on this basis have been pointed out.[116]

Definite progress has recently been made in the theory of supernovae.[166] The theoretical models that have been derived encourage the belief that formation of heavy elements may be possible in the central regions of expanding supernova remnants.

c) Nucleosynthesis of heavy elements with participation of neutron stars

Neutron stars are astrophysical objects associated with possible nucleosynthesis of r-elements. Formation of heavy elements in a neutron agglomeration was first discussed by Mayer and Teller.[118] This idea has recently been developed further. There are two models relating to the role of neutron stars in nucleosynthesis. Under the hypothesis of Ref. 128, a neutron star in a close binary system disintegrates if the other component is a black hole with large mass. The possibility of heavy-element nucleosynthesis in an outburst of this kind has been noted.[120,121] It seems to use that the formation of close binary systems of this type is improbable.

The other alternative is ejection of matter from the nonequilibrium layer of a neutron star, to which Secs. 4 and 5 were devoted.

The mechanism by which matter is ejected from the nonequilibrium layer has been discussed in the preceding sections. However, it is not now clear whether the ejection would be explosive or whether the matter would be effused with a characteristic hydrodynamic time. This would determine the chemical-evolution pathway of the matter.

Let us first consider a situation in which matter is ejected explosively from the neutron star. The agglomerations of matter containing superheavy elements expand rapidly after ejection from the neutron star,[94,95] so that the time t_H to expand to very low density is much smaller than the β-decay time t_β and the nucleus undergoes β decays from the boundary $Q_n=0$ until it has been transformed into a stable nucleus or has decayed (by α decay or fission). To estimate the expansion time, let us assume that an agglomeration with a diameter of ~10^3 cm has been ejected and that its outer boundary expands at a velocity close to the velocity of free fall v at the surface of the neutron star.

Expansion at constant velocity, which is characteristic of adiabatic expansion of a sphere in vacuum,[110] gives

$$\frac{1}{\rho}\frac{d\rho}{dt}=-\frac{3v}{r},\quad v=\frac{c}{3}\frac{r}{10^3+(ct/3)},\quad \rho=\rho_0\left(1+\frac{ct}{3\cdot 10^3}\right)^{-3}.$$

Then the characteristic time for expansion from a density of 10^{10} g/cm^3 to $\rho\sim 10^{-6}$ g/cm^3,[11)] when the r-process ceases, is ~$2\cdot 10^{-2}$ sec, which does not exceed the β-decay time. Thus, matter that has been ejected from a neutron star and is expanding at a high velocity due to a nulcear explosion[94,95] would expand very rapidly, and superheavy elements present in it would enter directly into the interstellar medium.

We can examine a situation with slow hydrodynamic expansion of matter from the nonequilibrium layer of a neutron star. Then the expansion would be accompanied by chemical evolution of matter in accordance with the ratios of the characteristic times: of expansion, β-decay, radiative neutron capture, and spontaneous fission (t_{fis}). Let us consider a model of this evolution from Ref. 117.

Let matter consist of nuclei (A, Z) and neutrons. As was shown in Sec. 5, the nuclei (A, Z) are on the stability boundary. Entering the atmosphere of the neutron star, the ejected matter expands. The electron Fermi energy $\varepsilon_{f,e}=mc^2(\rho/\mu_e\cdot 10^6)^{1/3}$ erg decreases. As Fig. 2 shows, the charge of the nuclei increases as a result of β decay and, finally, the nuclei enter the fission region. The neutron-capture and β-decay processes continue as long as $t_{\text{fis}}>t_\beta$. We note that the rate of neutron capture is always higher under these conditions than the β-decay rate, i.e., the nuclei are always on the stability boundary.

The evolution of nuclei during such expansion of matter was calculated in Ref. 117. The following assumptions were used. The expansion time is $t_H\sim 446/\sqrt{\rho(\text{g/cm}^3)}$ sec, the radiative neutron capture time is $t_n=m_p/\alpha\rho\langle\sigma v\rangle$, where m_p is the nucleon mass and α is the fraction of neutrons; the β-decay time is $t_\beta=6\Delta/10^{-4}W_0^6$ sec, where Δ is the level density in MeV^{-1} (Ref. 129) and W_0 is the β-decay energy in MeV with the condition $W_0\gg\varepsilon_f$ and $\lg(ft)=5$.[129] The parameter values used were $\sigma=10^{-25}$ cm^2 for the neutron capture cross section, $v=10^8$ cm/sec for the neutron velocity, and $\Delta=15$.

The half-life of the nuclei was determined, following Ref. 137, from

$$T_{1/2}=10^{-21}\cdot 10^{7.85E}\ (\text{sec}),$$

where $E=[19.0-0.36\,(Z^2/A)+\varepsilon]$ (MeV) and $\varepsilon=0.7$ for doubly even, $\varepsilon=0.4$ for even, and $\varepsilon=0$ for the other nuclei.

A computer was used to calculate the evolution of the

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11)A density value of 10^{-6} g/cm^3 is obtained from the formula $n=1/\langle\sigma v\rangle\,\tau$ with $\langle\sigma v\rangle=3\cdot 10^{-17}$ cm^3/sec[115] and $\tau=t_\beta=3\cdot 10^{-2}$ sec.

nuclei in accordance with the above equations. It was found that t_{fis} becomes smaller than t_β when the density becomes $\rho \sim 5 \cdot 10^6$ g/cm^3, irrespective of the initial conditions in matter. The rate of radiative neutron capture is always higher than the β-decay rate under the conditions indicated.

Recession of the fission limit to $A \approx 350$ owing to an excess of neutrons was reported in Ref. 119. Such nuclei are quite capable of existing under nonequilibrium-layer conditions. However, nuclei with still larger A may be formed in the ejection process under kinetic conditions (it is only necessary that $t_{fis} > t_\beta$). Although even $A = 350$ is sufficient for formation of "stability island" nuclei after a few β decays, fission of nuclei with larger A may also produce "stability island" nuclei directly in the fission prosess. Thus, matter expelled from the neutron star, and together with it, the heavy nuclei that have been formed, enter the atmosphere of the neutron star. Afterwards it can be dissipated in the interstellar medium by electrodynamic acceleration mechanisms. It may also be ejected utilizing nuclear-explosion energy.[94,95]

Let us evaluate quantitatively the possibilities of this heavy-element-forming mechanism in the Galaxy. We shall consider elements with $A > 150$, whose abundance on the sum amounts to 10^{-8} by weight, basically because of lead Pb^{206}, and elements with $A > 210$, whose total abundance is smaller by 2–3 orders of magnitude.[115] If we assume that a small fraction of the mass of the nonequilibrium layer, $\sim 10^{-3} M = 10^{-7} M_\odot$ from (5.13), is eventually transformed into heavy and superheavy elements, then 10^7–10^8 neutron stars in a Galaxy with a mass of $10^{11} M_\odot$ could have created a heavy-element abundance of 10^{-10}–10^{-11} by mass, which is sufficient to explain the elements with $A > 210$. But if the number of neutron stars reaches 10^9 and the efficiency $\sim 10^{-2} M = 10^{-6} M_\odot$ per neutron star, it would then also be possible to explain the origin of elements with $A > 150$ in terms of ejections from neutron stars.

d) Production of deuterium

Let us consider another crucial juncture in nucleosynthesis theory: that associated with deuterium production. The observed deuterium concentration $D/H \sim 10^{-4}$ (Ref. 122) appears to exceed the amount of deuterium produced in the hot model at the assumed cosmological density (Table IV). At the present time, the density of the Universe can confidently be put in the range (in a model with the cosmological constant $\lambda = 0$) $0.02\rho_c \leq \rho \leq 6\rho_c$ where $\rho_c = 3H^2/8\pi G = 6 \cdot 10^{-30}\ (H/55)^2$ (g/cm^3) and H is Hubble's constant in km/sec·Mpc. The lower bound is obtained from the total apparent mass of the galaxies and the upper bound from the condition that we do not observe a great deceleration in the expansion of the Universe. For the observed deuterium concentration in the interstellar medium to be of cosmological origin, it is necessary that the density of the Universe be $\rho < 0.05\rho_c$. This density is smaller than the average density obtained when the virial mass of the galaxies and the existence of haloes are considered.[139] Table IV lists deuterium concentrations for various densities of the Universe.

TABLE IV

Observations	Cosmic Rays	Nucleosynthesis of deuterium D in a hot universe[140]			
		ρ, g/cm^3	$5 \cdot 10^{-30}$	10^{-30}	$3 \cdot 10^{-31}$
$1.4 \cdot 10^{-4}$ *), $6 \cdot 10^{-3}$ **)	Small amount	D/H	$3.1 \cdot 10^{-5}$	$2.1 \cdot 10^{-5}$	$1.4 \cdot 10^{-4}$

*)According to Ref. 122.
**)From observations of HCN and DCN molecules.[151]

Because of this, the possibility of a noncosmological origin of the deuterium has been considered in recent years. In the most interesting model, deuterium is formed in supernova envelopes.[123] However, it has been shown[124,125] that these processes yield too high a Li content, one that is inconsistent with the observed lithium concentration in the interstellar gas. It is concluded in Ref. 152 on the basis of the calculated structure of the shock wave that nucleosynthesis of other light elements in supernova envelopes is also impossible.

It is possible for the observed deuterium concentration to have been produced without distortion of the cosmological nucleosynthesis of other elements in the evaporation of primary black holes.[138]

Another possibility arises in ejections from neutron stars, in which, together with heavy elements, free neutrons are expelled and may combine with protons that have entered the neutron star previously or are present in a more extended envelope to form deuterium without admixtures of other light elements. If we assume $\sigma v = 3 \cdot 10^{-17}$, $v = c/3$, and a neutron lifetime $\tau_n \sim 700$ sec, neutron capture requires a density of the surrounding protons $n_p = 1/\sigma v \tau \sim 5 \cdot 10^{13}$ cm^{-3}, which is possible with accretion of matter on a neutron star in a binary system. Formation of deuterium is also possible on the surface of a neutron star. Assuming that $\rho_H \leq 10^6$ g/cm^3 and that P_H and M_H are given by (5.9) and (6.7), the largest amount of hydrogen accreted on the star is $M_H = 0.1 \times P_H \approx 10^{22}$ g. The number of explosions during the lifetime of the nonequilibrium layer may be very large, $\sim 10^6$, and then a mass of 10^{28} g could be transformed into deuterium around the neutron star. In this case, deuterium concentrations of 10^{-3}–10^{-4} by mass could be reached in the neighborhood of neutron stars.[122,151] A total of $\sim 10^9$ neutron stars could produce an average deuterium abundance of $\sim 10^{-5}$ in the interstellar gas. Since an overabundance of deuterium is observed only in certain directions,[156] the hypothesis that it is produced in ejections from neutron stars merits attention.

9. CONCLUSION

It follows from the above exposition that nonequilibrium layers formed in the shells of neutron stars may be important in supporting the x-ray luminosity of single neutron stars, in nucleosynthesis of heavy elements and in explaining burst of γ radiation. To obtain more accurate estimates of the quantitative characteristics of the layer and of the phenomena associated with it, it is necessary to know the properties of heavy nuclei near the boundary $Q_n = 0$ and the cross sections for interaction of neutrons with such nuclei, and to calculate the kinetics of the nuclear reactions at temperatures transitional from nuclear equilibrium to a frozen number of nuclei. Solution of these physical problems will permit more definite calculations of various processes in the shells of neutron stars and of their observational manifestations.

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