Protomatter and EHE C.R.

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We explore a novel aspects expected from an assumed change of properties of space-time continuum in the central part of AGN in density range far above the nucleus. Suggesting the physics of protomatter we present a detailed theoretical and numerical study of the models which are an essential microscopic generalization of the standard black hole accretion (SBHA) models of AGNs. This approach is of vital importance for two crucial elements required for the GZK air shower events as far it predicts: 1) the large flux of EHE neutrinos above $10^{21}\,eV$, even after the neutrino trapping in the superdense medium, produced by the predominant neutrino cooling of the AGN-protomatter core via simple or nucleon-modified URCA processes, and the pionic reactions; 2) the relic cosmic background neutrinos (CBN) with light mass $\sim 1\,eV$ produced in hot primordial Big Bang phase, forming the hot dark matter (HDM).

KEYWORDS: Space-time physics, a protomatter, BHs, AGNs, EHE C.R.

§1. Introduction

The SBHA models were became a generally acceptable paradigm as the view on AGNs assumed to consist of central super-massive BHs surrounded by an accretion disks (ADs) of $\sim 10^9 M_\odot$. They have a number of advantages, but could not be regarded as a final word in the AGN physics. Within respect to SBHA models, note that such an approach suffers from some principle difficulties, among them, in particular, the theory breaks down inside the black hole and it is impossible to compute corresponding integral characteristics of AGNs in the sense of microscopic theory. In order to innovate the solutions to problems of AGN and EHE C.R. physics here we develop on recent novel microscopic approach to standard BH physics .[1] Within the general gauge principle (GGP) [2] (sec.2,3) we explore a considerable change of properties of space-time continuum in the central part of AGN in density range far above the nucleus when the superdense protomatter core (SPC), consisted of the central protomatter core (PC) and outer layers of ordinary matter, is formed with the thermodynamical properties differed strongly from the thermodynamics of ordinary compressed matter (sec.4). The SPCs surrounded by the ADs are considered as the AGNs. The external physics of suggested models in early stage of formation of SPCs is absolutely the same of the SBHA models. But, a milestone of difference from SBHA models just are both the stable SCPs with a number of integral characteristics and available metric singularity cutoff (MSC) effect (sec.4). Therefore, in due course, in strong contrast to standard BH physics, the most important next stage of evolution of the SPCs succeeds. This approach enables an insight to key puzzle of EHE C.R. (sec.5) when the particles have got unique possibility to be accelerated up to the energy range above GZK cutoff. Note that accelerating particles up to these energies is a challenge even for the most energetic astrophysical objects known. It predicts also the light relic neutrinos. This is not a final report on a closed subject, but it is hoped that suggested novel view point will serve as useful introduction and that it will thereby add the knowledge on the role of EHE C.R.s as a signature of existence of protomatter sources in the Universe. Of course, much remains to be done before one can determine whether this approach can ever contribute to the larger goal of gaining new insight into the AGN and EHE C.R. physics.

§2. Primer on GGP

This section contains some preliminaries on the GGP, the field equations of curvature (gravitation) and the inner distortion (ID) of space-time, as well the laws governing a phase transition of individual particle in the most simple case of one-dimensional space-like ID. We often suppress the indices without notice. The relevant steps of motivation are as follows:
Starting with an extension of the Minkowski space

$$\underset{x}{M_4} \to M_8 = \underset{x}{M_4} \oplus \underset{u}{M_4}, \tag{2.1}$$

in order to introduce the particle mass operator defined on the internal space $\underset{u}{M_4}$ of the inner degrees of freedom, we perform a following two-steps passage for each sample of the M_4:
a) A passage $M_4 \to M_6$ restores the complete equivalence between the three spatial and three time components

$$e_4 = (\boldsymbol{e}, e_0) \;\to\; \boldsymbol{e}_6 = (\boldsymbol{e}, \boldsymbol{e}_0) \in M_6. \tag{2.2}$$

b) A rotation $M_6 \overset{45^0}{\to} G_6$ of the basis vectors on the angle 45^0 provides an adequate algebra for quantization of the geometry [3]

$$\boldsymbol{e}_6 \overset{45^0}{\to} e_{(\lambda\alpha)}, \quad \lambda = \pm,\, \alpha = 1, 2, 3,$$

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$$e_{\pm\alpha} = \frac{1}{\sqrt{2}}(e_{0\alpha} \pm e_\alpha) = O_\pm \otimes \sigma_\alpha,$$
$$< O_\lambda, O_\tau >= {}^*\delta \equiv 1 - \delta_{\lambda\tau}, \quad < \sigma_\alpha, \sigma_\beta >= \delta_{\alpha\beta}. \tag{2.3}$$

Accordingly one gets $M_8 \to G_{12}$.

• Thus, within a simplified scheme we deal in terms of smooth differentiable manifold

$$G_{12} = \underset{\eta}{G}{}_6 \oplus \underset{u}{G}{}_6,$$
$$Dim\, G_{12} = 12, \quad Dim\, \underset{i}{G}{}_6 = 6 \quad (i = \eta, u). \tag{2.4}$$

Presumably one is allowed to perceive directly only the $\underset{\eta}{G}{}_6$ related to the space-time continuum, but not the $\underset{u}{G}{}_6$ displayed as a space of inner degrees of freedom. Let

$$e_{(\lambda,\mu,\alpha)} = O_{\lambda,\mu} \otimes \sigma_\alpha \quad (\lambda, \mu = 1, 2;\ \alpha = 1, 2, 3) \tag{2.5}$$

are the 12 basis vectors at the point p of G_{12}:

$$< O_{\lambda,\mu}, O_{\tau,\nu} >= {}^*\delta_{\lambda,\tau}{}^*\delta_{\mu,\nu}, \quad O_{\lambda,\mu} = O_\lambda \otimes O_\mu,$$
$$O_{\lambda,\mu} \leftrightarrow {}^*\mathbf{R}^4 = {}^*\mathbf{R}^2 \otimes {}^*\mathbf{R}^2, \quad \sigma_\alpha \leftrightarrow \mathbf{R}^3. \tag{2.6}$$

The metric on G_{12} is a

$$\hat{\mathbf{g}} : \mathbf{T}_p \otimes \mathbf{T}_p \to C^\infty(G_{12}) \tag{2.7}$$

section of conjugate vector bundle $S^2\mathbf{T}$. The decompositions

$$\underset{\eta}{G}{}_6 = \underset{x}{\mathbf{R}}{}^3 \oplus \underset{x_0}{\mathbf{R}}{}^3, \quad \underset{u}{G}{}_6 = \underset{u}{\mathbf{R}}{}^3 \oplus \underset{u_0}{\mathbf{R}}{}^3 \tag{2.8}$$

hold, when one has:

$$sgn(\underset{x}{\mathbf{R}}{}^3) = (+++), \quad sgn(\underset{x_0}{\mathbf{R}}{}^3) = (---), \tag{2.9}$$

and converse signatures for $\underset{u}{G}{}_6$. As far as all directions in $\underset{x_0}{\mathbf{R}}{}^3$ are equivalent in the case of flat manifold, then under a notion *time* one may imply the projection of time-coordinate on fixed arbitrary universal direction in $\underset{x_0}{\mathbf{R}}{}^3$. By this reduction

$$\underset{\eta}{G}{}_6 \to M_4 = \underset{x}{\mathbf{R}}{}^3 \oplus \underset{x_0}{\mathbf{R}}{}^1, \tag{2.10}$$

which clearly respects the physical ground, a passage to Minkowski space can be performed whenever it will be necessary. For a passage from six dimensional curved manifold $\widetilde{G}_6$ to four dimensional Riemannian geometry R_4 we reduce the three time components $\widetilde{e}_{0\alpha} = \frac{1}{\sqrt{2}}(\widetilde{e}_{(+\alpha)} + \widetilde{e}_{(-\alpha)})$ of basis six-vector $\widetilde{e}_{(\lambda\alpha)}$ at given point to the single one $\widetilde{e}_0$ in given universal direction, which merely fix a time coordinate. Actually, since Lagrangian of the fields defined on $\widetilde{G}_6$ is a function of the scalars as $A_{(\lambda\alpha)}B^{(\lambda\alpha)} = A_{0\alpha}B^{0\alpha} + A_\alpha B^\alpha$, then, taking into account that

$$A_{0\alpha}B^{0\alpha} = A_{0\alpha} < \widetilde{e}^{0\alpha}, \widetilde{e}^{0\beta} > B_{0\beta} =$$
$$A_0 < \widetilde{e}^0, \widetilde{e}^0 > B_0 = A_0 B^0, \tag{2.11}$$

one readily may perform a required passage. In this case one has

$$d\,\zeta^2 = d\,\eta^2 - d\,u^2 = 0,$$
$$d\,\eta^2\big|_{6\to4} \equiv d\,s^2 = g_{\mu\nu} d\,x^\mu d\,x^\nu = d\,u^2 = inv. \tag{2.12}$$

• A general distortion of $G_{12} \to \widetilde{G}_{12}$ under hidden Abelian local group $U^{loc}(1) = SO^{loc}(2)$ and one-dimensional trivial algebra $\hat{g} = R^1$ is considered within GGP .[2)] By the reflection of fields and their dynamics from Minkowski space to Riemannian a standard gauge principle of local internal symmetries is generalized, which is of the most importance for gravitational theories wherein the well-known difficulties stem from the fact that Riemannian geometry, in general, has not admitted a group of isometries. Namely, the Poincare transformations no longer act as isometries, and, for example, it is impossible to define energy-momentum as Noether currents related to exact symmetries. Effecting a reconciliation it is the guiding line framing GGP to explore a number of fascinating features of generating the gauge group of gravitation by hidden local internal symmetry. The most promising aspect in this approach so far is the fact that the energy-momentum conservation laws of gravitational interacting fields are formulated quite naturally by exploiting whole advantages of auxiliary flat space. Below we briefly reflect upon a few relevant points: Suppose a massless gauge field of distortion $a(\zeta)\,(\equiv a_{(\lambda,\mu,\alpha)}(\zeta))$ with the values in Lie algebra of $U^{loc}(1)$ is a local form of expression of connection in principle bundle $p : E \to G_{12}$ with a structure group $U^{loc}(1)$, where the coordinates ζ exist in the whole region $\mathcal{U}$ of flat space G_{12}. A collection of matter fields are defined as the sections of vector bundles associated with $U^{loc}(1)$ by reflection $\Phi : G_{12} \to E$ such that $p\,\Phi(\zeta) = \zeta$. They take values in standard fiber $\mathcal{F}_\zeta$ upon $\zeta : p^{-1}(\mathcal{U}) = \mathcal{U} \times \mathcal{F}_\zeta$, where an expansion into direct product $p^{-1}(\mathcal{U}) = \mathcal{U} \times G_{12}$ is defined upon $\mathcal{U}$. The various suffixes of Φ are left implicit. The fiber is Hilbert vector space on which a linear representation $U(\zeta) = \exp(-i\theta(\zeta))$ of group $U^{loc}(1)$ is given. This space regarded as Lie algebra of group $U^{loc}(1)$ upon which Lie algebra acts according to law of adjoint representation: $a \leftrightarrow ad\,a\ \Phi \to [a\,, \Phi]$. A distortion of geometry is implemented as follows:

• The 12-dimensional basis e transforms

$$\widetilde{e} = D(a)\, e \tag{2.13}$$

under the $a(\zeta)$, where the matrix $D(a)$ reads $D(a) = C(a) \otimes R(a)$, provided with the linear distortion transformations

$$\widetilde{O} = C(a)\, O, \quad \widetilde{\sigma} = R(a)\, \sigma. \tag{2.14}$$

The matrices $C(a)$ generate the group of distortion transformations of bi-pseudo-vectors:

$$C^{\tau,\nu}_{(\lambda\mu\alpha)} = \delta^\tau_\lambda \delta^\nu_\mu + \kappa a_{(\lambda,\mu,\alpha)}{}^*\delta^\tau_\lambda{}^*\delta^\nu_\mu, \tag{2.15}$$

but $R(a) \in SO(3)_{\lambda\mu}$ -the group of ordinary rotations of the planes involving two arbitrary basis of the spaces $R^3_{\lambda\mu}$ around the orthogonal third axes. The angles of rotations are determined from the special constraint imposed upon distortion transformations that a sum of distortions of corresponding basis vectors $O_{\tau,\nu}$ and σ_β must be zero for given τ,ν,β.

• We construct a diffeomorphism:

$$\widetilde{\zeta}(\zeta) : G_{12} \to \widetilde{G}_{12}, \tag{2.16}$$

where the holonomic functions $\widetilde{\zeta}(\zeta)$ imply

$$\widetilde{e}_A(a)\,\psi^A_B = e_B + \chi_B(a), \tag{2.17}$$

$A = (\lambda,\mu,\alpha)$ etc., and

$$\chi_B(a) = \widetilde{e}_A\chi^A_C = -\frac{1}{2}\widetilde{e}_A\int_0^{\zeta}(\partial_C D^A_B - \partial_B D^A_C)\,d\zeta^C, \tag{2.18}$$

which provides the criteria of integration $\partial_{\widetilde{C}}\psi^A_C = \partial_B\psi^A_{\widetilde{C}}$ and undegeneration $\|\psi\| \neq 0$ ($\psi^A_C \equiv \partial_C\widetilde{\zeta}^A$). So, out of the set of arbitrary curvilinear coordinates in $\widetilde{G}_{12}$ the real-curvilinear coordinates are distinguished, which satisfy eq. (2.17) under all the Lorentz (Λ) and gauge transformations. There is a single-valued conformity between corresponding tensors with various suffixes on $\widetilde{G}_{12}$ and G_{12} . Each transformation of real-curvilinear coordinates $\widetilde{\zeta} \to \widetilde{\zeta}'$ was generated by some Lorentz or gauge transformations. There would then exist preferred systems and group of transformations of real-curvilinear coordinates in $\widetilde{G}_{12}$. The wider group of transformations of arbitrary curvilinear coordinates in $\widetilde{G}_{12}$ would then be of no importance for the dynamics. If an inverse function ψ^B_A meets condition

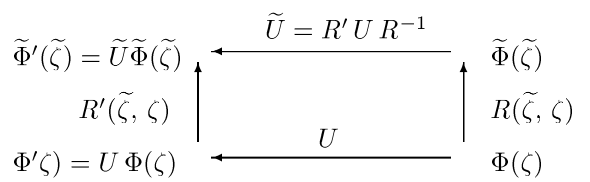
$$\frac{\partial\psi^B_A}{d\widetilde{\zeta}^C} \neq \Gamma^D_{A\,C}\,\psi^B_D, \tag{2.19}$$

where $\Gamma^D_{A\,C}$ is the usual Christoffel symbol agreed with a metric $g_{A\,B}$, then a curvature of $\widetilde{G}_{12}$ does not vanish.

• In pursuing the original problem further we are led to the principle point of drastic change of standard gauge scheme to construct a formalism of unitary reflection of the fields and their dynamics from G_{12} to $\widetilde{G}_{12}$ and vice versa. Namely, a single-valued, smooth, double-sided reflection of fields

$$R(a) : \Phi \to \widetilde{\Phi}(\Phi) : \quad (\mathcal{F}_\zeta \to \widetilde{\mathcal{F}}_{\widetilde{\zeta}}) \tag{2.20}$$

is assumed to hold where $\Phi \subset \mathcal{F}_\zeta$, $\widetilde{\Phi} \subset \widetilde{\mathcal{F}}_{\widetilde{\zeta}}$, $\widetilde{\mathcal{F}}_{\widetilde{\zeta}}$ is the fiber upon $\widetilde{\zeta} : p^{-1}(\widetilde{\mathcal{U}}) = \widetilde{\mathcal{U}} \times \widetilde{\mathcal{F}}_{\widetilde{\zeta}}$, $\widetilde{\mathcal{U}}$ is the region of base $\widetilde{G}_{12}$, $R(a)$ is the reflection matrix:



A scheme of general gauge principle.

The idea of GGP is framed into requirement of invariance of physical system of fields $\widetilde{\Phi}(\widetilde{\zeta})$ under the finite local gauge transformations $\widetilde{U} = R'UR^+$ of the Lie G_R group of gravitation generated by $U^{loc}(1)$ where $R' = R(a')$ if $a'(\zeta)$ is the gauge field transformed under $U^{loc}(1)$ in standard form. While the corresponding transformations of fields $\widetilde{\Phi}(\widetilde{\zeta})$ and their covariant derivatives are written

$$\begin{aligned} \widetilde{\Phi}'(\widetilde{\zeta}) &= \widetilde{U}(\widetilde{\zeta})\,\widetilde{\Phi}(\widetilde{\zeta}), \\ \left[g^A(\widetilde{\zeta})\,\widetilde{\nabla}_A\widetilde{\Phi}(\widetilde{\zeta})\right]' &= \widetilde{U}(\widetilde{\zeta})\left[g^A(\widetilde{\zeta})\,\widetilde{\nabla}_A\widetilde{\Phi}(\widetilde{\zeta})\right], \end{aligned} \tag{2.21}$$

where $\widetilde{U}^+\widetilde{U} = 1$, $\widetilde{\nabla}_A = \partial_A + \Gamma_A$, provided by the connection

$$\Gamma_A(\zeta) = \frac{1}{2}\,\Sigma^{C\,D}\,V^B_C(\zeta)\,\partial_A V_{D\,B}(\zeta), \tag{2.22}$$

the $\Sigma^{C\,D}$ are generators of Lorentz group, $V^A_C(\zeta)$ are the components of affine tetrad vectors in used coordinate net ζ^A. One has $g^A(\widetilde{\zeta}) \Rightarrow \widetilde{e}^A(\widetilde{\zeta})$ for fields of spin $(s = 0, 1)$; but $g^A(\widetilde{\zeta}) = V^A_B(x)\gamma^B$ for spinor field $(s = \frac{1}{2})$, where γ^B are the modified Dirac's matrices defined on G_{12}. Since the fields $\Phi(\zeta)$ no longer hold, the reflected ones $\widetilde{\Phi}(\widetilde{\zeta})$ will be regarded as the real physical fields. But a conformity $R(a) : \Phi \to \widetilde{\Phi}(\Phi)$ and eq. (2.21) enables $\Phi(\zeta)$ to serve as an auxiliary fields on the flat space G_{12}. These notions arise basically from the most important fact that a Lagrangian $\widetilde{L}(\widetilde{\zeta})$ of fields $\widetilde{\Phi}(\widetilde{\zeta})$ may be obtained under the reflection from a Lagrangian $L(\zeta)$ of corresponding $\Phi(\zeta)$ fields and vice versa. A Lagrangian $\widetilde{L}(\widetilde{\zeta})$ is as well an invariant under the wider group of arbitrary curvilinear transformations $\widetilde{\zeta} \to \widetilde{\zeta}'$:

$$J_\psi\widetilde{L}(\widetilde{\zeta})\Big|_{inv(G_R;\,\widetilde{\zeta}\to\widetilde{\zeta}')} = L(\zeta)|_{inv(\Lambda;\,U^{loc}(1))}\,, \tag{2.23}$$

where $J_\psi \equiv \|\psi\|\sqrt{g}$, g is the determinant of metric tensor on $\widetilde{G}_{12}$. A local internal gauge $U^{loc}(1)$ symmetry remains hidden symmetry as far it is screened by gauge group of gravitation G_R. Field equations are derived from an invariant action

$$S = S_a + \widetilde{S}_{\widetilde{\Phi}}, \tag{2.24}$$

where the S_a action of gauge field of distortion a defined on G_{12} is invariant under the Lorentz and $U^{loc}(1)$ gauge groups, but the $\widetilde{S}_{\widetilde{\Phi}}$ action of the rest of fields defined on $\widetilde{G}_{12}$ is invariant under the gauge group of gravitation G_R. In the sequel, a total action S is as well $U^{loc}(1)$ gauge invariant. Field equations followed at once in terms of Euler-Lagrange variations respectively on the G_{12} and $\widetilde{G}_{12}$. While the equation of distortion field a_A reads

$$\partial^B\partial_B a_A - (1-\zeta_0^{-1})\partial_A\partial^B a_B = J_A = -\frac{1}{2}\sqrt{g}\frac{\partial g^{BC}}{\partial a_A}T_{BC}, \tag{2.25}$$

where T_{BC} is the energy-momentum tensor, ζ_0 is the gauge fixing parameter. The curvature of manifold $\underset{\eta}{G} \to \underset{\eta}{\widetilde{G}}$ is a familiar distortion when

$$\begin{aligned} a_{(1,1,\alpha)} = a_{(2,1,\alpha)} &\equiv \frac{1}{\sqrt{2}}\,\underset{\eta}{a}_{\,(+\alpha)}, \\ a_{(1,2,\alpha)} = a_{(2,2,\alpha)} &\equiv \frac{1}{\sqrt{2}}\,\underset{\eta}{a}_{\,(-\alpha)}. \end{aligned} \tag{2.26}$$

The other regime of inner distortion (ID) presents at:

$$a_{(1,1,\alpha)} = -a_{(2,1,\alpha)} \equiv \frac{1}{\sqrt{2}} \underset{u}{a}_{(+\alpha)},$$
$$a_{(1,2,\alpha)} = -a_{(2,2,\alpha)} \equiv \frac{1}{\sqrt{2}} \underset{u}{a}_{(-\alpha)}. \quad (2.27)$$

A matter found in the ID-region of space-time continuum has undergone phase transition of II-kind, i.e., each particle goes off from the mass shell, a shift of mass and energy-momentum spectra occurs upwards along the energy scale. The matter in this state is called *protomatter* with the thermodynamics differed strongly from the thermodynamics of ordinary compressed matter (see below). We are interested in the case of a spherical-symmetric gravitational field $a_0(r)$ in presence of one-dimensional space-like ID-field $\widetilde{a}$, thus,

$$a_{(1,1,3)} = a_{(2,2,3)} = a_{(+3)} = \frac{1}{2}\left(-a_0 + \widetilde{a}\right),$$
$$a_{(1,2,3)} = a_{(2,1,3)} = a_{(-3)} = \frac{1}{2}\left(-a_0 - \widetilde{a}\right),$$
$$a_{(\lambda,\mu,1)} = a_{(\lambda,\mu,2)} = 0, \quad \lambda,\mu = 1,2. \quad (2.28)$$

Passing back from the $\underset{\eta}{\widetilde{G}_6}$ to R_4, it has the group of motions $SO(3)$ with two-dimensional space-like orbits S^2 where the standard coordinates are θ and φ. The stationary subgroup of $SO(3)$ acts isotropically upon the tangent space at the point of sphere S^2 of radius $\widetilde{r}$. So, the bundle $p : R_4 \to \widetilde{R}^2$ has the fiber $S^2 = p^{-1}(\widetilde{x})$, $\widetilde{x} \in R_4$ with a trivial connection on it, where $\widetilde{R}^2$ is the quotient-space $R_4/SO(3)$. One can then easily determine the basis vectors $\underset{i}{\widetilde{e}}$, where $tan\theta_{(\pm 3)} = \kappa\left(-a_0 \pm \widetilde{a}\right)$. They read

$$\widetilde{e}_0 = e_0(1-x_0) + \underset{u}{e}_3\, x, \quad \widetilde{e}_3 = e_3(1+x_0) - \underset{u}{e}_{03}\, x,$$
$$\widetilde{e}_1 = \frac{1}{2}\{(\cos\theta_{(+3)} + \cos\theta_{(-3)})e_1 + (\sin\theta_{(+3)} + \sin\theta_{(-3)})e_2 + (\cos\theta_{(+3)} - \cos\theta_{(-3)})\,\underset{u}{e}_{01} + (\sin\theta_{(+3)} - \sin\theta_{(-3)})\,\underset{u}{e}_{02}\},$$
$$\widetilde{e}_2 = \frac{1}{2}\{(\cos\theta_{(+3)} + \cos\theta_{(-3)})e_2 - (\sin\theta_{(+3)} + \sin\theta_{(-3)})e_1 + (\cos\theta_{(+3)} - \cos\theta_{(-3)})\,\underset{u}{e}_{02} - (\sin\theta_{(+3)} - \sin\theta_{(-3)})\,\underset{u}{e}_{01}\},$$
$$(2.29)$$

where $x_0 \equiv \kappa a_0, \quad x \equiv \kappa\widetilde{a}$. We assume an absence of transversal stresses and the transference of masses in R_4

$$T_1^1 = T_2^2 = T_3^3 = -\widetilde{P}(\widetilde{r}),$$
$$T_0^0 = -\widetilde{\rho}(\widetilde{r}), \quad \widetilde{r} \in \widetilde{R}^3 \quad (R_4 = \widetilde{R}^3 \oplus \widetilde{R}^0),$$
$$(2.30)$$

where $\widetilde{P}(\widetilde{r})$ and $\widetilde{\rho}(\widetilde{r})$ are taken to denote the internal pressure and macroscopic density of energy defined in proper frame of reference that is being used. In the case at hand, one has in Feynman gauge:

$$\Delta a_0 = \frac{1}{2}\left\{g_{00}\frac{\partial g^{00}}{\partial a_0}\widetilde{\rho}(\widetilde{r}) - \left[g_{33}\frac{\partial g^{33}}{\partial a_0} + g_{11}\frac{\partial g^{11}}{\partial a_0} + g_{22}\frac{\partial g^{22}}{\partial a_0}\right]\widetilde{P}(\widetilde{r})\right\},$$
$$\left(\Delta - \lambda_a^{-2}\right)\widetilde{a} = \frac{1}{2}\left\{g_{00}\frac{\partial g^{00}}{\partial \widetilde{a}}\widetilde{\rho}(\widetilde{r}) - \left[g_{33}\frac{\partial g^{33}}{\partial \widetilde{a}} + g_{11}\frac{\partial g^{11}}{\partial \widetilde{a}} + g_{22}\frac{\partial g^{22}}{\partial \widetilde{a}}\right]\widetilde{P}(\widetilde{r})\right\}\theta\left(\lambda_a - n^{-\frac{1}{3}}\right),$$
$$(2.31)$$

where n is the concentration of particles, $\lambda_a = \frac{\hbar}{m_a c} \simeq 0.4fm$ is the Compton lenght of the ID-field, but the results are, however, rather insensitive to its formal value since the real ID-effects occur far below it, a diffeomorphism $\widetilde{r}(r) : M_4 \to R_4$ is given by $r = \widetilde{r} - R_g/4$ where R_g is the gravitational radius.

• A distortion of the basis $\widetilde{e}$ in the ID regime, in turn, yields the transformations of Poincaré generators of translations. Given an explicit form eq. (2.29) of distorted basis vectors it is straightforward to derive the laws of phase transition for individual particle found in the ID-region $(x_0 = 0,\ x \neq 0)$ of the space-time continuum $\tan\widetilde{\theta}_3 = -x, \quad \widetilde{\theta}_1 = \widetilde{\theta}_1 = 0$. Accordingly, the Poincaré generators P_μ of translations undergone the transformations:

$$\widetilde{E} = E, \quad \widetilde{P}_{1,2} = P_{1,2}\cos\widetilde{\theta}_3, \quad \widetilde{P}_3 = P_3 - \tan\widetilde{\theta}_3 mc, \quad \widetilde{m} = \left|\left(m - \tan\widetilde{\theta}_3\,\frac{P_3}{c^2}\right)^2 + \sin^2\widetilde{\theta}_3\,\frac{P_1^2+P_2^2}{c^2} - \tan^2\widetilde{\theta}_3\,\frac{E^2}{c^4}\right|^{\frac{1}{2}}, \quad (2.32)$$

where $E, \boldsymbol{P}, m$ and $\widetilde{E}, \widetilde{\boldsymbol{P}}, \widetilde{m}$ are ordinary and distorted energy, momentum and mass at rest, $\widetilde{\theta}_3$ is the angle of space-like ID. The metric in holonomic coordinate basis takes the form

$$g_{00} = (1-x_0)^2 + x^2, \quad g_{\mu\nu} = 0 \quad (\mu \neq \nu),$$
$$g_{33} = -\left[(1+x_0)^2 + x^2\right], \quad g_{11} = -\widetilde{r}^2,$$
$$g_{22} = -\widetilde{r}^2\sin^2\theta. \quad (2.33)$$

We also bring an explicit form of the line element from the outside of configuration $\widetilde{r} > \widetilde{r}_b$, where $\widetilde{r}_b$ is the boundary of distribution of matter with the total mass M. In this case the eq. (2.31) in Feynman gauge reduced to $\nabla^2 a_0 = 0$, which has the solution $\frac{\kappa}{\sqrt{2}}a_0 = -\frac{GM}{r} = -\frac{r_g}{2r}$. Hence the basis vectors read

$$\widetilde{e}_0 = D_0^0 e_0, \quad \widetilde{e}_r = D_r^r e_r, \quad \widetilde{e}_\theta = e_\theta, \quad \widetilde{e}_\varphi = e_\varphi, \quad (2.34)$$

provided by $(a_0 \equiv a_{0r})$

$$D_0^0 = 1 + \frac{x_0}{\sqrt{2}}, \quad D_r^r = 1 - \frac{x_0}{\sqrt{2}},$$
$$e_0 = \xi_0 \otimes \sigma_r, \quad e_r = \xi \otimes \sigma_r, \quad e_\theta = \xi \otimes \sigma_\theta, \quad e_\varphi = \xi \otimes \sigma_\varphi,$$

$$\xi_0 = \frac{1}{\sqrt{2}}(O_+ + O_-), \quad \xi = \frac{1}{\sqrt{2}}(O_+ - O_-). \tag{2.35}$$

The coordinates $\widetilde{x}^\mu(\widetilde{t}, \widetilde{r}, \theta, \varphi)$ implying a diffeomorphism

$$\widetilde{x}^\mu(x)^l : M_4 \to R_4 \tag{2.36}$$

exist in the whole region $p^{-1}(\mathcal{U}) \in R_4$, where $\widetilde{x}^{0r} \equiv \widetilde{t}, \quad \widetilde{x}^{0\theta} = \widetilde{x}^{0\varphi} = 0$ and

$$\frac{\partial \widetilde{x}^\mu}{\partial x^l} \equiv \psi^\mu_l = \frac{1}{2}\left(D^\mu_l + \omega^m_l D^\mu_m\right). \tag{2.37}$$

The ω^m_l can be derived from the equation

$$\frac{\partial D^\mu_k}{\partial x^l} = \frac{\partial}{\partial x^k}\left(\omega^m_l D^\mu_m\right), \tag{2.38}$$

where, in general, $\omega = \omega^l_l$ is the Lorentz scalar function of the trace of curvature form Ω of connection a. That is

$$\omega = \omega(tr\,\Omega) = \omega(d\,tr\,\boldsymbol{a}), \tag{2.39}$$

provided by

$$\Omega = \sum_{l<k} F_{lk}\, dx^l \wedge dx^k, \quad \boldsymbol{a} = a_l\, dx^l, \tag{2.40}$$

where $F_{lk} = \partial_l\, a_k - \partial_k\, a_l$. Whence, a functional ω has a null variation derivative $\frac{\delta\omega(tr\,\Omega)}{\delta\, a} = 0$ at local variations of connection $a_l \to a_l + \delta a_l$, namely $\omega = inv(\Lambda,\, U^{loc}(1))$. A straightforward calculations give non-vanishing components

$$\psi^0_0 = \frac{1}{2}D^0_0, \quad \psi^0_1 = \frac{1}{2}\,t\,\partial_r D^0_0, \quad \psi^1_1 = D^r_r,$$
$$\psi^2_2 = \psi^3_3 = 0, \quad \omega^1_1 = \omega^2_2 = \omega^3_3 = 1, \quad \omega^0_1 = t\,\partial_r D^0_0. \tag{2.41}$$

where the corresponding components of Christoffel symbol are written $\Gamma^0_{01} = \frac{1}{2}g^{00}\partial_{\widetilde{r}}g_{00}$, $\Gamma^1_{00} = -\frac{1}{2}g^{11}\partial_{\widetilde{r}}g_{00}$, $\Gamma^1_{11} = \frac{1}{2}g^{11}\partial_{\widetilde{r}}g_{11}$ such that

$$\frac{\partial \psi^l_\mu}{\partial \widetilde{x}^\nu} \neq \Gamma^\lambda_{\mu\nu}\,\psi^l_\lambda, \tag{2.42}$$

namely a curvature of R_4 is not zero. Then, the line element is written down

$$d\,s^2 = (1 - x_0)^2 d\widetilde{t}^2 - (1 + x_0)^2 d\widetilde{r}^2 - \widetilde{r}^2(\sin^2\theta d\,\varphi^2 + d\,\theta^2), \tag{2.43}$$

§3. GGP for any spin

To facilitate writing in [2] the GGP was utilized only for the most important fields of spin 0, $\frac{1}{2}$, 1. But developed formalism may be readily extended to the field of arbitrary spin s, since latter will be treated as a system of $2s$ fermions of half-integer spin. To render our discussion more transparent, below we consider this case in detail in terms of four-dimensional space-time continuum. Certainly, the whole scheme outlined in section 2 will then hold provided we simply replace the 12-dimensional indices $A = (\lambda, \mu, \alpha)$ by the four-dimensional ones. Avoiding the inconsistencies arisen in standard approach to higher-spin field equations in curved space, below we proceed by adopting the most convenient Bargman-Wigner's wave functions [4] incorporating the results of section 2. The wave function of spin-$j = \frac{n}{2}$ particle with momentum $\boldsymbol{P}$ can be obtained by Lorentz transformation from the symmetric Dirac spinor of rank n corresponding to the particle in the rest $U_{\alpha_1\ldots\alpha_n}(0)$ implying $(\gamma_4 - 1)^{\beta'}_\beta U_{\beta'\beta_2\ldots\beta_n}(0)$ for each index

$$U_{\beta_1\ldots\beta_n}(\boldsymbol{P}) = S^{\beta'_1}_{\beta_1}(\alpha(\boldsymbol{P}))\ldots S^{\beta'_n}_{\beta_n}(\alpha(\boldsymbol{P}))U_{\beta'_1\ldots\beta'_n}(0), \tag{3.1}$$

while

$$\bar{U}'(\boldsymbol{P})U'(\boldsymbol{P}) = \bar{U}(\Lambda^{-1}\boldsymbol{P})U(\Lambda^{-1}\boldsymbol{P}). \tag{3.2}$$

A spin part is written $\Sigma_{\mu\nu} = \frac{1}{2}\sum_{r=1}^n \Sigma^{(r)}_{\mu\nu}$, where a matrix $\Sigma^{(r)}_{\mu\nu}$ acts only on the r-th index

$$\left(\Sigma^{(r)}_{\mu\nu}\right)^{\beta_1\ldots\beta_n}_{\alpha_1\ldots\alpha_n} = \delta^{\beta_1}_{\alpha_1}\cdots\delta^{\beta_{r-1}}_{\alpha_{r-1}}\left(\Sigma^{(r)}_{\mu\nu}\right)^{\beta_r}_{\alpha_r}\delta^{\beta_{r+1}}_{\alpha_{r+1}}\cdots\delta^{\beta_n}_{\alpha_n}, \tag{3.3}$$

provided that

$$\Sigma^{(r)}_{\mu\nu} = \frac{1}{4}[\gamma^r_\mu,\, \gamma^r_\nu],$$
$$\left(\gamma^r_\mu\right)^{\beta_1\ldots\beta_n}_{\alpha_1\ldots\alpha_n} = \delta^{\beta_1}_{\alpha_1}\cdots\delta^{\beta_{r-1}}_{\alpha_{r-1}}\left(\gamma^r_\mu\right)^{\beta_r}_{\alpha_r}\delta^{\beta_{r+1}}_{\alpha_{r+1}}\cdots\delta^{\beta_n}_{\alpha_n}. \tag{3.4}$$

Following to the GGP scheme let the spin-j field $\Phi_{\beta_1\ldots\beta_n}$ corresponding to the spinor eq. (3.1) is a section of vector bundle associated with group $U^{loc}(1)$ by reflection $\Phi : M_4 \to E$, the $R^{(r)}$ reflection matrix is defined for the r-th index (see eq. (3.9)), and the spin-j field $\widetilde{\Phi}_{\beta_1\ldots\beta_n}(\widetilde{x})$ takes values in the fiber $F_{\widetilde{x}}$ $(\widetilde{x} \in \widetilde{\mathcal{U}}$, $\widetilde{\mathcal{U}}$ is the region of the base $R_4)$, such that its Lagrangian will be invariant under a local gauge transformations

$$\widetilde{\Phi}'(\widetilde{x}) = \widetilde{U}^{(r)}\,\widetilde{\Phi}(\widetilde{x}),$$
$$\left(g^\mu_{(r)}(\widetilde{x})\,\nabla^{(r)}_\mu\widetilde{\Phi}(\widetilde{x})\right)' = \widetilde{U}^{(r)}\left(g^\mu_{(r)}(\widetilde{x})\,\nabla^{(r)}_\mu\widetilde{\Phi}(\widetilde{x})\right). \tag{3.5}$$

Provided that $g^\mu_{(r)}(\widetilde{x}) = V^\mu_\alpha(\widetilde{x})\gamma^\alpha_{(r)}$, $\nabla^{(r)}_\mu$ is the covariant derivative on R_4 defined only for the r-th index by the standard substitution [5]

$$\nabla^{(r)}_\mu \widetilde{U}_{\beta_1\ldots\beta_n} \to$$
$$\Lambda^{\alpha'}_\alpha(\widetilde{x})S^{\beta'_1}_{\beta_1}(\alpha(\Lambda))\ldots S^{\beta'_n}_{\beta_n}(\alpha(\Lambda))\nabla^{(r)}_\mu\widetilde{U}_{\beta_1\ldots\beta_n}, \tag{3.6}$$

where

$$\nabla^{(r)}_\alpha = V^\mu_\alpha\,(\partial_\mu + \Gamma^{(r)}_\mu),$$
$$\Gamma^{(r)}_\mu = \frac{1}{2}\,\Sigma^{\alpha\beta}_{(r)}\,V^\nu_\alpha(\widetilde{x})\,\partial_\mu V_{\beta\nu}(\widetilde{x}),\ \Gamma^{(r)}_\mu(\widetilde{x}) =$$
$$\frac{1}{4}\Delta^{(r)}_{\mu,\alpha\beta}\,\gamma^\alpha_{(r)}\,\gamma^\beta_{(r)}, \tag{3.7}$$

$\Delta^{(r)}_{\mu,\alpha\beta}(x)$ are the Ricci rotation coefficients. The eq. (3.5) holds if

$$\widetilde{\Phi}^{(r)} = R^{(r)}\,\Phi, \quad g^\mu_{(r)}\,\nabla^{(r)}_\mu\widetilde{\Phi}^{(r)} = R^{(r)}\,S^{(r)}\,(\gamma\,D\,\Phi). \tag{3.8}$$

where $D_l = \partial_l - ig\, a_l(x)$, $S(a)$ is the gauge invariant Lorentz scalar (see below). The reflection matrix $R^{(r)}$ is written

$$R^{(r)}(\widetilde{x},\, x) = R_f(x)\, R_g^{(r)}(\widetilde{x}) =$$
$$\exp\left(-i\Theta(x) - \widetilde{\Theta}^{(r)}(\widetilde{x})\right), \quad \Theta(x) = g\int_0^x a_l(x)dx^l,$$
$$\widetilde{\Theta}^{(r)}(\widetilde{x}) = \frac{1}{2}\int_0^{\widetilde{x}} R_f^+\left\{g^\mu_{(r)}\,\Gamma^{(r)}_\mu\, R_f,\, g^{(r)}_\nu\, d\widetilde{x}^\nu\right\}, \tag{3.9}$$

and

$$S(a) = \frac{1}{8K^{(r)}}\psi^l_\mu\left\{\widetilde{R^{(r)}}^+\, g^\mu_{(r)}\, R^{(r)},\, \gamma^{(r)}_l\right\} = inv, \tag{3.10}$$

where

$$K^{(r)} = \widetilde{R^{(r)}}^+ R^{(r)} = \widetilde{R^{(r)}}^+_g R^{(r)}_g =$$
$$\exp\left\{-\frac{1}{2}\int_0^{\widetilde{x}}\left[\left(R_f^+\,\widetilde{\Gamma^{(r)}}^+_\mu\, g^\mu_{(r)},\, g^{(r)}_\nu\, d\widetilde{x}^\nu\right) R_f\; + \right.\right.$$
$$\left.\left. R_f^+\left(g^\mu_{(r)}\,\Gamma^{(r)}_\mu\, R_f,\, g^{(r)}_\nu\, d\widetilde{x}^\nu\right)\right]\right\},$$
$$\widetilde{\Gamma^{(r)}}^+_\mu \equiv \gamma^0\,\Gamma^{(r)+}_\mu\,\gamma^0, \quad \widetilde{R}^{(r)} \equiv \gamma^0\, R^{(r)}\,\gamma^0. \tag{3.11}$$

Taking into account that $[R_f,\, g^{(r)}_\nu] = 0$, also [6]

$$\widetilde{\Gamma^{(r)}}^+_\mu\, g^\nu_{(r)} + g^\nu_{(r)}\,\Gamma^{(r)}_\mu = -\nabla^{(r)}_\mu\, g^\nu_{(r)} = 0 \tag{3.12}$$

we get $K^r = 1$, and

$$\widetilde{U^{(r)}}^+ U^{(r)} = \gamma^0 U^{(r)+}\gamma^0 U^{(r)} = 1. \tag{3.13}$$

A Lagrangian of the spin-j field can be written

$$L(x) = J_\psi\widetilde{L}(\widetilde{x}) = J_\psi\left\{\frac{i}{2}\left[\bar{\widetilde{\Phi}}^{(r)}(\widetilde{x})\, g^\mu_{(r)}(\widetilde{x})\nabla^{(r)}_\mu\widetilde{\Phi}^{(r)}(\widetilde{x}) - \right.\right.$$
$$\left.\left.(\nabla^{(r)}_\mu\bar{\widetilde{\Phi}}^{(r)}(\widetilde{x}))\, g^\mu_{(r)}(\widetilde{x})\widetilde{\Phi}^{(r)}(\widetilde{x})\right] - m\bar{\widetilde{\Phi}}^{(r)}(\widetilde{x})\widetilde{\Phi}^{(r)}(\widetilde{x})\right\} =$$
$$J_\psi\left\{S^{(r)}(a)\frac{i}{2}\left[\bar{\Phi}\,\gamma^l_{(r)} D_l\,\Phi - (D_l\bar{\Phi})\,\gamma^l_{(r)}\Phi\right] - m\bar{\Phi}\,\Phi\right\}. \tag{3.14}$$

Generalized Bargman-Wigner's equation for the spin-$j = \frac{n}{2}$ particle in curved space stems from a Lagrangian eq. (3.14):

$$\left(g'^\mu\,\nabla'_\mu\widetilde{\Phi}^{(r)} - m\right)^{\beta'}_\beta\,\widetilde{\Phi}_{\beta'\beta_2\ldots\beta_n}(\widetilde{\boldsymbol{P}}) =$$
$$[R'\,(S'\,D - m)]^{\beta'}_\beta\;\Phi_{\beta'\beta_2\ldots\beta_n}(\boldsymbol{P}) = 0, \tag{3.15}$$

where $R', S', \ldots$ refer to index β'.

§4. Brief outline of some issues in microscopic approach to the BH physics

In this section we explore the ID regime of space-time continuum to study equilibrium fermion configurations and generalized models of AGNs. The equations describing the equilibrium configurations include the gravitational and ID field equations eq. (2.31), the hydrostatic equilibrium equation

$$\frac{\partial\widetilde{P}}{\partial\widetilde{r}} + (\widetilde{P} + \widetilde{\rho})F = 0, \quad F = g^{00}\frac{\partial g^{00}}{\partial\widetilde{r}}, \tag{4.1}$$

and the state equation

$$\widetilde{\rho} = \widetilde{\rho}(\widetilde{P}) \tag{4.2}$$

specified for each domain of many layered configurations. Given the state equation eq. (4.2), the eq. (4.1) can be integrated, where the integration constant is determined from the condition of matching of internal and external metrics. Hence

$$g_{00}(r_f) = (1 - \frac{R_g}{2r^b})^2\exp\left[\int_0^{\widetilde{P}}\frac{2\widetilde{P}}{\widetilde{P} + \widetilde{\rho}}\right], \tag{4.3}$$

where $r^b = r(\widetilde{r}_b)$. We consider the equilibrium configurations of two most general classes with spherical-symmetric distribution of matter in many-phase stratified states .[7,8] The I-class configurations include:

1_I.*Domain* $\rho < \rho_{drip}$ - the shell made of cold catalized matter, which formed after nuclear burning in the density range below neutron drip $\rho < \rho_{drip} = 4.3 \times 10^{11}\, g\, cm^{-3}$. It consists of surface $\rho \le 10^6\, g\, cm^{-3}$, where the temperature and magnetic fields strongly affect the equation of state and outer crust ($\rho_{drip} \le \rho < 4.54 \times 10^{12}\, g\, cm^{-3}$)- a Coulomb lattice of heavy nuclei co-exist in β- equilibrium with relativistic electrons.

2_I. *Domain* $\rho_{drip} \le \rho < 4.54 \times 10^{12}\, g\, cm^{-3}$ inner crust-the electrons, nuclei and free neutrons co-exist in the medium; For a domain $7.86\, g\, cm^{-3} \le \rho < 4.63 \times 10^{12}\, g\, cm^{-3}$ we use the simple semiempirical formula of state equation given by Harrison and Wheeler .[9] The nuclear matter in high density range above nucleus density still remain not so well understood. A large number of representative models for the nuclear equation of state are available in literature, but all they are subject to many uncertainties, including such exotic processes as neutron and proton superfluidity, a pion condensation, a phase transition to quark matter. For a simplicity, above the density $\rho > 4.54 \times 10^{12}\, gcm^{-3}$ the I-class configurations are now thought to be composed of two phase of ideal cold n-p-e gas, which is mixture of neutrons, protons and electrons in complete β- equilibrium. The first phase state covers the intermediate density -

3_I. *Domain* $4.54 \times 10^{12}\, gcm^{-3} \le \rho < \rho_d = 2.6 \times 10^{16}\, g\, cm^{-3}$- which is the regular n-p-e gas in absence of ID. Second phase state is-

4_I. *Domain* $\rho > \rho_d$ - the n-p-e protomatter at short nucleon-nucleon distances $r_{NN} \le 0.4 fm$ in presence of ID. This regime needs in a special theoretical study .[7,8] Up to the density range $\rho \le \rho_{fl} = 4.09 \times 10^{14}\, g\, cm^{-3}$ to which the nucloun-nucleon distances $r_{NN} \le 1.6 fm$ correspond one has the same domains for the second class configurations (II). Above the density ρ_{fl} we consider an onset of melting down of hadrons when nuclear matter consequently turns to quark matter. In the domain of $\rho_{fl} \le \rho < \rho_{as} = m_n\,(0.25 fm)^3 = 1.07 \times 10^{17}\, g\, cm^{-3}$, where m_n is the neutron mass at rest, 0.25 fm is the string thickness, we consider two phase states of string flip-flop regimes:

4_{II}. *Domain* $\rho_{fl} \leq \rho < \rho_d$, to which the distances $0.46fm < r_{NN} \leq 1.6fm$ correspond- the regular string flip-flop when ID is absent, which is a kind of tunneling effect when the strings joining the quarks stretch themselves violating energy conservation and after touching each other they switch on to the other configuration;[10)]

5_{II}. *Domain* $\rho_d \leq \rho < \rho_{as}$- the string flip-flop regime in presence of ID at distances $0.25fm < r_{NN} \leq 0.4fm$ - a system is made of quark protomatter in complete β-equilibrium with rearrangement of string connections joining them.

6_{II}. *Domain* $\rho > \rho_{as}$- the system is made of quarks in one beg in complete β-equilibrium at presence of ID under the weak interactions and gluons, including the effects of QCD-perturbative interactions. Due to the screening of strong forces, the quarks are considered to be free inside the bag and to interact only in the surface region. The surface energy is estimated to be proportional to quark density. The quark protomatter is in overall color singlet ground state, which is a non-interacting relativistic Fermi gas found in the ID-region of the space-time continuum. A theoretical treatment of the last two domains is given in .[7,8)] Below we bring a brief discussion of properties of quark protomatter only in two domains, respectively, of asymptotic freedom and flip-flop. The QCD vacuum has a complicated structure, which is intimately connected to the glueon-glueon interaction. The confinement of quarks is a natural feature of the exercising a pressure B on the surface of the local region of the perturbative vacuum to which quarks are confined. This is just the main idea of phenomenological MIT quark bag model ,[11)] where quarks are assumed to be confined in a bag. The stability of the hadron is ensured by the vacuum pressure B and surface tension. The surface energy is estimated to be proportional to quark density. In most applications, sufficient accuracy is obtained by assuming that all the quarks are almost massless inside a bag. Now, our purpose is to prompt this physical picture with appropriate modifications into the framework describing a medium of quark protomatter regarded as non-interacting Fermi gas found in the ID-region of the space-time continuum, when $r_{NN} \leq 0.25fm$. We shall consider the case of quark protomatter of u, d and s flavors in complete β-equilibrium, namely

$$d \to u + l + \bar{\nu}, \quad s \to u + l + \bar{\nu}, \tag{4.4}$$

where l denotes electron (e^-) or muon (μ^-), $(\bar{\nu})$ is the associated antineutrino. Assuming β-equilibrium and that the neutrinos not to be retained in the medium, we get for the generalized chemical potentials $\widetilde{\mu} = \left[(\widetilde{P}_F c)^2 + (\widetilde{m}c^2)^2\right]^{\frac{1}{2}}$

$$\widetilde{\mu}_d \to \widetilde{\mu}_u + l + \widetilde{\mu}_l, \quad \widetilde{\mu}_s \to \widetilde{\mu}_u + \widetilde{\mu}_l, \tag{4.5}$$

and $\widetilde{\mu}_\nu = \mu_{nu} = 0$. For an extremely relativistic degenerate Fermi gas $\widetilde{\nu}_i \propto g_i \widetilde{\mu}_i$, where $g_i = 2$ for $i = l$, and $g_i = 6$ for each kind of quark $i = u, d, s$ in two spin and three color states. These equations lead to $\widetilde{\mu}_d = \widetilde{\mu}_s$ and $\widetilde{\nu}_d = \widetilde{\nu}_s$ when equilibrium between μ^- and e^- yields $\widetilde{\mu}_\mu = \widetilde{\mu}_e \equiv \widetilde{\nu}_l$ and $\widetilde{\nu}_\mu = \widetilde{\nu}_e \equiv \widetilde{\nu}_l$. The charge neutrality condition gives $\frac{2}{3}\widetilde{\nu}_u - \frac{1}{3}\widetilde{\nu}_d - \frac{1}{3}\widetilde{\nu}_s - \widetilde{\nu}_\mu - \widetilde{\nu}_e = 0$, or $\widetilde{\nu}_u - \widetilde{\nu}_s - 3\widetilde{\nu}_e = 0$. Whence $\left(\frac{\widetilde{\nu}_s}{6}\right)^{\frac{1}{3}} = \left(\frac{\widetilde{\nu}_u}{6}\right)^{\frac{1}{3}} + \left(\frac{\widetilde{\nu}_e}{6}\right)^{\frac{1}{3}}$. Then in terms of new variables $Z_e = \frac{\widetilde{\nu}_e}{\widetilde{\nu}_s}$ and $Z_u = \frac{\widetilde{\nu}_u}{\widetilde{\nu}_s}$ one gets

$$Z_u - 1 = 3Z_e, \quad 1 = Z_u^{\frac{1}{3}} + (3Z_e)^{\frac{1}{3}}, \tag{4.6}$$

which has a solution $Z_u = 1$ and $Z_e = 0$. Asymptotically $(\widetilde{P}_F >> \widetilde{m}c)$ one has

$$\widetilde{\nu}_u = \widetilde{\nu}_s = \widetilde{\nu}_d \equiv \widetilde{\nu}_b, \quad \widetilde{\nu}_e = \widetilde{\nu}_\mu = 0,$$
$$\widetilde{P}_{Fu} = \widetilde{P}_{Fs} = \widetilde{P}_{Fd} \equiv \widetilde{P}_{Fb}, \tag{4.7}$$

i.e., no leptons are presented in quark protomatter in β-equilibrium. Such a system is Λ-like protomatter with energy density $\widetilde{E}_0 = \frac{3}{4}(\widetilde{\nu}_u\widetilde{P}_{Fu} + \widetilde{\nu}_d\widetilde{P}_{Fd} + \widetilde{\nu}_s\widetilde{P}_{Fs})$. Now, we will get an overview of the QCD interaction effects in sufficient approximation with extension to quark protomatter some of those results and calculation have done in the case of ordinary quark matter. The first effect is the shift of the vacuum energy per unit volume. The bag constant $B \simeq 55\, MeV\, fm^{-3}$ of the MIT bag model [11)] must be added to the kinetic energy density. Including the gluon exchange perturbative interactions the energy density of quark protomatter is then given by the non-interacting Fermi contribution plus bag constant

$$\widetilde{E}_0 = \sum_i \frac{3}{4}\, \widetilde{\nu}_i\, \widetilde{P}_{Fi}\, \widetilde{b}_I\, (\widetilde{N},\, \alpha_c) + B, \tag{4.8}$$

where $\widetilde{P}_{Fi}$ is the distorted Fermi momentum of i flavor, $\widetilde{N}$ is the number of flavors present. The $\widetilde{N}$ and running coupling constant α_c takes into account the QCD perturbative interactions. The first correction to the free ground state is the ordinary exchange energy corresponding to the second order closed loop diagrams. [12)] Next correction is coming up from the sum of different ring diagrams. With equal numbers of quarks of each flavor presented, according to [13)] the modified function $\widetilde{b}_I(\widetilde{N}, \alpha_c)$ is written

$$\widetilde{b}_I(\widetilde{N}, \alpha_c) = \widetilde{N}[1 + \frac{2\alpha_c}{3\pi} +$$
$$\frac{\alpha_c^2}{3\pi^2}(\widetilde{N} \ln \frac{\alpha_c \widetilde{N}}{\pi} + 0.02\widetilde{N} + 6.75)]. \tag{4.9}$$

For numerical calculations it is sufficient to make use of the value $\alpha_c \simeq 2.2$ fitting the MIT bag model, where the quarks will be taken fully relativistic $m_i \to 0$.

In [10)] the calculations of flip-flop processes in dense matter have been carried out without QCD-corrections and the effects of pions. A similar calculations are made in [7,8)] but for the quark protomatter, wherein we have concerned with the individual particle approximation (Hartree approximation), while the Hartree potential is almost linearly proportional to the string length. The Y shape string is the most convenient for calculations, because the center of it almost equals to the center of gravity. We employ the tunneling effect of quantum fluctuation of string and the negative potential energy caused by such a quantum jump. The basic technique adopted for calculation of transition matrix element $\widetilde{K}$ is the instanton technique (semi-classical treatment). Due to quantum string flip-flop, an attractive interaction between quarks is presented, when during the quantum transition from a state ψ_1 of energy $\widetilde{E}_1$ to another one ψ_2 of energy $\widetilde{E}_2$ the lowering of energy of system takes place. The quark matter acquires $\Delta\widetilde{E}$ correction to the classical string energy such that the flip-flop energy lowers the energy of quark matter, consequently by lowering the critical density or critical Fermi momentum. If one, for example, looks for the string flip-flop transition amplitude of simple system of $q\bar{q}q\bar{q}$ described by the Hamiltonian $\widetilde{H}$ and invariant action $\widetilde{S}$, then one has

$<$  $\Big| \, e^{-\widetilde{H}T} \, \Big|$  $> = < \int [d\,\widetilde{\sigma}] \; e^{-\widetilde{S}} >, \qquad (4.10)$

where T is a (imaginary) time interval, $[d\,\widetilde{\sigma}]$ is the integration over all the possible string motion. The action $\widetilde{S}$ is proportional to the area $\widetilde{A}$ of the surface swept by the strings in the finite region of ID-region of R_4. The strings are initially in the  configuration and finally in the  configuration. Note that the maximal contribution to the path integral eq. (4.10) comes from the surface σ_0 of the minimum surface area ("instanton"). A computation of the transition amplitude is strightforward by summing over all the small vibrations around σ_0. Hence

$$\widetilde{K} \simeq \pi\sqrt{a_0}\, e^{-a_0 \Delta\widetilde{A}},$$

where $\Delta\widetilde{A}$ is the increase of the 1-instanton surface area over the 0-instanton one, which actually equals to the area the string sweeps in making the transition. Note that string has a finite thickness d, and the width of the area $\Delta\widetilde{A}$ cannot be less than d. This cutoff introduces a factor $\exp(-a_0\, d\, r_{NN})$, (where r_{NN} is the distance between two separated centers) in the amplitude $\widetilde{K}$ resulting in the finite-ranged potential. The interaction energy between two centers has a range of order $2\widetilde{\bar{r}}$ due to overlap of wave functions. For $r_{NN} < \widetilde{\bar{r}}$ the interaction potential reads

$$\widetilde{V_{cc}} = \begin{cases} -\sqrt{a_0} & r_{NN} < \widetilde{\bar{r}}, \\ 0 & \text{otherwise}, \end{cases} \qquad (4.11)$$

where $\widetilde{r}_i$ is the distance from the center, a_0 is the string tension estimated from charmonium spectroscopy $a_0 \simeq 0.1\, GeV$. One simplifies the calculations by assuming that the centers are uniformly distributed with a concentration $\widetilde{n}_b$. Hence, the average of distance to the nearest center equals to

$$\widetilde{\bar{r}} = \Gamma(\frac{4}{3}) \left(\frac{4\pi}{3}\widetilde{n}_b\right)^{-\frac{1}{3}}, \qquad (4.12)$$

where Γ is the gamma function. This leads to expression of the interaction energy per center (or per baryon)

$$\widetilde{U}(\widetilde{n}_b) = \frac{1}{2}\,\widetilde{n}_b \int \widetilde{V_{cc}}\, d^3 r_{NN} \simeq -\frac{1}{2}\,\widetilde{n}_b\,\frac{4\pi}{3}\widetilde{\bar{r}}^3. \qquad (4.13)$$

The potential energy, as usual, is of the Wigner type and caused a surface tension to the quark protomatter: Surface tension $\simeq \frac{\pi}{8}\widetilde{n}_b^2\,\widetilde{\bar{r}}^4\,\sqrt{a_0}$. However, the quark flip-flop regime and bag model are only phenomenological, and it remains to see in future whether or not such quark systems really could exist.

4.1 State equations

The hydrostatic equilibrium equation eq. (4.1) can be rewritten in the terms of a new variable ν

$$n_{OV} = 7.9 \times 10^{55} e^{\nu}, \quad \rho_{OV} = 7.2 \times 10^{-4} e^{\nu}. \qquad (4.14)$$

It reads

$$\nu' = -(s_1 + s_2)\frac{1}{2}\left(\ln g_{00}\right)', \qquad (4.15)$$

provided the $(')$ means $\frac{\partial}{\partial r}$, $s_1 = \frac{\widetilde{P}_{OV}\,\nu'}{P'_{OV}}$ and $s_2 = \frac{\widetilde{\rho}_{OV}\,\nu'}{\rho'_{OV}}$. The state equations are specified below for each domain step-by-step away from the domain of lower density up to the domain of higher density .[7,8] For the I-class configurations one has respectively-

domain: $-27.2 \le \nu < -15.5$

$$P_{OV} = 4.7 \times 10^{-25}\left(1.9 \times 10^5 \rho_{OV}^{\frac{1}{3}} - 1.4\right)^5 - 2.3 \times 10^{-26},$$
$$s_1 = \frac{1.5 \times 10^{-7}\rho_{OV}^{-\frac{1}{3}}\left[\left(1.9 \times 10^5 \rho_{OV}^{\frac{1}{3}} - 1.4\right)^5 - 1\right]}{\left(1.9 \times 10^5 \rho_{OV}^{\frac{1}{3}} - 1.4\right)^4},$$
$$s_2 = 6.6 \times 10^{18}\rho_{OV}^{\frac{2}{3}}\left(1.9 \times 10^5 \rho_{OV}^{\frac{1}{3}} - 1.4\right)^{-4}; \qquad (4.16)$$

domain: $-15.5 \le \nu < -2.8$:

$$P_{OV} = 0.03\rho_{OV}^{\frac{5}{4}}\left(1 + 2.8 \times 10^{-5}\rho_{OV}^{-\frac{1}{2}}\right)^{-\frac{5}{6}},$$
$$s_1 = \frac{0.08\left(1 + 2.8 \times 10^{-5}\rho_{OV}^{-\frac{1}{2}}\right)}{\left(1 + 3.9 \times 10^{-5}\rho_{OV}^{-\frac{1}{2}}\right)},$$
$$s_2 = 3.2\rho_{OV}^{-\frac{1}{4}}\left(1 + 2.8 \times 10^{-5}\rho_{OV}^{-\frac{1}{2}}\right); \qquad (4.17)$$

domain: $-2.8 \le \nu < -0.11$:

$$P_{OV} = 1.8 \times 10^{-5}\rho_{OV}^{\frac{2}{3}}\left(1 + 1.4\rho_{OV}^{1/6})^6\right),$$
$$s_1 = \frac{1.5\left(1 + 1.4\rho_{OV}^{\frac{1}{6}}\right)}{\left(1 + 3.5\rho_{OV}^{\frac{1}{6}}\right)},$$
$$s_2 = 8.4 \times 10^4\rho_{OV}^{\frac{1}{3}}\left(1 + 1.4\rho_{OV}^{\frac{1}{6}}\right)^{-5}\left(1 + 3.5\rho_{OV}^{\frac{1}{6}}\right)^{-1}; \qquad (4.18)$$

domain $\nu \ge -0.11$

$$\widetilde{\rho} = \frac{\widetilde{m}_e c^2 \chi(\widetilde{y}_e)}{\widetilde{\lambda}_e^3} + \frac{\widetilde{m}_p c \chi(\widetilde{y}_p)}{\widetilde{\lambda}_p^3} + \frac{\widetilde{m}_n c^2 \chi(\widetilde{y}_n))}{\widetilde{\lambda}_n^3},$$
$$\widetilde{P} = \frac{\widetilde{m_e} c^2 \varphi(\widetilde{y}_e)}{\widetilde{\lambda}_e^3} + \frac{\widetilde{m_p} c^2 \varphi(\widetilde{y}_p)}{\widetilde{\lambda}_p^3} + \frac{\widetilde{m_n} c^2 \varphi(\widetilde{y}_n)}{\widetilde{\lambda}_n^3}. \qquad (4.19)$$

A following notational conventions will be used throughout:

$$\chi(\widetilde{y}) = \frac{1}{8\pi^2}\left\{\widetilde{y}(1+\widetilde{y}^2)^{\frac{1}{2}}(1+2\widetilde{y}^2) - \ln\left[\widetilde{y}+(1+(1+\widetilde{y}^2)^{\frac{1}{2}}\right]\right\},$$

$$\varphi(\widetilde{y}) = \frac{1}{8\pi^2}\left\{\widetilde{y}(1+\widetilde{y}^2)^{\frac{1}{2}}(\frac{2}{3}\widetilde{y}^2-1) + \ln\left[\widetilde{y}+(1+(1+\widetilde{y}^2)^{\frac{1}{2}}\right]\right\},$$

$$\widetilde{m} = \mid \eta \mid^{\frac{1}{2}}, \quad \eta = 1 - x^2 - \frac{xy}{\sqrt{3}} - \frac{y^2x^4}{6(1+x^2)},$$

$$y = \frac{P_F}{mc} = (3\pi^2)^{\frac{1}{3}}\lambda n^{\frac{1}{3}}, \quad \widetilde{y} = \frac{\widetilde{P}_F}{\widetilde{m}c} = (3\pi^2)^{\frac{1}{3}}\widetilde{\lambda}\widetilde{n}^{\frac{1}{3}},$$

$$\widetilde{P_F} = P_F\zeta^{\frac{1}{2}}, \quad \zeta = y^2\left[1 - \frac{2x^2}{3(1+x^2)}\right] + \frac{2xy}{\sqrt{3}} + x^2, \tag{4.20}$$

$\lambda = \hbar/mc$, $\widetilde{\lambda} = \hbar/\widetilde{m}c$, $\widetilde{P}_F$ and P_F are the distorted and ordinary Fermi momenta, $\widetilde{n}$ is the distorted concentration of particles. Further simplification gives

$$\widetilde{y}_p = \left(\frac{\widetilde{a}_1 + \widetilde{a}_2\widetilde{y}_n^2 + \widetilde{a}_3\widetilde{y}_n^4}{1+\widetilde{y}_n^2}\right)^{1/2}, \tag{4.21}$$

where

$$\widetilde{a}_1 = (1/4)\left[\left(\widetilde{Q}/\widetilde{m}_p\right)^2 - \left(\widetilde{m}_e/\widetilde{m}_p\right)^2\right] \times \left[(1+\widetilde{m}_p/\widetilde{m}_n)^2 - (\widetilde{m}_e/\widetilde{m}_n)^2\right],$$

$$\widetilde{a}_2 = (1/2)\left[(\widetilde{m}_n/\widetilde{m}_p)^2 - 1 - (\widetilde{m}_e/\widetilde{m}_p)^2\right],$$

$$\widetilde{a}_3 = (\widetilde{m}_n/2\widetilde{m}_p)^2, \quad \widetilde{Q} = \widetilde{m}_n - \widetilde{m}_p. \tag{4.22}$$

The ratio of proton-neutron distorted concentrations takes form

$$\frac{\widetilde{n}_p}{\widetilde{n}_n} = \left(\frac{\widetilde{m}_p}{\widetilde{m}_p\widetilde{y}_n}\right)^3 \left[\left(\widetilde{a}_1 + \widetilde{a}_2\widetilde{y}_n^2 + \widetilde{a}_3\widetilde{y}_n^4\right)\left(1+\widetilde{y}_n^2\right)\right]^{\frac{3}{2}}. \tag{4.23}$$

For the intermediate density domain of regular n-p-e gas in absence of ID, the proton-neutron ratio initially decreases, as the density increases, and reaches a maximum value of 0.0026 at $\rho_0 \simeq 7.8\times10^{11}\, g\, cm^{-3}$, and afterwards rises monotonically to $\frac{1}{8}$ for high densities .[9)]

In the case of the II-class configurations one has: domain- $8.5 \leq \nu < 9.9$- the potential energy per quark reads:

$$\widetilde{V}_{\text{per q}} = a_0\widetilde{\bar{r}}\beta = \widetilde{M}(\widetilde{P}_{Fb})\beta, \tag{4.24}$$

where

$$\widetilde{M}(\widetilde{P}_{Fb}) = 1.4\,\frac{a_0}{c^3\,\widetilde{P}_{Fb}} = \frac{0.2}{y_b}\,m_n, \tag{4.25}$$

or

$$\widetilde{M}(\widetilde{P}_{Fb}) = \frac{M(P_{Fb})}{\widetilde{\zeta}_b^2},$$

$$\widetilde{\zeta}_b = \frac{\zeta_b}{y_b^2} = 1 - \frac{2x^2}{3(1+x^2)} + \frac{2x}{\sqrt{3}y_b} + \frac{x^2}{y_b^2},$$

$$y_b = \frac{P_{F\,b}}{m_nc} = \frac{(3\pi^2)^{\frac{1}{3}}\hbar n_b^{\frac{1}{3}}}{m_nc}, \tag{4.26}$$

m_n is the neutron mass at rest, $\widetilde{M}(\widetilde{P}_{Fb})$ is the effective distorted mass of the quark. It will be convenient to make numerical calculations in the approximation that quarks have small ordinary mass $m_i \simeq m_u = 5MeV$: $\widetilde{m} = \widetilde{m}_u\,\widetilde{\eta}_u^{\frac{1}{2}}$, where

$$\widetilde{\eta}_u = \left(\mid \eta_u \mid^{1/2} + 5441.6/\zeta_u^{\frac{1}{2}}\right)^2,$$

$$\eta_u = 1 - x^2 - \frac{xy_u}{\sqrt{3}} - \frac{y_u^2x^4}{6(1+x^2)},$$

$$\zeta_u = y_u^2\left[1 - \frac{2x^2}{3(1+x^2)}\right] + \frac{2xy_u}{\sqrt{3}} + x^2. \tag{4.27}$$

The density and internal pressure are as follows:

$$\widetilde{\rho} = 3\,\widetilde{m}c^2\,\frac{\chi(\widetilde{y})}{\widetilde{\lambda}^3} - 56.3MeV\,\widetilde{n}_b, \quad \widetilde{P} = 3\,\widetilde{m}c^2\,\frac{\varphi(\widetilde{y})}{\widetilde{\lambda}^3}, \tag{4.28}$$

Finally, for domain $\nu \geq 9.9$, one has

$$\widetilde{\rho} = \sum_i \widetilde{m}c_i^2\,\frac{\chi(\widetilde{y}_i)\widetilde{b}_I}{\widetilde{\lambda}_i^3} + B, \quad \widetilde{P} = \sum_i \widetilde{m_i}c^2\,\frac{\varphi(\widetilde{y}_i)\widetilde{b}_I}{\widetilde{\lambda}_i^3}, \tag{4.29}$$

where the quarks will be taken fully relativistic $m_i \to 0$, i is the flavor $i = u, d, s$.

4.2 Neutron stars

In this case the ID-field equals zero $x \equiv 0$. The remaining ones are integrated numerically leading from the center of configuration up to the surface, where the internal pressure tends to zero. Each configuration is defined by the unique free parameter of the central value of particle concentration $n(0)$ or equivalent to it central density $\rho(0)$. The interior gravitational potential $x_0^{int}(r)$ should be matched into the exterior one $x_0^{ext}(r)$. Checking it out one introduces the dimensionless sewing parameter

$$D(r_b) = \frac{\left|x_0^{int}(r_b) - x_0^{ext}(r_b)\right|}{x_0^{int}(r_b)}, \tag{4.30}$$

when one must choose two unknown constants to satisfy the subsidiary sewing condition $D(r_b) = 0$ imposed at the boundary of configuration. The results of the numerical integration are as follows: The maximum values of the masses

$$M = \frac{4\pi}{c^2}\int_0^R \widetilde{\rho}r^2\,dr, \quad M_1 = \frac{4\pi}{c^2}\int_V \widetilde{\rho}\,dV,$$

$$M_0 = Nm_n = \int_V m_nN\,dV, \tag{4.31}$$

are

$$\frac{M}{M_\odot} \simeq 1.1, \quad \frac{M_1}{M_\odot} \simeq 1.23, \quad \frac{M_0}{M_\odot} \simeq 1.07,$$

where $dV = 4\pi(-g_{33})^{1/2}r^2\,dr$ is the volume element. Note that the macroscopic energy includes apart from the energy at rest also the energy of motion of particles and the energy of their interaction per $1cm^{-3}$. That is why the total mass M is not equal to the sum of masses of volume elements M_1, and since $(-g_{33})^{1/2} > 1$, then $M < M_1$. The other mass M_0 is calculated without taking into account the binding energy of gravitational interaction between the particles. The possible maximal total masses and radii, for example, of the II-class configurations are

$$\frac{M}{M_\odot} \simeq 0.9 \div 1.1, \quad R \simeq (8.1 \div 9.5)km,$$

while the stable configurations of $x_0(0) = 0.01$ correspond to

$$\frac{\rho(0)}{\rho_0} \simeq 4.1 \div 15.8, \quad \frac{P(0)}{P_0} \simeq 6.4 \times 10^5 \div 2.7 \times 10^6,$$

where $\rho(0)$ and $P(0)$ are central values of density and pressure, $\rho_0 = 2.8 \times 10^{14}\, g\, cm^{-3}$ and $P_0 = 10^{33}\, erg\, cm^{-3}$, $N_0 = 10^{57}$. The Fig. 1 corresponds to the II-class configurations with $x_0(0) = 0.01$ where one makes use of special units $P_{OV} = 6.469 \times 10^{36}\, erg\, cm^{-3}$, $\rho_{OV} = 7.194 \times 10^{15}\, g\, cm^{-3}$ and $r_{OV} = 13.68$. Reflecting upon the results far obtained in this subsection we draw a statement that discussed gravitational theory is consistent with the general relativity up to the limit of neutron stars. It is remarkable that the inferred from the astrophysical observations values have been in a variaty of ways, are well within the range of theoretically allowed masses of neutron star models discussed here.

4.3 AGNs

In presence of the ID-mechanism each configuration now will be defined by two free parameters of central values of particle concentration $\widetilde{n}(0)$ and ID-field $x(0)$. The central value of the gravitational potential $x_0(0)$ can be found by reiterating integrations when the sewing condition holds. An elaborated scenario is as follows:

• The SPC with thermodynamical properties differed strongly from the thermodynamics of ordinary matter are formed, where, in PC, the energy density and internal pressure have sharply increased proportional to gravitational forces of compression in about 18-20 order of magnitude with respect to corresponding central values of neutron star. This counteracts the collapse and equilibrium holds even for the maximum masses $M_{max} \simeq 3.48 \times 10^8 M_\odot$.

• The SPCs are found inside the even horizon (EH) sphere and could be observed only in presence of accreting matter in their close vicinities. Versus the central values of the parameters of SPCs surrounded by the ADs such configurations are considered as the AGNs.

• A singularity of metric no longer holds because of action of the MSC effect. It can be significant for the AGNs to which the accreting matter has steadily filled the inside of EH sphere and in due course has formed a protomatter disk (PD) around the PC, which gradually becomes very thin at reaching out the edge of the EH. In the small region of intersection of PD with the EH the ID-field is switched on and a metric singularity disappeared (eq.(2.33)). Along with it a sharply increase of gravitational forces has ceased, and the particles may

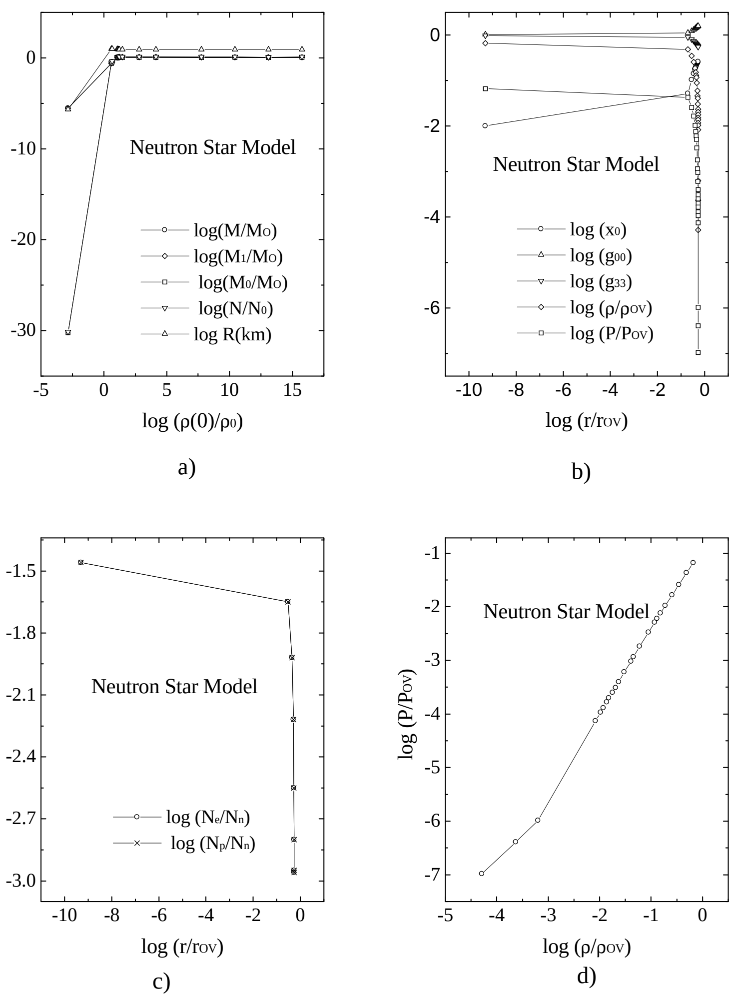


Fig. 1. *a) The masses, number of baryons and radius against the central density; b) the radial profiles of x_0, pressure and density; c) the radial profiles of the electron-neutron and proton neutron ratios; d) the state equation.*

escape, in principle, through this vista to the outside world.

As an aid to clarify the problem, some detailed numerical results are tabulated. The equilibrium configurations could exist only if a sewing condition $D << 1$ holds. We also present the Fig. 2 of the I-class configurations with

$$\frac{M}{M_\odot} \simeq 1.1\times10^7 \div 5.3\times10^7, R \simeq (8.6\times10^6 \div 3.9\times10^7)\, km,$$

where the stable configurations correspond to

$$\frac{\rho(0)}{\rho_0} \simeq 0.01 \div 118.0, \quad \frac{P(0)}{P_0} \simeq 756.8 \div 7.5 \times 10^8,$$

and the Fig. 3 of the II-class configurations of

$$\frac{M}{M_\odot} \simeq 5.9\times10^7 \div 3.3\times10^8, \; R \simeq (4.4\times10^7 \div 2.4\times10^8)\, km,$$

where for the stable configurations one has

$$\frac{\rho(0)}{\rho_0} \simeq 6.6\times10^5 \div 1.2\times10^6, \; q\frac{P(0)}{P_0} \simeq 1.1\times10^{13} \div 1.5\times10^{13}.$$

Table I. I-class configurations.
Some integral characteristics of the AGN models.

$\max(\frac{M}{M_\odot})$	$R(km)$	$R_g(km)$	$\frac{N}{N_0}$	D
4.8×10^{2}	4.3×10^{2}	1.4×10^{3}	4.4×10^{3}	1.4×10^{-2}
1.4×10^{3}	1.2×10^{3}	4.2×10^{3}	1.8×10^{3}	4.3×10^{-2}
1.6×10^{4}	1.3×10^{4}	4.8×10^{4}	4.1×10^{5}	5.1×10^{-2}
2.1×10^{5}	1.6×10^{5}	6.1×10^{5}	1.1×10^{7}	2.3×10^{-2}
1.4×10^{6}	1.1×10^{6}	4.2×10^{6}	1.4×10^{8}	1.8×10^{-2}
7.3×10^{6}	5.5×10^{6}	2.2×10^{7}	1.3×10^{9}	6.3×10^{-2}
1.1×10^{7}	8.5×10^{6}	3.4×10^{7}	2.3×10^{9}	2.8×10^{-2}
7.7×10^{7}	5.7×10^{7}	2.3×10^{8}	2.9×10^{10}	9.2×10^{-2}
9.7×10^{7}	7.2×10^{7}	2.9×10^{8}	3.9×10^{10}	2.1×10^{-2}
1.2×10^{8}	8.8×10^{7}	3.5×10^{8}	5.3×10^{10}	3.0×10^{-2}
2.3×10^{8}	1.7×10^{8}	6.8×10^{8}	1.3×10^{11}	1.3×10^{-2}
2.7×10^{8}	2.0×10^{8}	8.0×10^{8}	1.6×10^{11}	8.1×10^{-2}
3.3×10^{8}	2.5×10^{8}	9.9×10^{8}	2.1×10^{11}	1.5×10^{-2}
3.4×10^{8}	2.5×10^{8}	1.0×10^{9}	2.1×10^{11}	9.3×10^{-3}
3.4×10^{8}	2.5×10^{8}	1.0×10^{9}	2.2×10^{11}	6.5×10^{-3}
3.4×10^{8}	2.6×10^{8}	1.0×10^{9}	2.2×10^{11}	5.5×10^{-3}
3.5×10^{8}	2.6×10^{8}	1×10^{9}	2.1×10^{11}	4.9×10^{-3}
3.5×10^{8}	2.6×10^{8}	1.0×10^{9}	2.2×10^{11}	1.5×10^{-3}

Table II. II-class configurations.
Some integral characteristics of the AGN models.

$\max(\frac{M}{M_\odot})$	$R(km)$	$R_g(km)$	$\frac{N}{N_0}$	D
4.7×10^{2}	4.7×10^{2}	1.4×10^{3}	5.7×10^{3}	3.8×10^{-2}
1.8×10^{3}	1.6×10^{3}	5.5×10^{3}	3.1×10^{4}	2.0×10^{-2}
2.8×10^{4}	2.2×10^{4}	8.2×10^{4}	9.2×10^{5}	5.5×10^{-2}
3.7×10^{5}	2.8×10^{5}	1.1×10^{6}	2.7×10^{7}	3.1×10^{-2}
9.6×10^{5}	7.3×10^{5}	2.8×10^{6}	9.5×10^{7}	6.4×10^{-2}
3.3×10^{7}	2.5×10^{6}	9.8×10^{6}	4.9×10^{8}	3.9×10^{-3}
5.7×10^{6}	4.3×10^{6}	1.7×10^{7}	5.7×10^{6}	2.8×10^{-2}
1.7×10^{7}	1.2×10^{7}	4.9×10^{7}	4.2×10^{9}	3.3×10^{-2}
2.0×10^{7}	1.5×10^{7}	6×10^{7}	8.6×10^{9}	1.6×10^{-3}
4.2×10^{7}	3.1×10^{7}	1.3×10^{8}	6.6×10^{10}	1.9×10^{-2}
5×10^{7}	3.7×10^{7}	1.5×10^{8}	4.2×10^{10}	6.9×10^{-2}
5.9×10^{7}	4.4×10^{7}	1.7×10^{8}	1.4×10^{12}	1.0×10^{-3}
9.1×10^{7}	6.7×10^{7}	2.7×10^{8}	1.4×10^{12}	1.2×10^{-2}
1.4×10^{8}	1.1×10^{8}	4.2×10^{8}	1.8×10^{11}	3.0×10^{-2}
2.6×10^{8}	1.9×10^{8}	7.6×10^{8}	2.5×10^{11}	1.9×10^{-2}
3.2×10^{8}	2.3×10^{8}	9.4×10^{8}	3.5×10^{11}	1.8×10^{-2}
3.3×10^{8}	2.4×10^{8}	9.7×10^{8}	3.3×10^{11}	1.3×10^{-2}
3.4×10^{8}	2.5×10^{8}	1.0×10^{9}	8.0×10^{14}	5.1×10^{-2}

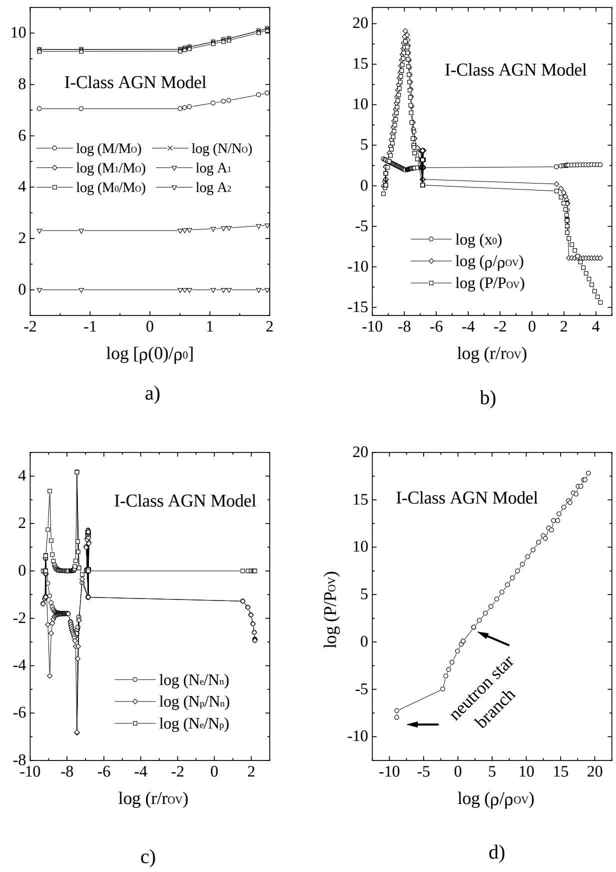


Fig. 2. a) The masses, number of baryons and gravitational packing coefficients versus central density; b) the radial profiles of x_0, pressure and density; c) the radial profiles of the electron-neutron, proton-neutron and electron-proton ratios; d) the state equation.

§5. Prediction of two crucial elements for the GZK air showers

Although more than three decades passed since the famous first detection of cosmic rays with huge energies above $10^{19}eV$ by the Volcano Ranch group led by John Linsley, nevertheless, the solution to this outstanding puzzle had not been achieved yet, and this principle problem was ever since much the same as now. At present, about 20 events above $10^{20}eV$ have been reported (e.g. [14]) worldwide by experiments such as the High Resolution Fly's Eye, AGASA, Fly's Eye, Haverah Park, Yakutsk, and Volcano Ranch. Neutrinos are perhaps the most interesting of all elementary particles, well-known to play an important role in the Universe. They can penetrate cosmological distances in the Universe and their trajectories are not deflected as they are neutral. Therefore, the most economical among the various proposals initiated by the EHE C.R. puzzle suggests that the extra-galactic EHE neutrinos may escape the GZK cutoff and travel on cosmic distances hitting local light relic neutrinos clustered in dark halos and form EHE C.R. through the hadronic Z and W- bosons decays .[15] The weakest link in this hypothesis are probably both unknown accelerating mechanism of the primary neutrinos up to this huge energies and their large flux required at the resonant energy well above the GZK cutoff. Such a flux severely challenges conventional source models. Furthermore, any concomitant photon flux must not violate existing upper limits. Note that the physics of light mass $\sim eV$ of relic neutrinos is also an open question.

5.1 EHE extra-galactic AGN-neutrino fluxies

Hereafter the quantities denoted by wiggles refer to SCP while the corresponding quantities of quark-star or

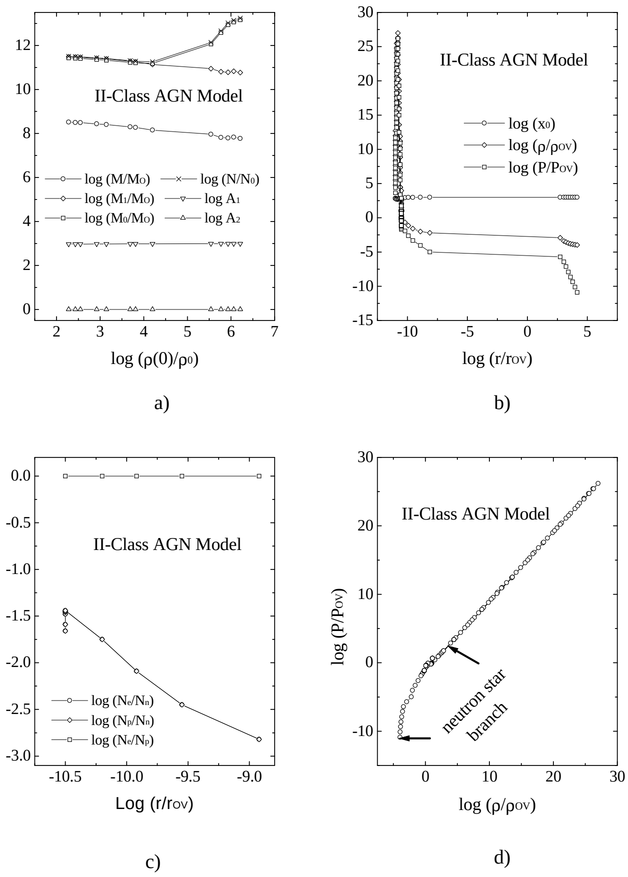
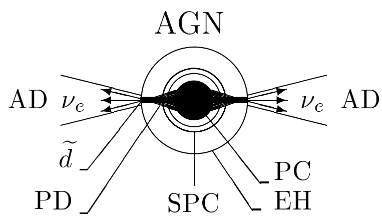


Fig. 3. a) The masses, number of baryons and gravitational packing coefficients versus central density; b) the radial profiles of x_0, pressure and density; c) the radial profiles of the electron-neutron, proton-neutron and electron-proton ratios; d) the state equation.

neutron star are left without wiggles. In this subsection we should see why the extra-galactic sources-AGNs can be strong EHE neutrino emitters. After formation of the SPC, which accommodates the highest energy scale ($> 10^{21} eV$) in the central massive PC, the radius of it equals $\sim \frac{1}{4} R_g$, thus, the SPC is always found inside the EH sphere and could be observed only in presence of AD. At this early stage the compact SPCs do work in AGNs absolutely in the same way as central massive BHs, the engines thought to be responsible for the activity of AGNs. Here, of course, a key open question is to enlighten the mechanisms that trigger the activity, and how a large amount of matter can be funneled to the central regions, to fuel this activity. In high luminosity AGNs the large-scale internal gravitational instabilities drive gas towards the nucleus that trigger big starbursts, and the coeval compact cluster just formed. It seemed they have some connection to the nuclear fueling through mass loss of young stars as well their tidal disruption, and supernovae. A rough quantitatively estimate of fueling can be given assuming that (e.g. [16]): 1) the typical luminosity and power of an active nucleus to be of the order or higher than $10^{46}\, erg\, s^{-1}$; 2) a mass-to-energy conversion efficiency $\epsilon \sim 10\%$ ($L = \frac{dM}{dt} c^2 \epsilon$). Then the mass accretion rate $\frac{dM}{dt}$ should be:

$$\frac{dM}{dt} \sim 1.7 \frac{0.1}{\epsilon} \left(\frac{L}{10^{46}\, erg\, s^{-1}} \right) \frac{M_\odot}{yr}.$$

That is, a total mass up to $2 \times 10^8\, M_\odot$ should be available inside of EH during the duty cycle of the AGN is of order of $10^8\, yr$. It must lose its large angular momentum, which can be due to viscouse torques in a geometrically thin AD (e.g. [17]). Then, the most important next stage of evolution of the SPCs succeeds, which is in strong contrast to standard BH physics. Namely, during the time $\leq 10^8\, yr$ this infalling matter has formed a PD around the PC, which gradually becomes very thin at reaching out the edge of EH. The MSC effects acts in the small region of intersection of PD with the EH, and hence, the particles may escape, in principle, through this vista to the outside world. The predominant cooling mechanism of the SPC is neutrino emission, namely the neutrino created leave the SPC carrying away energy and thus cooling of SPC. Both the quark- and the pion condensed PC cool much more rapidly than n-p-e-PC, because the simple URCA process can occur in both cases, while in the n-p-e PC it is very inefficient for phase space reasons. In the latter the nucleon-modified URCA processes can occur only when the number of participating degenerate fermions is two larger than for simple URCA processes.



EHE Neutrino Cooling of the SPC.

To determine the cooling rate of the SPC, one first need to find the surface temperature, because the temperature that determines the thermal emission from the SPC is that at the surface T_e^{SPC} rather than the interior temperature $\widetilde{T} > \frac{10^{26}\, eV}{k_B} \sim 1.2 \times 10^{30} K$ in the center of PC, or the temperature of surface of PC: $\widetilde{T}_\nu > \frac{10^{21}\, eV}{k_B} \sim 1.2 \times 10^{25} K$. The SPC interiors are to a good approximation isothermal, but near the surface the temperature drops rapidly. We can estimate the factor $\frac{T_e^{SPC}}{\widetilde{T}_\nu}$ by notifying that Thomson scattering already dominates the opacity at $T \geq 2 \times 10^8 K$. Taking the SPC surface composition to be pure $^{56}_{26}Fe$ we have that $X = 0, \quad Z = 1, \quad \mu_e = \frac{56}{26}$ and $\mu = \frac{56}{27}$, where X is the mass-fraction of hydrogen, Z is the mass-fraction of "heavy" elements. The appropriate approximation to the opacity is [9] $\kappa = \kappa_T = \frac{0.40}{\mu_e} cm^2 g^{-1}$. We apply a standard discussion of the degenerate-nondegenerate transition region of a neutron star to this region of SPC. Namely, the roughly isothermal interior of SPC is completely degenerate, while the surface layers of SPC are nondegenerate and considered to be in radiative equilibrium. Equating then the luminosity at this transition layer of PC

$$\widetilde{L} \sim \left(1.6 \times 10^{23}\, erg\, s^{-1} \mu\right) \frac{\widetilde{M}}{M_\odot} \widetilde{T}_\nu^{\frac{3}{2}} \quad (5.1)$$

to blackbody photon emission

$$\widetilde{L}_\gamma \sim \left(7 \times 10^{36}\, erg\, s^{-1}\right) \left(\frac{\widetilde{R}}{10\, km}\right)^2 \left(T_{e\,7}^{SPC}\right)^4 \quad (5.2)$$

from the surface at an effective surface temperature T_e^{SPC}, one has

$$\frac{T_e^{SPC}}{\widetilde{T}_\nu} \simeq 1 \times 10^{-2} \widetilde{T}_{\nu 9}^{-\frac{5}{8}} \left(\frac{\widetilde{M}}{M_\odot}\right)^{\frac{1}{4}} \left(\frac{\widetilde{R}}{10\, km}\right)^{-\frac{1}{2}}, \quad (5.3)$$

where $T_{e\,7}^{SPC}$ is the temperature in units of $10^7 K$. Hence,

$$\frac{T_e^{SPC}}{\widetilde{T}_\nu} \sim 7.8 \times 10^{-12}, \quad T_e^{SPC} \sim 8.8 \times 10^{13} K. \quad (5.4)$$

Given the L_ν^q luminosity of the neutron star of mass M and a uniform density ρ with no muons [18] due to the neutrino emission nucleon-modified URCA process it is straightforward to derive the same for the SPC_I. Actually, taking into account that for each degenerate species, only a fraction $\frac{\widetilde{T}}{\widetilde{T}_F}$ effectively contribute to the cooling rate, and, thus, there are two such initial species and three such final species, we can readily write down the resulting luminosity of a uniform density SPC with no muons due to the neutrino emission nucleon-modified URCA process:

$$\widetilde{L}_{\nu\,\varepsilon}^{URCA} = \varepsilon \frac{\widetilde{M}}{M_\odot} \left(\frac{\widetilde{\rho}}{\rho_{nuc}}\right)^{\frac{4}{3}} \left(\frac{T_e^{SPC}}{T_e^n}\right)^3 L_\nu^{URCA}, \quad (5.5)$$

where

$$\varepsilon = \varepsilon_d\, \varepsilon_{trap}, \quad \varepsilon_d = \frac{2\,\pi\, R_g\, \widetilde{d}}{4\,\pi\, R_g^2} = \frac{\widetilde{d}}{2R_g}, \quad (5.6)$$

$\widetilde{d}$ is the thickness of the PD at the edge of even horizon, ε_{trap} is the neutrino "trapping" coefficient, which is due to the fact that as the neutrinos are formed in PC at super-high densities they experience greater difficulty escaping from the PC before being dragged along with the matter, namely the neutrinos are "trapped" comove with matter. Neutrino trapping forces the liberated energy to be emitted on a much larger diffusion timescale for neutrinos to diffuse out of the PC, which can be estimated by assuming that the coherent scattering in PC is the dominant opacity source. Namely, the heavy particle of PC acts nonlinearly as a single particle, thus,

$$\widetilde{t}_{diff} \sim \frac{\widetilde{\lambda}^{coh}\, \widetilde{N}_{scatt}}{c},$$

where $\widetilde{\lambda}^{coh}$ is the mean free path of a typical neutrino, $\widetilde{N}_{scatt}$ is the number of scatterings experienced by the neutrino prior to escape. Coherent scattering induces a random-walk trajectory for the neutrino in the PC and PD without slightly changing its energy before reaching the surface of SPC :[19] $\widetilde{\lambda}^{coh}\, \widetilde{N}_{scatt}^{\frac{1}{2}} \sim \widetilde{R}$. Hence

$$\frac{\widetilde{t}_{diff}}{t_{diff}} = \left(\frac{\widetilde{R}}{R_{nuc}}\right)^2 \frac{\lambda^{coh}}{\widetilde{\lambda}^{coh}} \sim 4 \times 10^4 \frac{\lambda^{coh}}{\widetilde{\lambda}^{coh}}.$$

The mean energy of neutrinos generated via electron capture is comparable to the electron Fermi energy, thus,

$$\left(\widetilde{\lambda}^{coh}\right)^{-1} \sim 3.9 \times 10^{-5}\, cm^{-1} \widetilde{\rho}_{12},$$

$\widetilde{\rho}_{12}$ is the density in units of $10^{12}\, g\, cm^{-3}$. Therefore,

$$\frac{\widetilde{t}_{diff}}{t_{diff}} \sim \left(\frac{\widetilde{\rho}_{12}}{\rho_{12}}\right)^{\frac{5}{3}} = \left(\frac{\widetilde{\rho}}{\rho}\right)^{\frac{5}{3}},$$

and

$$\varepsilon_{trap} = \frac{\widetilde{t}_{coll}}{\widetilde{t}_{diff}} \sim 9.6 \times 10^{-10}, \quad (5.7)$$

where $\widetilde{t}_{coll}$ is the hydrodynamical collapse timescale is in order of magnitude the free-free timescale: $\widetilde{t}_{coll} \sim (G\, \widetilde{\rho})^{-\frac{1}{2}}$, $\widetilde{\rho}$ is the mean density of the collapsing core. Assuming that the PD is thin at EH - $\varepsilon_d \sim 10^{-6}$, the eq. (5.5) gives the spectral flux:

$$\frac{d\, \widetilde{L}_{\nu\,\varepsilon}^{URCA}}{d\, y_1} = \left(2.0 \times 10^{31}\, erg\, s^{-1}\right) \times$$
$$\left(\frac{3\pi^4 y_1^4}{8} + \frac{5\pi^2 y_1^5}{12} + \frac{y_1^7}{24}\right) \left(e^{y_1} + 1\right)^{-1}, \quad (5.8)$$

where $y_1 \equiv \widetilde{E}_\nu \widetilde{\beta}$, and $\widetilde{\beta} \equiv (k_B\, \widetilde{T})^{-1}$. The total flux is

$$\widetilde{L}_{\nu\,\varepsilon}^{URCA} \sim 1.8 \times 10^{34}\, erg\, s^{-1}. \quad (5.9)$$

In the same manner we get for the pionic reactions:

$$\widetilde{L}_{\nu\,\varepsilon}^{\pi} = \varepsilon \frac{\widetilde{M}}{M} \frac{\rho}{\widetilde{\rho}} \left(\frac{\widetilde{T}_F}{T_F}\right)^3 \left(\frac{T_e^{SPC}}{T_e^n}\right)^3 L_\nu^\pi, \quad (5.10)$$

where L_ν^π refers to the neutron star ,[20] which gives

$$\widetilde{L}_{\nu\,\varepsilon}^{\pi} \sim 4.2 \times 10^{36}\, erg\, s^{-1}. \quad (5.11)$$

The URCA-energy-loss rate of the SPC_{II} due to the neutrino emission process $d \to u + e^- + \bar{\nu}_e$ can be written

$$\widetilde{\epsilon}_{\bar{\nu}\,\varepsilon}^{q} = 6V^{-1} \varepsilon \left(\prod_{i=1}^{4} V \int \frac{d^3 \widetilde{p}_i}{(2\pi)^3}\right) \widetilde{E}_{\bar{\nu}} V (2\pi)^4 \delta^{(4)} (\widetilde{p}_d -$$
$$\widetilde{p}_{\bar{\nu}_e} - \widetilde{p}_u - \widetilde{p}_e) \frac{|\, M\, |^2}{\prod_{i=1}^{4} 2\widetilde{E}_i V} f_d (1 - f_u)(1 - f_e), \quad (5.12)$$

where the four-vectors $\widetilde{p}_i$ are numbered as $i = 1, 2, 3, 4 \equiv d, \bar{\nu}_e, u, e^-$, V is the normalization volume, $|\, M\, |^2$ is the squared invariant amplitude averaged over the initial d-quark spin and summed over the final spins of the u quark and the electron,

$$|\, M\, |^2 = \frac{1}{2} \sum_{\sigma_1\, \sigma_2\, \sigma_3} |\, M_{fi}\, |^2 \simeq 64 G^2 \cos^2 \theta_C (\widetilde{p}_1 \cdot \widetilde{p}_2)(\widetilde{p}_3 \cdot \widetilde{p}_4),$$

$f \equiv \left[\exp \widetilde{\beta}(\widetilde{E}_i - \mu_i) + 1\right]^{-1}$, the blocking factors $1 - f_i$ accounting for the distribution of final states reduce the reaction rate, which ensure that the exclusion principle is obeyed; the factor 3 takes account of three color degrees of freedom, and 2 is the spin of the initial d quark. It is straightforward to evaluate the cooling rate integral eq. (5.12) by notifying that the neutrinos are produced thermally, thus, we can neglect the neutrino momentum in the momentum conservation law, also, for each degenerate species only a fraction $\frac{\widetilde{T}}{\widetilde{T}_F}$ effectively contribute to the cooling rate. There are one such initial species and two such final species. Using the standard technique for calculating nonequilibrium processes for degenerate fermions [9] we finally arrive at the explicit mean spectral flux of the EHE antineutrinos above the GZK cutoff from SPC$_{II}$:

$$\frac{d\widetilde{L}^q_{\bar\nu\,\varepsilon}}{d\,y_2} = \left(8.1\times10^{41}\,erg\,s^{-1}\right) y_2^3\left(\pi^2+y_2^2\right)\left(e^{y_2}+1\right)^{-1}, \tag{5.13}$$

where $y_2 \equiv \widetilde{E}_{\bar\nu}\widetilde{\beta}$. Continuing along this line we get the flux of the AGN-neutrinos due to the processes $u+e^- \to d+\nu_e$

$$\frac{d\widetilde{L}^q_{\nu\,\varepsilon}}{d\,y_1} = \left(1.5\times10^{41}\,erg\,s^{-1}\right) y_1^4\left(\pi^2+y_1^2\right)\left(e^{y_1}+1\right)^{-1}. \tag{5.14}$$

A different $\widetilde{E}_\nu$ dependence arisen in the latter is due to the partial restriction of the electron's phase space which introduces an extra factor of $(\widetilde{E}_e\widetilde{\beta})^{-1}$. Since eq. (5.14) is independent of the electron's Fermi energy a potential source of ambiguity is eliminated. The resulting total mean EHE neutrino luminosity of the SPC$_{II}$ will then be dramatically larger than of quark star: [21]

$$\widetilde{L}^q_{\nu\,\varepsilon} \sim 1.4\times10^{44}\,erg\,s^{-1}. \tag{5.15}$$

Note that the nuetrino luminosity of the SPC could achieve its theoretically maximum possible value when the total binding energy would be completely emitted as neutrinos in presence of neutrino trapping:

$$\widetilde{L}^{max}_\nu = \varepsilon\,\frac{\frac{G\widetilde{M}^2}{\widetilde{R}}\,(G\,\widetilde{\rho})^{-\frac{1}{2}}}{\frac{GM^2_\odot}{R_{nuc}}\,(G\,\rho_{nuc})^{-\frac{1}{2}}}L^{max}_\nu \sim 1.4\times10^{58}erg\,s^{-1}, \tag{5.16}$$

where $L^{max}_\nu \sim 10^{57}\,erg\,s^{-1}$ is the maximum possible value of neutrino luminosity of the neutron star in absence of neutrino trapping .[6] This exhibits the strong inequality $\widetilde{L}^q_{\nu\,\varepsilon} << \widetilde{L}^{max}_\nu$, which demonstrates the inability in reality of reaching this limit.

5.2 *Cooling time*

Inserting the corresponding luminosities into cooling equation

$$\frac{d\,\widetilde{U}}{d\,t} = -(\widetilde{L}_{\nu\,\varepsilon}+\widetilde{L}_\gamma), \tag{5.17}$$

and integrating gives the time for the SPC to cool from an initial interior temperature $\widetilde{T}(i)$ to a final temperature $\widetilde{T}(f)$. The thermal energies reside almost exclusively in degenerate fermions:

$$\begin{aligned}\widetilde{U}_n(SPC_I) &\simeq \frac{\widetilde{M}}{M}\left(\frac{\rho}{\widetilde{\rho}}\right)^{-\frac{1}{3}}\left(\frac{\widetilde{T}_{F\,9}}{T_{F\,9}}\right)^2 U_n \sim 1\times10^{57}\,erg,\\ \widetilde{U}_q(SPC_{II}) &\simeq \frac{\widetilde{M}}{M}\left(\frac{\rho}{\widetilde{\rho}}\right)^{-\frac{2}{3}}\left(\frac{\widetilde{T}_{F\,9}}{T_{F\,9}}\right)^2 U_q \sim 2\times10^{56}\,erg,\end{aligned} \tag{5.18}$$

Neglecting interactions, in the case if $\widetilde{L}_{\nu\,\varepsilon}$ dominates, we get respectively for the nucleon-modified URCA processes, pionic reactions, quarks and neutrino pair bremsstrahlung:

$$\Delta\,t(\text{URCA}) \sim 1\times10^{15}\,yr\left(\frac{\widetilde{\rho}}{\rho_{nuc}}\right)^{-\frac{1}{3}}\widetilde{T}_9^{-6}(f)\left\{1-\left[\frac{\widetilde{T}_9(f)}{\widetilde{T}_9(i)}\right]^6\right\},$$

$$\Delta\,t(\text{pions}) \sim 6.6\times10^{8}\,yr\,\theta^{-2}\left(\frac{\widetilde{\rho}}{\rho_{nuc}}\right)^{\frac{1}{3}}\widetilde{T}_9^{-4}(f)\left\{1-\left[\frac{\widetilde{T}_9(f)}{\widetilde{T}_9(i)}\right]^4\right\},$$

$$\Delta\,t(\text{quarks}) \sim 1.2\times10^{11}\,yr\left(\frac{\widetilde{n}}{n_{nuc}}\right)^{-\frac{1}{3}}\widetilde{T}_9^{-4}(f)\left\{1-\left[\frac{\widetilde{T}_9(f)}{\widetilde{T}_9(i)}\right]^4\right\},$$

$$\Delta\,t(\text{brem}) \sim 2\times10^{15}\,yr\frac{\widetilde{M}}{M_{cr}}\left(\frac{\widetilde{\rho}}{\rho_{nuc}}\right)^{-\frac{2}{3}}\widetilde{T}_9^{-4}(f)\left\{1-\left[\frac{\widetilde{T}_9(f)}{\widetilde{T}_9(i)}\right]^4\right\},$$

where $\theta \sim 0.3$ is an angle measuring the degree of pion condensation, M_{cr} is the mass of the crust.

5.3 *Light relic neutrinos*

The suggested microscopic approach predicts as well the relic CBN with light mass produced in hot primordial Big Bang phase. Actually, due to the phase transition of particles in protomatter, the CB relic neutrinos liberated from the primordial early Universe protomatter-plasma, when the Universe was only one second, have acquired distorted mass stemming from eq. (2.32):

$$\widetilde{m}_\nu\,c^2 = \widetilde{E}_\nu\,x\,k_\perp\sqrt{1-\frac{1}{1+x^2}} \tag{5.19}$$

where ordinary m_ν mass is zero, $\widetilde{\boldsymbol{P}}_\nu \equiv \frac{\widetilde{E}_\nu}{c}\boldsymbol{k}$. Expansion of the Universe cools the ultrarelativistic neutrino protomatter, and, as soon as it reaches the boundary of protomatter-ordinary matter, the neutrinos still have retained a residual mass at rest $\widetilde{m}^d_\nu$. Taking the energies at this boundary to be of average central value of neutron star $E_d \sim 100\,MeV$, in isotropic case one gets

$$\widetilde{m}^d_\nu\,c^2 \simeq \sqrt{\frac{2}{3}}\,E_d\,x_d^2, \tag{5.20}$$

where $x_d \sim 1.2 \times 10^{-4}$ is the boundary value of the ID-potential. Therefor

$$\widetilde{m}^d_\nu c^2 \sim 1\, eV. \quad (5.21)$$

Such relic light neutrinos cluster into large galactic halos to form HDM.

§6. Conclusions

• Discussed gravitational theory is consistent with the general relativity in description of realistic models of neutron stars ($x \equiv 0$).
• The SPCs with thermodynamical properties strongly different from the thermodynamics of ordinary compressed matter are formed, where, in PC, the energy density and internal pressure have sharply increased proportional to gravitational forces of compression in about 18-20 order of magnitude with respect to corresponding central values of neutron star. This counteracts the collapse and equilibrium holds even for the maximum masses $M_{max} \simeq 3.48 \times 10^8 M_\odot$.
• After formation of the SPC, which accommodates the highest energy scale ($> 10^{21}\, eV$) in the central massive PC, the radius of it equals $\sim \frac{1}{4} R_g$, thus, the SPC is always found inside the EH sphere and could be observed only in presence of a matter to infall to the nucleus. At this early stage the compact SPCs do work in the AGNs absolutely in the same way as the central massive BHs. Here, of course, a key open question is to clear up the mechanisms that trigger the activity, and how a large amount of matter can be funneled to the central regions, to fuel this activity. Then, the most important next stage of evolution of SPCs succeeds, which is in strong contrast to standard BH physics. Namely, during the time $\leq 10^8\, yr$ this infalling matter has formed a PD around the PC, which gradually becomes very thin at reaching out the edge of EH. The MSC effects acts in the small region of intersection of PD with the EH, i.e.: the ID-field is switched on and a metric singularity disappeared, along with it a sharply increase of gravitational forces has ceased, and the particles may escape, in principle, through this vista to the outside world.
• The suggested microscopic approach to BH physics predicts a large flux of primary EHE extra-galactic AGN-neutrinos above $10^{21} eV$, even after the neutrino trapping in the superdense medium, produced by the predominant neutrino cooling of the SPC via simple or 'modified' URCA processes, and pionic reactions. The part of it may lose to produce, further, the secondary EHE electrons in the AD and in a torus of hot gas surrounding the SPC, which, in turn, may give rise a secondary flux of the γ-rays. It supports the idea that AGNs are strong EHE γ-ray emitters. This approach predicts as well the mean density of the relic CBN with the mass $\sim 1\, eV$ produced in hot primordial Big Bang phase.
• Final observations: All the previous assertions made in simple case of ideal Fermi neutron gas configurations [1] are due directly to global properties of space-time continuum, it is why they hold irrespective of more realistic state equations used in .[7,8] Certainly, unlike the neutron stars, the integral characteristics of AGNs have weakly depended of central values of density and pressure, namely the global properties of the space-time continuum just are really responsible for all the processes occurred at short distances. It is also remarkable that such a view point is developed within a microscopic approach to the Standard Model of particle physics and its SUSY-extension [3] employing the multiworld geometry.

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