yield important information on the structure of the nuclear surface. This circumstance was noted earlier in [6] and [2].

As indicated above, the cross sections for the excitations of the levels satisfying the condition (b) do not decrease with increasing energy [3]. Thus, formulas (6) and (7) connect the cross sections of the most noticeably excited levels in high-energy hadron beams with the corresponding electromagnetic-excitation cross sections.

We note in conclusion that the method used by us is actually equivalent to the eikonal approximation in the distorted wave method, which has been sucessfully used to describe elastic scattering of protons by C^{12} and O^{16} nuclei in [7, 8]. There, however, model formulas were used for the wave functions of the nuclei, and no connection was established between the cross sections of the electromagnetic and hadronic processes.

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NONEQUILIBRIUM COMPOSITION OF NEUTRON-STAR SHELLS AND NUCLEAR ENERGY SOURCES

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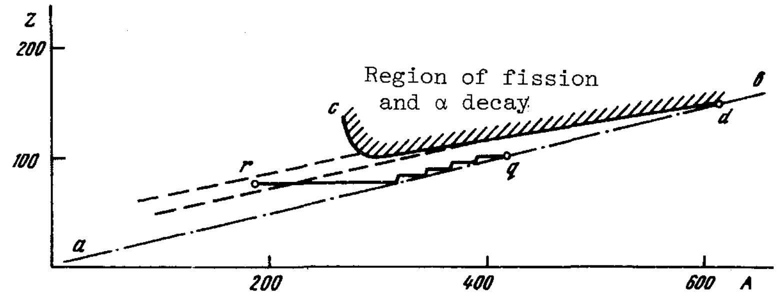
1. The question of the chemical composition of matter at densities $\rho = 10^{10} - 10^{14}$ g/cm^3, corresponding to shells of neutron stars, was considered recently [1, 2]. However, it is assumed in these as well as in earlier papers that the matter of the neutron star is in a lower energy state. We show in the present paper, on the basis of an analysis of the process of neutron-star formation, that the state of complete minimum of the energy is not reached in neutron-star shells at densities $10^{14} \leq \rho \leq 10^{11}$ g/cm^3. A large excess of nonequilibrium neutrons is produced, and their reserve of nuclear energy could ensure a neutron-star luminosity at a level $\sim 10^{36}$ erg/sec for $\sim 3\times 10^5$ years.

2. According to present-day concepts, the formation of a neutron star should occur during a collapse process accompanied by a subsequent expansion of the shell and observable supernova explosion [3]. The matter making up the shell of the neutron star should go during the collapse proces through a stage of temperatures T and densities ρ much higher than in a stationary star. During the stage of higher ρ and T there are produced neutrons and superheavy nuclei, the neutrons making up the bulk of the matter [4].

During the course of the expansion of the outer layers of the neutron stars and establishment of the stationary state, the nuclei rapidly capture neutrons, undergo β^- decay, and heavy nuclei are produced [5, 6]. At large A and Z, the nuclei become unstable against spontaneous fission and alpha-particle emission. An increase in the number of neutrons in the nucleus stabilizes these instabilities, but when the neutron excess becomes very strong the energy Q_n of detachment of the last neutron becomes equal to zero and further neutron capture becomes impossible. Under normal conditions, such nuclei are β^--unstable, but in the presence of degenerate electrons with large Fermi energy ε_{fe} the β^- decay is forbidden.

The figure shows on the (A, Z) plane the line ab where $Q_n \simeq 0$. The values of A and Z on this line were calculated for several nuclei in [7]; it is usually assumed [3] that $Z = A/4$ on



Production of superheavy nuclei. The nuclear evolution is represented by the stepped line rq. $Q_n = 0$ on the dash-dot line ab; the dashed line cd delimits the region of spontaneous nuclear fission; on the dashed lines, $\varepsilon_\beta = Q_p - Q_n$ = const and $\varepsilon_{\beta_1} > \varepsilon_{\beta_2}$. The line $Q_n = 0$ enters the fission region at the point d.

ab, just as for the He^8 nucleus. The line cd separates the region of nuclei that are unstable against fission and α decay and was drawn approximately with allowance for the data of [8], where the stabilizing effect of neutron shells on fission and decay of nuceli was demonstrated. Using a semi-empirical formula for the half-life [9] and the values of the proton energy Q_p on the boundary ab from [6], we obtain, putting $T_{1/2} \approx 3\times10^7$ years at the point d, the following parameters: Z 153, A 612, and Q_{p0} 8 MeV. We note that Q_p decreases along the line ab with increasing A and Z.

The time of neutron capture by the nuclei is much shorter than the β-decay time. Therefore an ordinary nucleus (A, Z) captures neutrons until a nucleus is produced at the boundary $Q_n \approx 0$. Further neutron capture is impossible, so that β^- decay takes place, after which neutrons are captured again. At a constant pressure corresponding to a given value of ε_{fe}, the process continues until $Q_p = \varepsilon_{fe}$ if $Q_{p0} < \varepsilon_{fe}$, or until the track of the nucleus goes off into the fission region at $Q_{p0} > \varepsilon_{fe}$. The track of the nucleus for the case $Q_{p0} < \varepsilon_{fe}$ is represented in the figure by the stepped line rq. Thus, nuclei with $Q_p = \varepsilon_{fe}$ are produced at $\varepsilon_{fe} > Q_{p0}$ and are stable under these conditions, so that an appreciable fraction of the neutrons remains stored in the shell. At $\varepsilon_{fe} < Q_{p0}$ the fission and α-radioactivity increases the number of priming nuclei, until all neutrons are transformed into nuclei and the matter reaches a state close to the minimum of the total energy. The limiting density ρ_1 corresponding to $\varepsilon_{fe} = Q_{p0}$ is determined from the known formula for a degenerate Fermi gas [3]

$$\rho_1 = \mu_e \cdot 10^6 \left(\frac{Q_{p0}}{m_e c^2}\right)^3 \approx \mu_e 4 \cdot 10^9 \ \text{g/cm}^3, \tag{1}$$

where μ_e is the number of nucleons per electron.

At densities $\sim10^{14}$ g/cm^3 the minimum energy corresponds to a neutron liquid with a small addition of protons and electrons [1, 2]. The upper density limit ρ_2 at which nonequilibrium excess neutrons exist is therefore lower than 10^{14} g/cm^3, and lies apparently in the range $\rho_2 = 10^{12} - 10^{14}$ g/cm^3.

3. Thus, the outer part of the neutron-star shell, $\rho < \rho_1$, is at the minimum, while the inner part, $\rho_1 < \rho < \rho_2$ has a large reserve of nuclear energy in the form of neutrons and superheavy nuclei. Let us estimate the stored energy. The shell mass M_{sh} of a neutron star of mass M and of radius $R = 10^6$ cm, contained between the densities ρ_1 and ρ_2 corresponding to pressures P_1 and P_2, is expressed by the simple formula

$$M_{sh} = \frac{4\pi R^4}{GM}(P_2 - P_1) \approx 0{,}1(P_2 - P_1) \ \text{g}. \tag{2}$$

Assuming that $\sim7\times10^{-3}$ of the rest mass[1)] is released when neutrons are transformed into nuclei, we obtain the energy reserve E_{nuc}

$$E_{nuc} \approx 6 \cdot 10^{17} (P_2 - P_1) \ \text{erg}. \tag{3}$$

Since the pressure in the shell is due to degenerate electrons, the pressure in the ultrarelativistic region is [3]

$$P \approx 10^{15} (\rho/\mu_e)^{4/3}.$$

[1)] In fact, the energy release is less, owing to the high ε_{fe}, and vanishes at a density ρ_2. This, however, does not change the order of magnitude of E_{nuc} from formulas (3) and (4).

For $\rho_2 = 10^{12} - 10^{14}$ g/cm^3 we have

$$E_{nuc} = (6 \cdot 10^{32}/\mu_e^{4/3})(10^{16} - 5 \cdot 10^{18}) \text{ erg.} \quad (4)$$

In view of the large number of neutrons, the quantity μ_e can be large: $\mu_e \approx 10$. Considering the average estimate, we obtain $E_{nuc} \approx 10^{49}$ erg. This energy reserve is sufficient to ensure a neutron-star luminosity at a level $10^{35} - 10^{37}$ erg/sec for $3\times10^4 - 3\times10^6$ years.

4. If the matter from the internal part of the shell somehow finds its way to the outer part, then the nuclear energy stored in the neutrons is immediately released within a time on the order of the time t_β of the slowest β^- decay among the reactions (see the figure). Owing to the large number of heavy-nucleus levels, the β^- decay can be regarded as allowed with respect to the quantum numbers with log ft = 3 - 5 and β^--decay energy $\sim$10 MeV; we then have at $\varepsilon_\beta >> \varepsilon_{fe}$

$$t_\beta \approx \frac{30ft}{\ln 2}(\epsilon_\beta/m_e c^2)^{-5} \approx (10^{-2} - 1) \text{ sec.} \quad (5)$$

The hydrodynamic time at $\rho = 10^{10} - 10^{12}$ g/cm^3 is

$$t_H \approx \frac{10^3}{\sqrt{\rho}} \approx (10^{-2} - 10^{-3}) \text{ sec.} \quad (6)$$

The transformation of neutrons into nuclei therefore takes place within a time longer than or of the order of the hydrodynamic time. Since energy is released when matter expands, the considered transition layer between shells will be unstable against the buildup of oscillations with a characteristic period $T \sim t_H \approx 10^{-2} - 10^{-3}$ sec. The oscillations in the transition layer and the ensuing release of nuclear energy are accompanied by the propagation of compression waves that pass through a medium of decreasing density, so that their amplitude increases and their emergence to the surface in the form of shock waves leads to heating, constant ejection of the outer layers of the shells, and enrichment of the space with plasma. Neutrons can go over to the outer shell also by diffusion, since a neutron-concentration gradient exists at the layer boundary, or by mixing of the convective type due to the thermal instability of the source in the layer.

5. The considered energy source can play an important role in the explanation of the emission from x-ray sources, particularly single sources [10] of the Sco-X-1 type. The irregular type of variability observed in this source may be due to different types of instability developing in the vicinity of the layer source, or to processes occurring in the vicinity of the neutron star and instabilities in the surrounding plasma. The luminosity of the source GX340 + 0 which is interpreted as a hot neutron star [11] can also be due to the reserve of nuclear energy in the shell.

No optical or x-radiation was observed from the majority of radio pulsars. In view of the possible energy sources, it would be of interest to observe the pulsating and non-pulsating radiation from pulsars in the soft x-ray and ultraviolet regions. Self-oscillations can enrich the pulsar magnetosphere with quasineutral plasma, in contrast to the enrichment with particles of like charge [12] by unipolar effects; this may be of fundamental significance for the mechanism of the pulsar radio emission.

We note in conclusion that the possible existence of nuclear energy sources in neutron stars was indicated in [13], but the concrete mechanisms of energy release mentioned in [13] are ot at all similar to those considered in the present article.

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THREE-REGGEON DESCRIPTION OF INCLUSIVE PROTON SPECTRA

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The study of $pp \to pX$ hadronic spectra near the kinematic boundary $s \gg M^2 \gg m^2$, $-t \lesssim m^2$ within the framework of the three-reggeon (TR) scheme [1, 2] is of interest both from the point of view of a test of scaling in inclusive reactions (i.e., the dependence of the spectra on only $x = 2p/\sqrt{s} \simeq 1 - M^2/s$) and in connection with the possibility of obtaining information on the three-reggeon vertices. Particular interest attaches to the three-pomeron vertex, the behavior of which at small t is essential for the understanding the character of the asymptotic behavior of strong interactions [3 - 5].

In earlier analyses, the data at $s = 20 - 60$ GeV^2 [6, 7] and the new ISR data at $s = 940$ and 1995 GeV^2 [8] were described differently assuming appreciable PPR and RRP contributions and small or nonexistence of PPP and RRR contributions (here P is the Pomeranchuk pole and R stands for the secondary trajectories P', ρ, ω, A_2, etc.).

It is shown in the present paper that a simultaneous analysis of these data leads to an entirely different relation between the vertices, namely, to appreciable PPP and RRR, the presence of RRP, and smallness or nonexistence of PPR.

The inclusive cross section is given in the TR model by [1, 2]

$$E\frac{d^3\sigma}{d^3p} = \sum_{ijk} G_{ijk}(t)(1-x)^{\alpha_k(0)-\alpha_i(t)-\alpha_j(t)}\left(\frac{s}{s_0}\right)^{\alpha_k(0)-1} = \Sigma(ijk)\,;$$

$$G_{ijk}(t) = \frac{1}{16\pi^2}\beta_i(t)\beta_j(t)\beta_k(0)\eta_i(t)\eta_j^*(t)\,g_{ijk}(t);\quad i, j, k = P, R\ldots \tag{1}$$

where $\beta_i(t)$ are the vertices of the connection of the Regge poles with

$$p\left(\sigma_{tot}^{pp} = \sum_i \beta_i^2(0)(s/s_0)^{\alpha_i(0)-1}\right),$$

$\eta_i(t)$ is the signature factor, $g_{ijk}(t)$ is the three-reggeon vertex[1)], and $s_0 = 1$ GeV^2.

A characteristic feature of the proton spectrum at $s = 940$ GeV^2 and $s = 1995$ GeV^2 is the sharp peak at $x \gtrsim 0.9$ and $0.2 \le p^2 \le 1$ (GeV/c)^2. This peak was described in [8] by the term PPR. However, the PPR contribution increases like $s^{-1/2}$ with decreasing s, and use is made of the data at $s = 20 - 60$ GeV^2, we can show that its upper bound at $s = 1995$ GeV and $x \simeq 0.9 - 0.97$ is $\le 15\%$. This makes it necessary to introduce PPP. As a background for the latter we use RRR and RRP and consider variants with and without PPR. The contributions of the RRP and PRR interference terms are neglected since, first, there are indications that they are small [9] and, second, our analysis shows that allowance for these terms can change g_{PPP} by only $\sim 20\%$.

We present the results of the data reduction. The quantities $G_{iik}(t)$ were parametrized in the form

$$G_{iik}(t) = G_{ik}(0)\,e^{R_{ik}^2 t},$$

and the values chosen for the Regge-pole parameters were $\alpha_P(0) = 1$, $\alpha_R(0) = 0.5$, $\alpha_P'(0) = 0$ and 0.15 (GeV/c)$^{-2}$ and $\alpha_R'(0) = 0.75$ (GeV/c)$^{-2}$. The decrease in the slopes takes into account the

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[1)] The quantities $g_{ijk}(t)$ are connected in our normalization with the r(t) used in [5] by the following relation: $g_{PPP}(t) = \sqrt{8\pi}\,r(t)$.