Birth and life of massive black holes

V. I. Dokuchaev

Institute of Nuclear Research of the Academy of Sciences of the USSR

(Submitted November 19, 1990)
Usp. Fiz. Nauk **161**, 1–52 (June 1991)

The problems of massive black holes in galactic nuclei of different types are reviewed. The dynamical evolution of compact star systems ends naturally in a gigantic concentrated mass of gas, containing an admixture of surviving stars, that unavoidably collapses into a black hole. The subsequent joint evolution of the remnant star system with a massive black hole at the center leads either to the phenomenon of a bright central source in the nuclei of active galaxies and quasars or to the opposite case of a "dead" frozen black hole in the nucleus of a normal galaxy.

1. INTRODUCTION

The nuclei of active galaxies (immense phenomena in the observable universe) and quasars (their most impressive variant) are connected with the release of gravitational energy in very massive sources. Practically the only candidate advanced for such sources by the modern physical paradigm is an accreting massive black hole (BH). The most striking characteristics of the nuclei of active galaxies and quasars—enormous luminosity, long duration of the active stage, rapid variability, directed emissions in the form of gigantic collimated jets—can be explained in a consistent fashion only by introducing massive black holes.[1-7] Alternative models, such as supermassive stars[8-11] or compact clusters of colliding and exploding stars,[12] also can serve as powerhouses for the nuclei of active galaxies. However such objects have too short a lifetime and they are only an intermediate stage of the collapse of large masses of matter into a black hole. If all galaxies have a common origin, then massive black holes can also be found in the nuclei of normal galaxies, but in the passive state of a "dead" quasar or in a state of weak activity owing to weak accretion flows or extremely low mass-to-radiation conversion efficiency.

Aside from heuristic arguments and theoretical models of accreting black holes, observational evidence of the presence of compact massive formations in both the nuclei of the nearest normal galaxies,[13-19] including our own galaxy,[20,21] and the nuclei of active galaxies[22-29] is also gradually accumulating. It must nonetheless be stressed that there is still no decisive proof confirming the relativistic character of the central sources in galactic nuclei and making it possible to identify the most important characteristic of a black hole—the event horizon.

Beginning with the discovery of quasars[30] in 1963 and the first theoretical works in the 1960s[31-34] the process of working out a solution to the problem of the activity of galactic nuclei has developed into an entire "industry" of observations in different wavelength ranges, data analysis, modeling, and theoretical research, rooted in virtually every branch of physics. The overall modern picture of the activity phenomenon is as follows. The central source of the nucleus of an active galaxy or quasar is a massive black hole with mass $M_h \sim 10^6$–$10^{10}\ M_\odot$. These massive black holes either are of a primordial nature, with respect to the galaxies, and their history goes back to the early stages of expansion of the universe[35-38] or, more likely, they arise in the course of the natural evolution of galactic nuclei.[1-7] The luminosity of the central source can be close to the Eddington limit $L_E \approx 1.3\cdot 10^{46} M_h/10^8 M_\odot$ ergs·s^{-1}. Accretion rates up to $\dot{M}_h \sim 1$–$10^2\ M_\odot \mathrm{yr}^{-1}$ with gravitational energy being converted into radiation with the efficiency $\eta = L/\dot{M}_h c^2 \sim 0.1$, which can be achieved, for example, by means of disk accretion, are required in order to maintain the gigantic luminosities. The non-steady-state processes occurring in the flow of the accreted plasma are accompanied by fluctuations of the luminosity at minimum scales of the order of the time required for light to traverse a distance equal to the gravitational radius of the central source $r_g = GM_h/c^2 \approx 1.5\cdot 10^{13} M_h/10^8 M_\odot$ cm. The accreted gas, carrying angular momentum, forms around the black hole a disk that is embedded in a hot outflowing corona. Owing to dynamo effects, a magnetosphere with a regular magnetic field, partially threading through the hole, forms in the accretion disk. As a result of the interaction of the black hole and the rotating magnetized plasma disk rarefied funnels oriented along the common rotation axis of the hole and disk are formed. Oppositely directed relativistic jets are generated in these axisymmetric funnels. According to observations the collimated jets or emissions from the central source in galactic nuclei often propagate over a distance many times greater than the size of the galaxies. The enormous linear sizes of the jets indicate that the active stage of galactic nuclei lasts for up to 10^8 yr and possibly even longer. Cosmic rays must inevitably be generated in the jets themselves near the black hole or on the shockfronts forming when the jets interact with the cosmic medium.

There are many detailed reviews and monographs concerning the problems of black holes (see, for example, Refs. 1–7 and 39–43 and the references cited there). We shall confine our attention primarily to the theoretical aspects of the formation, evolution, and activity of massive black holes, whose masses significantly exceed the typical stellar masses $m \sim M_\odot$, in galactic nuclei.

The most natural evolution of galactic nuclei consists of the following successive stages: a contracting compact star cluster, a short-lived supermassive star or disk, a massive

 0038-5670/91/060447-24$01.00

black hole at the center of the remnant of the star system, and a single supermassive black hole at the center of the galaxy. Dissipative processes occurring in the evolving stellar system accelerate the contraction of the system and increase the mass of the core that is capable of collapsing into a black hole. According to this evolutionary scheme, massive black holes are typical residents of galactic nuclei and they are either in an active state of a powerful central source abundantly fed by accretion or in a passive state of a "dead" quasar with meager accretion. Only low-mass and rarefield stellar systems in galactic nuclei with the parameters of globular star clusters are incapable of producing a black hole with mass significantly greater than the star type mass because of large mass loss owing to dispersal of stars out of the system in the process of evolutionary contraction.

2. FORMATION OF MASSIVE BLACK HOLES IN EVOLVING STELLAR SYSTEMS

2.1. Characteristics of central stellar systems in galactic nuclei

Galactic nuclei, which are the bottom of a deep gravitational potential well, are the most appropriate location for concentration of large masses of matter and collapse of this matter into a massive black hole. Additional mass is transported into the central part of the galaxy by stars or gas, for example, in the form of a "cooling" flow accreted from the intergalactic medium onto the galaxy and its nucleus.[44,45] In addition, galaxies with active nuclei often exhibit explicit vestiges of merging with another galaxy in the past or their shapes are strongly distorted by tidal interactions with a nearby companion galaxy.[46,47] Significant mass growth of galactic nuclei is most likely associated with similar rare shocks to the parent galaxies. Massive galaxies, especially those found in rich clusters, evidently undergo several mergings and the activity of their nuclei can be recurrent.

The gas collecting in galactic nuclei contracts rapidly, as compared with their life cycle time, and forms, together with the stars that are already present there, a central star cluster. The contraction of this cluster is accompanied by acceleration of all processes occurring in it, dynamical isolation of the cluster from the rest of the galaxy, and relatively rapid and independent evolution. Galactic nuclei naturally evolve in the direction of secular growth of the star velocities and, correspondingly, of the gravitational potential at the center of the galactic nuclei, and this process leads to the birth of a massive black hole. External actions on galactic nuclei in the form of settling of galactic and intergalactic matter on the nucleus occurs slowly compared with the intrinsic dynamical evolutionary cycles of central star systems. Separate actions, however, such as intense tidal interaction and merging of galaxies, occur only rarely. They are all equivalent to changes in the initial conditions and give rise to different paths of evolution of galactic nuclei.

Self-gravitating star systems do not reach a state of complete statistical equilibrium, since the total energy does not have a minimum in the Newtonian limit. This is manifested in at least two ways. First, as a result of star interactions some stars acquire velocities that exceed the velocity of escape from the cluster and such stars escape (evaporate) from the system.[48,49] Second, some stars can combine into gravitationally bound subsystems, which in many respects evolve independently. Instability of a star system resulting in the dynamical isolation and rapid, as compared with the system as a whole,[50] contraction of its center is called the gravothermal catastrophe.[51] This regime of contraction begins when the difference of the star density between the core and the periphery of the cluster is sufficiently large and the relaxation time of moving stars at the center is substantially different from that at the periphery of the system. Because a state of statistical equilibrium is impossible and because the star interactions are of a collective character the systematic investigation of the dynamics of star systems requires a kinetic approach,[52] i.e., the kinetic equation for the distribution function of the stars must be solved simultaneously with the Poisson equation, or the equations of motion must be directly integrated numerically. Some qualitative features of the dynamical evolution are manifested in the homological approximation, also called the gross-dynamical approximation, when the integral parameters of the system, such as the effective size, the velocity dispersion of the stars, the mass of the core, etc., are studied as a function of time, eschewing the evolution of the spatial structure of the system. In the analysis of the dynamics of systems consisting of very compact objects, for example, neutron stars and black holes of stellar masses, the post-Newtonian corrections must be included in the equations of motion or a completely relativistic approach must be employed.[53–57] In realistic star systems, however, dissipative processes associated with the finite sizes of the stars begin to play a decisive role long before the onset of the relativistic stage of evolution. These processes include, in particular, tidal interactions and disruptive collisions of stars, which accelerate the contraction of the central part of the stellar system and transform it, even at the Newtonian stage, into a self-gravitating compact cloud of gas—a supermassive star.

2.2. Nondissipative contraction of a star system

The classical scheme of dynamical evolution of star systems of the type of galactic nuclei and globular star clusters[58,59] consists of the fact that the exchange of energy between single stars in an isolated star cluster results in expansion and evaporation of the outer part of the system, while the central part undergoes unlimited contraction. As a result, the star system acquires a "compact core—rarefied halo" structure, and in addition the core contraction time is only several tens of times greater than the relaxation times.

If the total mass of the stars in the core of a cluster in dynamic equilibrium is equal to M, where for simplicity the cluster consists of identical stars of mass m, and the velocity dispersion of the stars, averaged over the entire cluster, is equal to v_0, then the effective radius of the core is equal to $R = GM/2v_0^2$. This relation is essentially a definition of R. According to the virial theorem,[60] in the case of dynamical equilibrium the total energy of the system can be written in the form $E = -Nmv_0^2/2$, where $N = M/m$ is the number of stars in the core. The parameters of central star clusters depend on the types of galaxies and vary over wide limits. For the nuclei of normal galaxies the characteristic values are $M \sim 10^6$–$10^8\ M_\odot$ and $v_0 \sim 10^2\ \mathrm{km \cdot s^{-1}}$, and the most likely values in the nuclei of the most active galaxies are $M \geqslant 10^8 M_\odot$ and $v_0 \geqslant 10^3\ \mathrm{km \cdot s^{-1}}$.

Because of the Coulomb nature of the interactions of the stars the change in the trajectories of separate stars in the system is determined primarily by distant encounters and

occurs in a diffusion manner with the characteristic relaxation time[61,62]

$$T = \left(\frac{2}{3}\right)^{1/2} \frac{v^3}{3\pi G^2 m^2 n \Lambda}, \tag{2.1}$$

where $\Lambda = \ln(N/2)$ is the Coulomb logarithm; v and n are, respectively, the local values of the velocity dispersion of the stars and the star density; and, m is the mass of a separate star. In particular, in the disks of spiral galaxies and in most of the volume of elliptical galaxies the relaxation time T is significantly longer than the age of the universe. Relaxation times $T < 10^{10}$ yr occur only in the nuclei of galaxies as well as in the central parts of rich globular star clusters; this is why these objects evolve much more rapidly than the galaxy as a whole. Comparing the relaxation time T with the characteristic dynamical traversal time of the stellar system $t_{\rm dyn} = R/v$, we obtain $T/t_{\rm dyn} \sim N/\ln N \gg 1$. For this reason, the trajectories of stars in a self-gravitating dynamically equilibrium system are regular orbits whose parameters vary slowly over times $t \ll T$.

The qualitative features of the evolution of the central parts of star systems can be traced in the homological approximation, starting from the equation of evolution

$$\frac{\dot{R}}{R} = -\left(\frac{\dot{E}}{E}\right) + 2\left(\frac{\dot{M}}{M}\right), \tag{2.2}$$

which follows from the virial theorem. This equation describes the temporal change occurring in the integral parameters of a self-gravitating system in dynamical equilibrium and is valid for times $t \gg t_{\rm dyn}$. The main factor of the dissipation-free evolution of a star system, approximated by $N = M/m$ point-like bodies, is evaporation of stars out of the system,[48,49,62,63] which occurs with the rate

$$\dot{N}_{\rm ev} = \alpha \frac{N}{T_0}, \tag{2.3}$$

where the numerical constant $\alpha \approx 7.4 \cdot 10^{-3}$ with a Maxwellian star velocity distribution function, $T_0 = T(v_0, n_0)$, and v_0 and n_0 are the velocity dispersion of stars and the star density averaged over the system. Stars whose velocities exceed the escape velocity $v_{\rm esc} \approx 2v_0$ leave the system. The total energy of each such star is close to zero. For this reason, in the star evaporation process the total energy of the stellar system $E \approx$ const and, as follows from Eq. (2.3), the system contracts as $R \propto N^2$. In the process the velocity dispersion of the stars remaining in the system increases as $v_0 \propto N^{-1/2}$. Integrating simultaneously Eqs. (2.2) and (2.3) with $E =$ const and neglecting, as is done everywhere below, the slow change in the logarithmic factor L gives

$$R(t) = R(0)\left[1 - \left(\frac{t}{\tau_{\rm ev}}\right)\right]^{4/7}, \tag{2.4}$$

where $\tau_{\rm ev} = (2/7\alpha)\, T_0(0) \simeq 40T(0)$. In a star system consisting of very compact stars, star evaporation unavoidably leads to a relativistic contraction regime, when $v_0 \to c$ and the Newtonian approximation becomes inapplicable. At the relativistic stage the contraction of the system is now controlled not by star evaporation but by gravitational radiation; this stage ends with the collapse of the central part of the system and the formation of a massive black hole. If, however, the star system consists of normal stars or even degenerate stars of the type of white dwarfs, whose escape speed from the surface of the star $v_p = (2Gm/r_*)^{1/2} \ll c$, where m is the mass of the star and r_* is the radius of the star, then even at the Newtonian stage dissipative processes associated with the finite sizes of the stars start to have a dominant effect on the fate of the system.

2.3. Role of tidal interactions of stars in the evolution of star systems

Because of their dissipative character inelastic interstellar interactions associated with the tidal action of stars on one another "cool" the star system and, generally speaking, accelerate the evolutionary contraction of the system.[64] At the same time inelastic interactions facilitate the separation of the system into gravitationally bound subsystems of stars, and this has a nontrivial consequence—self-production of intrinsic sources of "heat" in the system.[65–67] These are the most easily forming subsystems—binary stars. They exchange energy with the system as a whole and the energy of some of them decreases, and in the process the energy of the rest of the system increases. The intermediaries in this exchange are, most often, separate stars, which acquire in the form of recoil from the binary star excess energy, which is then gradually collectivized by the entire star collection. Under certain conditions the heating of a star system by pairs forming in the system can radically intervene in the evolutionary history of the system and prevent the formation of a massive black hole.[68–70]

Binary stars in rarefied star systems are primordial. They are born in the process of fragmentation of rotating clouds of gas on separate protostars. In dense star systems, however, binary stars are formed owing to the tidal interaction of the stars in two-body encounters and with a much lower probability in nondissipative three-body encounters.[71,72] The tidal mechanism of pair formation operates if the energy dissipated in encounters between stars does not exceed the characteristic binding energy of the stars. This requirement is satisfied only in systems in which the velocity dispersion of the stars $v_0 < v_p$. Because of their relatively large mass binary stars are concentrated at the center of the system. Only hard pairs, whose binding energy x is greater than the mean kinetic energy of the stars $\beta^{-1} = mv_0^2/3$, where m is the mass of a separate star and v_0 is the velocity dispersion of the stars, participate in the heating of the system. However soft pairs, whose binding energy $x < \beta^{-1}$, mainly break up in interactions with separate stars[63,67] without significantly affecting the dynamics of the system as a whole. The fact that the tidal formation of pairs is accompanied by dissipation of energy from the star system while hardening of the pairs, conversely, heats the system makes the self-consistent investigation of the role of binary stars in the evolution of star systems a quite difficult problem, which has not been completely solved even with the help of numerical modeling.[73–77] In what follows we shall study the effect of the tidal interaction of stars on the evolution of a star system using the example of a simple analytical model.

2.4. Dissipative cooling of a star system and the formation of binary stars

In the case of close fly-bys of stars the tidal gravitational interaction excites in the passing stars nonradial oscillations.

The energy for these oscillations comes from the kinetic energy of the stars, and it is gradually dissipated into heat and radiated out.[71] Model calculations based on the polytropic approximation for the internal structure of the stars[78] give for the energy dissipated in a single encounter of identical stars with masses m and radii r_* the value

$$\Delta E \approx 2^{7/3}\frac{Gm^2}{r_*}\cdot\left(\frac{r_*}{R_{\min}}\right)^{10}, \tag{2.5}$$

where $R_{\min}$ is the periastron radius, corresponding to the distance between the stars at the moment of closest approach. If in a single encounter of stars energy $\Delta E > mu^2/4$, where u is the relative velocity of the stars before the encounter, is dissipated in each of them, then the stars form a gravitationally bound pair. Encounters in which less energy is dissipated do not lead to binding of the stars and are equivalent to friction. By analogy to atomic processes, we shall call such encounters free-free (f–f transitions).

Tidal interactions resulting in the binding of the stars, as a rule, engender a pair with nonzero orbital eccentricity e and binding energy x, small compared with the mean kinetic energy of the stars in the system $\beta^{-1} = mv_0^2/3$. The orbital angular momentum J of such a pair is significantly higher than the intrinsic rotational angular momentum of the stars comprising the pair. For this reason subsequent encounters of stars in a pair that are accompanied by an increase in binding energy and circularization of the orbit (decrease of e) will occur with $J \approx$ const. In this case the ratio

$$\frac{x}{1-e} = \frac{G^2m^5}{4J} = \frac{Gm^2}{4R_{\min_0}}, \tag{2.6}$$

where $R_{\min_0}$ is the distance to the periastron at the moment of formation of the pair, will remain constant. Hence the maximum binding energy that a pair can acquire as a result of circularization of its orbit owing to energy dissipation in many subsequent passes through the periastron will be equal to

$$x_c = \frac{x_0}{1-e_0^2}, \tag{2.7}$$

where x_0 and e_0 are, respectively, the binding energy and the eccentricity of the orbit of the pair at the moment of formation. We shall call free-bound (f–b) transitions those encounters which result in a hard pair whose binding energy is $x > \beta^{1}$ already in a single encounter of the stars and which will certainly survive in the system. We shall also classify as f–b processes encounters of stars that end in collisions of the stars. Free-bound-bound processes, in which a soft pair forms first and then transforms into a hard pair owing to further tidal dissipation of energy, are also possible. According to Eqs. (2.6) and (2.7), for this the eccentricity e and the binding energy x of the pair must satisfy the condition $1 - e^2 < \beta x$. An additional restriction follows from the fact that the tidal consolidation time of such a pair to a binding energy $x_c = \beta^{-1}$ corresponding to the maximum hardness must be short compared with the breakup time of a soft pair into single stars[63,67]

$$t_s(x) = \frac{3^{3/2}}{20\pi^{1/2}}\cdot G^{-2}m^{-7/2}n^{-1}\beta^{-1/2}x \sim \beta x\Lambda T. \tag{2.8}$$

Some soft pairs, which do not have enough time to overcome the consolidation limit before they break up, give f–b–f transitions. Nondissipative triple encounters of stars also can engender a bound pair. Most of these pairs do not reach the consolidation limit in the subsequent tidal circularization of the orbits and break up (b–b–f transitions). Only a small fraction of these pairs cross the consolidation limit and survive in the star system (b–b transitions). It is found that consolidated pairs form in a star system predominantly by means of tidal pair interactions of stars.[64,69,72]

The total rate of tidal dissipation of kinetic energy of stars in a star system is

$$\dot{E}_{\text{dis}} = -\frac{N}{2}\cdot\int_0^\infty d^3u f(u) u \int_0^\infty dp\cdot 2np\Delta E(p, u). \tag{2.9}$$

In this expression n is the star density, $f(u)$ is the distribution of stars over their relative velocities u, p is the impact parameter which is equal to

$$p = R_{\min}\left[1 + \left(\frac{4Gm}{u^2R_{\min}}\right)\right]^{1/2} \tag{2.10}$$

and, finally, $\Delta E(p, u)$ is the dissipated kinetic energy of the stars, which is equal to the energy of tidal dissipation (2.7) for f–f transitions, and $\Delta E(p,u) = mu^2/4$ for f–b and f–b–b transitions. In the case of a Maxwellian distribution function $f(u)$ we obtain[64,69]

$$\dot{E}_{\text{dis}} = \left(\frac{2}{\pi}\right)^{1/2}\left(\Gamma(0,9) + \frac{8}{5}\cdot\Gamma(0,8)g\right) g\beta^{-1}\frac{N}{\Lambda T_0}; \tag{2.11}$$

Here $\Gamma(0.9) \approx 1.069$ and $\Gamma(0.8) \approx 1.164$ are gamma functions, and

$$g = 2^{37/30}(Gm^2\beta/r_*)^{-9/10}. \tag{2.12}$$

The contributions of f–b, f–f, and f–b–f processes in Eq. (2.11) are equal to 85, 10, and 5%,[69] respectively, while the contributions of b–b–f and b–b processes are negligibly small. Detailed analysis shows that the optimal conditions for the formation of hard pairs with binding energies, after circularization of the orbits, in the interval $\beta^{-1} < x < x_{\max}$ are satisfied only if the velocity dispersion of the stars falls into the range

$$0{,}52(\beta x_{\max})^{-5/9} < v_0/v_p < 0{,}52, \tag{2.13}$$

where the right-hand side of the inequality corresponds to the condition $g < 1$ in Eq. (2.12). In this case, which is equivalent to the approximation linear in g, the rate of formation of hard pairs in the system is equal to

$$\dot{N}_b \approx \left(\frac{2}{\pi}\right)^{1/2}\left(\Gamma(0,9) - (\beta x_{\max})^{-1/10}\right) g\frac{N}{\Lambda T_0}. \tag{2.14}$$

The maximum binding energy $x_{\max}$ of the pairs corresponds to the limit of merging of the stars in the pair and depends on the types of stars.

2.5. Dissipative evolution taking into account the heating by binary stars

On the average, the binding energy of hard pairs of stars that form in dissipative encounters or arise in a star system at the stage when stars in the system are formed (primordial

pairs) increases only in interactions with separate stars, and in the process such pairs transform into sources of heat in the system. The rate of change of the binding energy x of a hard pair by an amount y in an interaction with an incident separate star is given by the transfer function[67]

$$Q(x, y) \approx 45m^{7/2}\beta^{1/2}G^2 \begin{array}{ll} \times x^{-2}\exp(\beta y) & \text{if } y<0, \\ \times x^{5/2}(x+y)^{-9/2} & \text{if } y>0. \end{array} \quad (2.15)$$

Hence, for $\beta x \gg 1$ the average energy transferred to a separate star is equal to

$$\langle y\rangle = \int_{-\infty}^{+\infty} Q(x, y)y\,dy \left(\int_{-\infty}^{+\infty} Q(x, y)dy \approx \frac{2}{5}x\right)^{-1}. \quad (2.16)$$

In the case of stars having the same mass two-thirds of this energy is transferred to the separate star and one-third goes to changing the motion of the center of mass of the hard pair. For this reason, the interaction of separate stars with hard pairs with $x > x_1 = (135/4)\beta^{-1}$ results in ejection of the former from the system and for $x > x_2 = (315/4)\beta^{-1}$ the hard pair itself is ejected from the system as a result of recoil. Such superhard pairs contribute to the heating of the system only in the case of a single interaction with separate stars. According to Eq. (2.16), the binding energy of a hard pair grows at the rate

$$\left(\frac{dx}{dt}\right)_h = n\int_{-\infty}^{+\infty} Q(x, y)y\,dy \approx \frac{36}{7}G^2m^{7/2}\beta^{1/2}n. \quad (2.17)$$

The emission of gravitational waves in the case when the orbit of the pair is circular results in increase of the binding energy of the pair at the rate[60]

$$\left(\frac{dx}{dt}\right)_{grav} = \frac{2^{11}}{5}\frac{x^5}{Gm^5c^5}. \quad (2.18)$$

Hard pairs participate in the heating of the system only if $(dx/dt)_h > (dx/dt)_{grav}$ or

$$\beta x_{max} \le \beta x_{grav} \approx 4{,}7(Gm)^{3/5}\frac{cn^{1/5}}{v_0^{11/5}}. \quad (2.19)$$

Using the virial theorem we obtain the estimate $\beta x_{grav} \sim (c/v_0)N^{-2/5}$. The binding energy of pairs is also limited by the merging of the stars in the pair, which happens[79] when the length of the semi-major axis of the pair $\alpha \leqslant 1.2r_*$, which corresponds to the limiting binding energy of the pair

$$\beta x_{max} < \beta x_{tide} \approx 0{,}6(v_p/v_0)^2. \quad (2.20)$$

Pairs no longer heat the system if $x_{max} = \min(x_2, x_{grav}, x_{tide})$. The rate at which a star system with velocity dispersion of the stars in the interval (2.13) is heated by the hard pairs formed is equal to

$$\dot{E}_h = \left(\frac{2}{\pi}\right)^{1/2}\left(A(x_{max})g - B(x_{max})\right)\beta^{-1}\frac{N}{\Lambda T_0}, \quad (2.21)$$

where the tidal factor g is given by Eq. (2.12)

$$A(x_{max}) = \begin{cases} \frac{7}{5}\beta x_{max}[\Gamma(0,9) - (\beta x_{max})^{-1/10}] & \text{if } \beta^{-1} < x_{max} < x_1, \\ 15[(\beta x_2)^{-1/10} - (\beta x_{max})^{-1/10}] + 37{,}57 & \text{if } x_{max} > x_2, \end{cases} \quad (2.22a)$$

$$B(x_{max}) = \begin{cases} \frac{1}{10}\left[\ln(\beta x_{max}) + C\right] & \text{if } \beta^{-1} < x_{max} < x_1, \\ 0{,}242 & \text{if } x_{max} > x_2, \end{cases} \quad (2.22b)$$

where $C \approx 0.577$ is the Euler constant. The formulas for $A(x_{max})$ and $B(x_{max})$ in the case $x_1 < x_{max} < x_2$ are very complicated and are not presented here. For star velocities outside the interval (2.13) the rate of formation of hard pairs and the rate of heating of the system associated with them are too low to affect significantly the dynamical evolution of the system. In this case we shall formally set $A(x_{max}) = B(x_{max}) = 0$. As a result the rate of change of the total energy of the system $\dot{E} = \dot{E}_{dis} + \dot{E}_h$ is the sum of the dissipative cooling (2.11) and the heating by pairs (2.21). With the help of Eqs. (2.12) and (2.22) it can be shown that the condition $\dot{E} > 0$, when heating by pairs transforms into dissipative cooling, can be satisfied only if $\beta x_{max} > 4$. From the equation of evolution (2.2) and Eqs. (2.11), (2.13), and (2.21) it follows that heating of the star system by hard pairs can compete with dissipative cooling only if $v_0 \ll v_p$, which is equivalent to values of the parameter $g \ll 1$.

Binary stars radically change the fate of the central part of a star system, if the evolutionary growth of the velocity dispersion of the stars stops as a result of the heating of the system by hard pairs. Substituting into the equation of evolution (2.2) the rate of change of the energy of the system owing to dissipative cooling (2.11) and heating by the pairs (2.21) as well as the rate of change of the mass of the system owing to the star evaporation we find that the condition $\dot{v}_0 = 0$ is satisfied if $v_0 = v_{cr}$, where the critical velocity is given by

$$v_{cr} = \left(\frac{3}{2}\right)^{1/2} 2^{-67/54}\left[\frac{3(\pi/2)^{1/2}\alpha\Lambda + B(x_{max})}{A(x_{max}) - \Gamma(0,9)}\right]^{5/9} v_p. \quad (2.23)$$

The dependence of v_{cr} on the total number of stars in the system $N = M/m$, under the assumption that x_{max} is determined by the gravitational radiation of the pairs, is shown in Fig. 1. The value of v_{cr} is minimum for $x_{max} > x_2$. The nu-

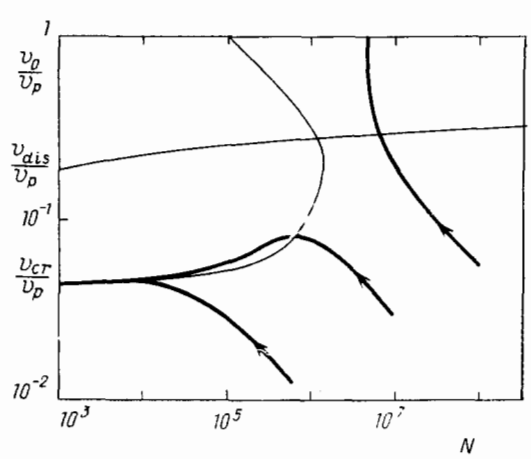

FIG. 1. The evolutionary trajectories (heavy lines) of star systems in the plane (velocity dispersion v_0 of stars, number of stars N). The fine lines represent the critical velocity v_{cr} under the assumption that the maximum binding energy of pairs participating in the heating of the system is limited by their gravitational radiation, and v_{dis}, which determines the limit of accelerated contraction of star systems.

merical value of v_{cr} in this case depends logarithmically on N and for $N \sim 10^7$ is equal to

$$v_{cr} \approx 30(m/M_\odot)^{1/2}(r_*/R_\odot)^{-1/2}\ \mathrm{km \cdot s^{-1}}. \tag{2.24}$$

For $\dot{E} > 0$ the evolutionary equation (2.2) can be written in the form

$$\frac{\dot{R}}{R} = \left[\left(\frac{v_0}{v_{cr}}\right)^{9/5}(1+\Delta) - (2+\Delta)\right]\frac{\alpha}{T_0}, \tag{2.25}$$

where the parameter Δ is given by

$$\Delta = \frac{B(x_{max})}{3(\pi/2)^{1/2}\alpha\Lambda}. \tag{2.26}$$

The simultaneous solution of Eqs. (2.23) and (2.25) for $x_{max} > x_2$ and neglecting, as has already been stipulated, the weak change in the logarithmic factors is[70]

$$\rho = \nu[\vartheta - (\vartheta - 1)\nu^{9(1+\Delta)/10}]^{10/9}, \tag{2.27}$$

$$t = \frac{T_0(0)}{\alpha}\int_\nu^1 z[\vartheta - (\vartheta - 1)z^{9(1+\Delta)/10}]^{5/3}dz, \tag{2.28}$$

$$v_0(t) = v_0(0)(\nu/\rho)^{1/2}, \tag{2.29}$$

where the numerical value of the parameter is $\Delta \approx 0.6(\Lambda/15)^{-1}$ and the dimensionless variables $\rho = R(t)/R(0)$ and $\nu = N(t)/N(0)$ and the parameter $\vartheta = (v_0(0)/v_{cr})^{9/5}$ were introduced. The classical regime of nondissipative contraction of the system (2.4) is obtained from here by setting $\Delta = \vartheta = 0$. An important feature of Eq. (2.25) and its solution is the presence of negative feedback owing to the competition between the cooling and heating processes. The negative feedback controls the changes in the radius of the system R and the velocity dispersion v_0 of the stars in the system. When for $v_0 > v_{cr}$, more accurately, when $\vartheta > 2$, R increases, v_0 and the heating decrease. As a result the system reaches some maximum size and starts to contract. Conversely, when for $v_0 < v_{cr}$ R decreases, v_0 increases and correspondingly the heating intensifies, so that the contraction of the system slows down. In any case the negative feedback between the changes in R and v_0 drive the system into a contraction regime with $R \propto N$ and $v_0 \to v_{cr} \approx$ const. In contrast to the classical nondissipative contraction of the system (2.4) the slower contraction of the system accompanying heating by binary stars does not lead to unlimited growth of the velocity dispersion of the stars, and this makes it impossible for a massive black hole to form. If the mass of the system continuously decreases owing to star evaporation, the velocity dispersion of the stars in the system with solar type stars freezes at the level $v_0 = v_{cr} \approx 30\ \mathrm{km \cdot s^{-1}}$. This regime of moderated contraction is evidently realized in the central parts of globular star clusters, whose cores evaporate completely without collapsing into a black hole. The process of formation of a dense core in globular star clusters, the subsequent contraction of the core, and complete evaporative dissipation can be of a recurrent character; this is confirmed by model numerical calculations.[76,77] If, however, the star system contains quite a large number of primordial hard pairs, then these pairs generally give rise to an overall expansion of the system,[80,81] which is accompanied by a decrease of the velocities of separate stars.

Conditions under which energy dissipation processes predominate over heating by hard pairs, when $E < 0$, are favorable for formation of massive black holes. Such conditions are realized in central star clusters of galactic nuclei which are significantly more massive and dense than globular star clusters. In the limit, when dissipative cooling predominates over heating by pairs the evolutionary equation assumes the form

$$\frac{\dot{R}}{R} = -\left[\left(\frac{v_0}{v_{dis}}\right)^{9/5} + 2\right]\frac{\alpha}{T_0}, \tag{2.30}$$

where

$$v_{dis} = \left(\frac{3}{2}\right)^{1/2} 2^{-67/54}\left[\frac{3(\pi/2)^{1/2}\alpha\Lambda}{\Gamma(0,9)}\right]^{5/9} v_p. \tag{2.31}$$

The first term in Eq. (2.30) is related with the tidal dissipation of energy and the second term is related with star evaporation. The simultaneous solution of Eqs. (2.3) and (2.30) is

$$\rho = \nu\left[(1+\vartheta_{dis})\nu^{9/10} - \vartheta_{dis}\right]^{10/9}, \tag{2.32}$$

$$t = \frac{T_0(0)}{\alpha}\cdot\int_\nu^1 z\left[(1+\vartheta_{dis})z^{9/10} - \vartheta_{dis}\right]^{5/3}dz, \tag{2.33}$$

where the parameter $\vartheta_{dis} = (v_0(0)/v_{dis})^{9/5}$ and as before $v_0(t) = v_0(0)(\nu/\rho)^{1/2}$. The nondissipative contraction regime is obtained from here for $\vartheta_{dis} = 0$. A characteristic feature of the solutions (2.32) and (2.33) is that a finite fraction of stars

$$\nu_{dis} = \left[\vartheta_{dis}/(1+\vartheta_{dis})\right]^{10/9}, \tag{2.34}$$

remain in the star system as $R \to 0$. This feature of dissipative contraction can be interpreted as formation of a massive central object with the mass $M_h = N(0)m\nu_{dis}$—a possible ancestor of a massive black hole—in the remnant of the star system. In the case $\vartheta_{dis} \ll 1$, i.e., when $v_0 \ll v_{dis}$, the evolution of the star system, as long as the number of stars in it satisfies $\vartheta_{dis}^{10/9} < \nu < 1$, follows the classical scenario $\rho \approx \nu^2$. When the fraction of stars remaining in the system is $\nu < \vartheta_{dis}^{10/9}$, the contraction deviates appreciably from the classical law and proceeds much more rapidly. In the case $\vartheta_{dis} \gg 1$ the contraction of the system proceeds already at the outset much more rapidly than in the classical case. Thus dissipative processes substantially accelerate the evolutionary contraction of star systems with a sufficiently large velocity dispersion of the stars: $v_0 > v_{dis} \approx 190(\Lambda/15)^{5/9}(m/M_\odot)^{1/2} \times (r_*/R_\odot)^{-1/2}$ $\mathrm{km \cdot s^{-1}}$.

When the stars participating in the heating of the system have the maximum binding energy $x_{max} < x_2$ the evolutionary equation (2.2) does not have a simple analytical solution because of the complicated character of the competition between the cooling and heating by the vapors which are now no longer ejected from the system but rather cease to exist owing to the radiation of gravitational waves and merging. However, by using the critical velocity v_{cr}, separating the regimes of increasing and decreasing velocity dispersion v_0 of the stars in the system, it is possible to follow qualitatively the evolutionary trajectory of the system in this case also. Some such trajectories are shown in Fig. 1. The

conditions for rapid dissipative contraction and formation of a dense massive remnant are realized when the number of stars in the system is $N = M/m \geqslant 10^7$ and the velocity dispersion of the stars is $v_0 \geqslant 190$ km·s^{-1}. Destructive collisions of stars, resulting ultimately in the formation of a supermassive star or disk[4,83-85] or, finally, a gas shell with a massive black hole at the center, play a significant role in the subsequent evolution of the remnant of the star system.

2.6. Evolution at the stage when star collisions predominate

Collisions between stars approaching one another with a relative velocity u less than the escape velocity from the surface of the star $v_p = (2Gm/r_*)^{1/2}$ will result in merging or complete disruption of the stars, possibly accompanied by an explosion of the supernova type. A frontal collision with $u > v_p$ necessarily ends with the disruption of both stars. Following Ref. 82, we shall assume that the disruptive collisions of stars occur more efficiently for encounters at distances less than the radius r_* of the stars. The cross section of such a process is $\sigma = \pi r_*^2 [1 + 2(v_p^2/u^2)]$. Collisions have a predominant effect on the evolution of a star system for $v_0 > v_p$, when gravitational focusing is weak and the collision cross section is close to its geometric value. In this limit the rate of disruptive collisions of stars in the system is equal to

$$\dot{N}_{\rm coll} \approx \left(\frac{v_0}{v_p}\right)^4 \frac{N}{\Lambda T_0}. \tag{2.35}$$

In Eq. (2.35) and similar expressions presented below the rate of collisions between stars, which is determined only by close encounters, is expressed in terms of the relaxation time $T_0 = T(v_0, n_0)$, which is related with distant encounters and is given by the expression (2.1), in order to be able to make a direct comparison with the rates of other evolutionary processes.

A star system in which $v_0 > v_p$ consists of a decreasing number of stars and an increasing number of gas clouds—the products of star collisions. The gas clouds participate in the virial balance of the system equally with the stars, so long as they are small compared with the size of the system. In each disruptive collision the total kinetic energy of the two stars is dissipated, so that the total rate of dissipation of the virial energy in the system is equal to $\dot{E} = 2E\dot{N}_{\rm coll}/N$ and makes a predominant contribution to the evolutionary equation (2.1). Star collisions do not change the total mass of the system, but in the general case they do change the total number N of its constituent objects: stars and separate clouds of gas. Energy dissipation brings about rapid contraction of the star system over a characteristic time $\tau_{\rm coll} \sim [v_p/v_0(0)]^4 \Lambda T_0(0) < N t_{\rm dyn}$, where $t_{\rm dyn} = R/v_0$. As an illustrative example of the temporal evolution of the parameters of the system we shall examine the approximation when two clouds with approximately equal masses form on the average as a result of disruptive collisions of stars so that $\dot{N} = \dot{M} = 0$. Integrating Eqs. (2.2) and (2.35) then gives the law of contraction of the star system

$$\rho = \left(1 - t\tau_{\rm coll}^{-1}\right)^{2/7}, \tag{2.36}$$

where the collisional evolutionary time $\tau_{\rm coll} = (v_p/v_0(0))^4 \Lambda T_0(0)/7$ itself is virtually independent of the number clouds formed in each star collision. The system continues to contract according to the law (2.36) until its size reaches $R \sim N^{1/2} r_*$, at which point $\tau_{\rm coll} \sim t_{\rm dyn}$, and the system ceases to exist as a collection of separate stars and transforms into a compact massive cloud of gas. The subsequent gas-dynamic stage of evolution of this gravitationally strongly bound cloud is accompanied by the formation of a supermassive star or disk and ultimately, even when the disk or cloud is partially fragmented, should end in collapse into one or several massive black holes.[5,83-90]

Thus the classical scheme of the dynamical evolution of star systems, regarded as a collection of $N \gg 1$ point-like bodies, predicts the possibility of the formation of massive black holes in virtually all compact star clusters in which the relaxation time is less than the age of the universe and the deviations from the rotational spherical symmetry are quite small. The central contraction of mass in galactic nuclei and in globular star clusters pertain precisely to such systems. This simple scheme, however, becomes much more complicated when inelastic star interactions, which are associated with the finiteness of the sizes of the stars, are taken into account. Inelastic star interactions have the character of tidal forces and, apart from purely dissipative "cooling" of star systems, they can lead to the formation of binary stars, which in turn "heat up" the star systems. In particular, the formation of pairs by tidal forces and the energy pumping into the star system associated with the pairs apparently prevent the formation of massive black holes in globular star clusters. At the same time, dissipative star interactions accelerate contraction and collapse of the most compact galactic nuclei, thereby increasing the mass of the massive black holes formed in them.

3. SUPERMASSIVE STAR WITH A BLACK HOLE AT THE CENTER

3.1. Structure and evolution of a supermassive star

Disruptive collisions of stars in a star system with velocity dispersion $v_0 > v_p = (2Gm/r_*)^{1/2}$ lead to rapid contraction of the system into a dense cloud owing to dissipation of star kinetic energy. The cloud consists of a mixture of gas, surviving stars, and possibly black holes of star type masses. The further fate of this cloud, generally speaking, is not unique and depends on parameters of the cloud, such as its mass, size, and total angular momentum. A cloud will contract due to efflux of gas and loss of energy to radiation and it will ultimately collapse into a black hole or, if rapid rotation is present or global instability develops, it will fragment into separate parts. In the latter case the exchange of energy between the separate fragments can cause the main mass of the starting cloud to break apart. Even with this development of events the growth of the gravitational potential of the remnant of the initial cloud is nonetheless unavoidable. In any variant, at the intermediate stage of collapse into a black hole a supermassive star forms in any star system that has passed through the stage of disruptive star collisions.[8-11] The collisional stage and the supermassive-star stage will not occur under conditions of evolutionary contraction and collapse only in the case of nonrelativistic systems consisting entirely of neutron stars and black holes of star type masses.

The structural characteristics of massive stars having masses $M_s > 100 M_\odot$ are determined by the fact that deep in these stars the radiation pressure is higher than the gas pressure.[1,2,91,92] In the case when a massive star consists entirely

of hydrogen the ratio of the average gas pressure P_g in the star to the total of the gas pressure plus the radiation pressure $P_g + P_r$ is equal to

$$\delta = \frac{P_g}{P_g + P_r} \approx 8{,}6(M_s/M_\odot)^{-1/2}. \quad (3.1)$$

The internal structure of a massive star in which $\delta \ll 1$ is described by the standard Eddington model with a polytropic equation of state and specific-heat ratio $\gamma = (P/\rho)^{-1}(\partial P/\partial \rho) \approx 4/3$. The heat flow from the deep interior of a massive star is transferred to the outside convectively,[92] since the temperature gradient is greater than its adiabatic value. Energy is transferred by means of radiant heat conduction only in the outer layers of the star, where the gas density is low. The condition of equilibrium of the surface layers leads to the luminosity of the stationary massive star (Eddington's condition or limit) $L_E = M_s c^2/t_E$, where $t_E = \sigma_T c/4\pi G m_P \approx 5 \cdot 10^8$ yr determines at the same time the absolute upper limit of the lifetime of such a star. GTR effects as well as electron-position pair production[1,2] result in hydrodynamic instability and collapse of the massive star at the stage when the radius of the star R_s is still significantly greater than its gravitational radius $R_g = GM_s/c^2 \approx 1.5 \cdot 10^{13}(M_s/10^8 M_\odot)$ cm. The radius of instability in the case of a nonrotating supermassive star is equal to, in the post-Newtonian approximation,[93]

$$R_{hi} \approx 2{,}4 \cdot 10^5 (M_s/M_\odot)^{3/2} \text{ cm}. \quad (3.2)$$

When the mass is sufficiently large the collapse of the star occurs under conditions when nuclear reactions, which could be the last obstacle to the formation of a black hole, are still not ignited in it. The density of a collapsing supermassive star which contracts to the gravitational radius is comparatively low and no problems arise with the behavior of the superdense matter that would make it difficult to analyze the collapse of neutron stars and formation of supermassive black holes in the early universe. Because the nuclear sources of energy are ineffective the lifetime of a supermassive star at the stage of stationary contraction up to the moment of instability is short and is equal to $\tau \sim t_E(R_g/R_{hi})$. Rotation, turbulence, and the intrinsic magnetic field stabilize a supermassive star, and this decreases the radius of instability R_{hi} and correspondingly increases the lifetime of the star.[9-11,94-96]

3.2. A central black hole as an energy source of a supermassive star

A supermassive star can also be stabilized by a massive black hole of mass $M_h < M_s$ at its center, if accretion on the black hole makes an appreciable contribution to the energetics of the parent star. A supermassive star can contain a massive black hole, which is formed either in the process of merging and accretional growth of a black hole of star type masses in the primary star cluster or with the collapse of part of the matter of the compact gas cloud remaining from the self-disruptive star system. In order for a supermassive star to be in a steady state the radius of the star must exceed the radius of instability (3.2) and the luminosity must be maintained at the Eddington level. When the mass of the black hole is $M_h < M_s$ the hole will provide the required release of energy only if a regime of accretion above the Eddington limit relative to it is realized. It is obvious that the radiation of the black hole in the steady-state laminar regime of accretion cannot handle the transfer of the required energy flow because of excess radiation pressure acting on the accreted gas. In the deep interior of a supermassive star, however, energy is transferred by means of convection and the density of the convective heat flux is higher than the local Eddington value. For this reason, in contrast to the case of spherical accretion of rarefied gas, in a dense gas with a high heat capacity the excess gradient of the radiation pressure can be a source of instability. It leads to the development of turbulent convection, sufficient for transfer of super-Eddington (relative to the black hole) heat flux from the deep interior of the supermassive star to the outside. Such a possibility, though it agrees with the model of convection with an effective mixing length,[92] must be regarded as a plausible hypothesis because of the fact that there does not exist a strict theory of turbulent heat transfer in plasma.

We shall determine the radius of a steady-state supermassive star under the assumption that its equilibrium is maintained by means of energy released in the process of supercritical turbulent accretion on a central black hole. The gas pressure and density at the center of a polytropic star with the ratio of specific-heat capacities γ are equal to[2,91,92]

$$P_c = H_1(\gamma) G M_s^{2/3} p_c^{4/3}; \quad (3.1)$$

$$P_c = H_2(\gamma) M_s / R_s^3, \quad (3.2)$$

where with $\gamma = 4/3$ the constants of numerical integration are equal to $H_1(4/3) = 0.364$ and $H_2(4/3) = 12.9$. The effect of a black hole with mass $M_h \ll M_s$ on the distribution of gas in a star with radius R_s is limited to a region of the radius

$$r_h = \frac{GM_h}{a_c^2} = \frac{1}{\gamma H_1 H_2^{1/3}} \frac{M_h}{M_s} R_s, \quad (3.3)$$

where $a_c = (\gamma P_c/p_c)^{1/2}$ is the velocity of sound at the center of the star. If $M_h \ll M_s$, the parameters of the gas in the region $r \simeq r_h$ are close to their values for the center of the star, so that the rate of spherical accretion on the black hole (Bondi accretion) will be equal to

$$\dot{M}_h \simeq \alpha(\gamma) r_g^2 c p_c (a_c/c)^3, \quad (3.4)$$

where the gravitational radius of the black hole is $r_g = GM_h/c^2$ and the numerical constant $\alpha(4/3) \simeq 5.6$. As the supermassive star contracts the accretional luminosity of the black hole at the center of the star $L = \eta \dot{M}_h c^2$, where η is the matter-to-radiation conversion efficiency, increases. Contraction stops, if the rate of energy release of such a central source reaches the Eddington value for a supermassive star $L = L_E$. The equations (3.1), (3.2), and (3.4) with $L = L_E$ make it possible to find the steady-state radius of a supermassive star with a central black hole as the source of energy

$$R_{st} \approx 2 \cdot 10^{20} (\eta/0{,}1)^{2/3} (M_h/M_\odot)^{4/3} (M_s/M_\odot)^{-1} \text{ cm}. \quad (3.5)$$

The condition of stability of such a nonrotating supermassive star relative to collapse $R_{st} > R_{hi}$ can be written, with the help of Eqs. (3.2) and (3.5), in the form

$$M_h/M_\odot > 5{,}6 \cdot 10^3 (\eta/0{,}1)^{-1/2} (M_s/10^8 M_\odot)^{15/8}, \quad (3.6)$$

and therefore a black hole with mass $M_h < M_s$ can potentially be a central source of energy in a supermassive star with mass $M_s < 7\cdot 10^{12}(\eta/0.1)^{4/7}M_\odot$. Stabilizing factors, such as rotation and the magnetic field, additionally weaken this restriction. Since the mass of a central black hole can only increase, the established steady state of the supermassive star $R_{st} > R_{hi}$ will also be maintained in the future, and in addition the steady-state radius R_{st} of the star increases with time.

High mass-to-radiation conversion efficiency, at a level $\eta \sim 0.1$, can apparently be achieved only if the accretional flux deviates substantially from spherical symmetry in the relativistic region near the event horizon of the black hole; this should be unavoidable in a rotating supermassive star. In this case, in the region of the principal release of energy near the black hole the accretion picture will be analogous to that of a thick disk[97-100] with toroidal isodenses and low effective velocity of radial flow of plasma, as compared with the sound velocity. The last condition must be satisfied in order to realize convective transfer of the released heat to the outside. The lifetime of a supermassive star with a black hole as a central source does not depend on the mass of the star and is equal to $\tau_{st} = \eta t_E \sim 5\cdot 10^7$ yr. During this period of time a supermassive star gradually expands and is "eaten away" by its central black hole, bypassing completely the stage of catastrophic collapse.

4. STRUCTURE OF A STAR SYSTEM WITH A BLACK HOLE AT THE CENTER

4.1. Tidal breakup of stars by a massive black hole

A massive black hole with mass M_h, formed at the center of the remnant of a star cluster with mass $M = Nm$ and effective radius $R = GM/2v_0^2$, affects, by means of its gravitational field, the motion of stars within the radius of influence

$$r_h = \frac{GM_h}{v_0^2}. \tag{4.1}$$

If the mass of the black hole is greater than the total mass of the stars in the system, then the radius of influence r_h is also the radius of the entire star system. If $M_h \ll M$, the sphere of influence of the black hole contains only a small part of the volume and mass of the star system.

A star having mass m and radius r_*, passing by at a distance from the black hole less than the tidal radius

$$r_t \approx 2r_*(M_h/m)^{1/3}, \tag{4.2}$$

will be broken apart by the tidal forces of the hole.[101] Numerical simulation shows[102] that in the process of tidal breakup of a star some fraction of the star's mass is ejected from the neighborhood of the black hole, but a significant part of the star's mass remains in the region of tidal breakup and can form around the black hole a gaseous accretion disk. If the infall time of the matter from such a disk onto the hole is shorter than the time interval between successive events of tidal breakup of stars, then the activity of the black hole will be manifested in the form of separate bright outbursts separated by long time intervals and lasting for several years.[103-105] In the opposite case accretion disks will accumulate. The interaction of disks with one another will result in compensation of their angular momenta and quasispherical accretion on the black hole or formation of a single disk in the case of a star system whose rotational angular momentum is sufficiently large.

High accretional activity of a black hole owing to tidal breakup of stars can be achieved, if the stars break up outside the event horizon of the hole.[101] More precisely, the necessary condition is $r_t > r_{mb}$, where r_{mb} is the capture radius or the minimum radius of stable orbits for stars falling on the hole from large distances. In the case of a nonrotating black hole for parabolic orbits $r_{mb} = 4GM_h/c^2$. In this case the condition $r_t > r_{mb}$ can be represented with the help of Eq. (4.2) in the form

$$M_h < M_{max} = m(c/v_p)^3, \tag{4.3}$$

where $v_p = (2Gm/r_*)^{1/2}$ is the escape speed at the surface of the star. If $M_h > M_{max}$, the stars are swallowed up whole by the black hole; this is not accompanied by an outburst of activity. In particular, $M_{max} \sim 10^8 M_\odot$ for solar type stars and $M_{max} \sim 10^5 M_\odot$ for white dwarfs with $m \sim M_\odot$ and $r_* \sim 10^{-2} R_\odot$. If $M_h < M_{max}$, stars whose velocity vectors are oriented in the direction of the black hole and which fall within the loss cone with half-angle

$$\theta_{lc} = \frac{r_t}{r}\left[1 + \left(\frac{2GM_h}{v^2 r_t}\right)\right]^{1/2} \tag{4.4}$$

are subject to tidal breakup. In realistic star systems the loss cone is always significantly larger than its purely geometric size and its size is determined by the second term in the brackets in the expression (4.4), which is associated with the focusing of the orbits of stars in the gravitational field of a massive black hole. If $M_h > M_{max}$, when stars located inside the loss cone are swallowed whole by the black hole, in the expression (4.4) and everywhere below the tidal radius r_t must be replaced by the capture radius r_{mb}.

4.2. Rate of tidal breakup of stars

The frequency with which stars are broken up by tidal forces or are swallowed up depends on the distribution and character of their orbits near the massive black hole as well as on the rate at which the loss cone is filled. In particular, collective star interactions in the gravitational field of a black hole under conditions when the loss cone is filled result in the accumulation of stars in the region $r < r_h$. For this reason a kinetic approach must be employed in order to find the star distribution near the black hole and the star flux onto the hole.

The star distribution near a black hole is, generally speaking, nonequilibrium. This is attributable to the flow (draining) of stars into the tidal breakup sphere or below the horizon of the black hole. The form of the nonequilibrium distribution function and the rate of tidal breakup of stars are determined by the interactions of the stars with one another. An adequate approximation for solving this problem is the kinetic equation, derived by Landau,[106] for a system of particles which interact by means of the Coulomb potential. We shall represent the collision integral in this approximation as the divergence of the flux in momentum space

$$I_s = -\frac{\partial}{\partial p^i} j^i, \quad i = 1, 2, 3, \tag{4.5}$$

where in the case of a star system consisting of identical stars each of mass m,

$$j^i = 2\pi(Gm^2)^2\Lambda\int_0^\infty d^3p' w^{ik}\left(\frac{\partial f'}{\partial p'^k}f - \frac{\partial f}{\partial p^k}f'\right), \tag{4.6}$$

$$w^{ik} = \frac{1}{u^3}\left(u^2\delta^{ik} - u^i u^k\right); \tag{4.7}$$

here p^i is the momentum of a star, $u^i = (p'^i - p^i)/m$ is the relative velocity of two stars, δ^{ik} is the Kronecker delta function, and $\Lambda = \ln(N/2)$ is the gravitational Coulomb logarithm. The collision integral (4.5) takes into account only pair interactions of stars, so that the Fokker-Planck equation is the corresponding kinetic equation. The appearance of the Coulomb logarithm Λ in the collision integral reflects the fact that distant encounters of stars, in which the trajectories of the stars change insignificantly, play the dominant role in the interactions. The orbit of a separate star changes significantly over the relaxation time (2.1), which for $N \gg 1$ greatly exceeds the characteristic dynamical time $t_{dyn} = R/v$. For this reason, over the time t_{dyn} the distribution function of the stars $f(\mathbf{r}, \mathbf{p}, t)$, which in the general case is a function of seven variables, can be assumed to depend only on the time and the integrals of motion. In the case of a spherically symmetric stellar system we have $f = f(E, J, t)$, where E is the total energy and J is the angular momentum of a separate star. The change produced in the orbits of the stars by the dominant distant encounters occurs dynamically slowly and has the form of diffusion in (E, J) space.

If $M_h \ll Nm$, all stars in the system can be divided, according to the type of orbit, into infinite and finite, relative to the black hole. Infinite stars are not gravitationally bound with the black hole and move primarily outside the hole's sphere of influence of radius r_h. The distribution and character of the orbits of these stars are virtually independent of the presence of a massive black hole at the center of the star system. Finite stars, however, are bound gravitationally with the black hole and their orbits lie entirely inside the hole's sphere of influence. These stars, which are transferred by interactions to orbits which are gravitationally bound with the hole, form a reservoir for the finite stars.

Because of the existence of a flow of stars onto the black hole the loss cone is filled with stars or, in other words, completely vanishes only outside some kinetic region of radius r_{cr} around the black hole.[107] In the region $r < r_{cr}$ there is no longer enough time for the interactions of stars to fill with new stars the loss cone which was emptied by the tidal breakups or trapping of stars by the hole. The radius of the critical orbit r_{cr}, determined from the equality

$$\theta_{lc}(r_{cr}) = \left(\int_{r_t}^{r_{cr}} \frac{dr}{vT}\right)^{1/2}, \tag{4.8}$$

where the size of the loss cone θ_{lc} is given by Eq. (4.4), can be written in the form

$$r_{cr} \approx b r_h \left(\frac{r_t v_0 T_0}{r_h^2}\right)^q, \tag{4.9}$$

where the numerical constant obeys $b \sim 1$ and the exponent q depends on the star distribution near the black hole. In particular, if $r_h \ll r_{cr} \ll R$, when the loss cone is empty within the radius of influence of the black hole r_h, the exponent is $q = 1/3$.

In the case $r_{cr} \ll R$ there is enough time for the interactions to isotropize the distribution function of the infinite stars within most of the volume of the system. The rate of tidal breakup of such stars is equal to their rate of flow in the direction of the black hole within the loss cone (4.4). For a Maxwellian velocity distribution of infinite stars this rate of flow is equal to[101]

$$\dot{N}_{inf} = (6\pi)^{1/2} n_0 r_t r_h v_0 \simeq 1{,}5\left(\frac{M_h}{m}\right)^{4/3}\left(\frac{v_0}{v_p}\right)^2 / \Lambda T_0, \tag{4.10}$$

where n_0, v_0, and T_0 are, respectively, the average (over the system) density, velocity dispersion, and relaxation time (2.1), employed in order to make it convenient to compare with the rates of other processes. The flow rate of infinite stars (4.10) can be expressed directly in terms of the critical radius r_{cr} as $\dot{N}_{inf} \sim (r_{cr}/R)^3 N/\Lambda T_0$. The condition $r_{cr} \ll R$, which is a necessary condition for the existence of the rate of flow of stars (4.10), can be represented with the help of Eq. (4.9) and the virial theorem (irrespective of the ratio of M_h and Nm) in the form $M_h \ll M_{dif}$, where[108]

$$M_{dif} \approx 0{,}5\left[(v_p/v_0)^2 \Lambda N\right]^{3/4} m. \tag{4.11}$$

In the particular case $M_h \ll Nm$ the quantity M_{dif} depends strongly only on the size R of the star system:

$$M_{dif} \approx 2\left(\frac{R\Lambda}{r_*}\right)^{3/4} m \approx 10^7\left(\frac{\Lambda}{10}\right)^{3/4}\left(\frac{R}{1\ \text{ps}}\right)^{3/4}\left(\frac{r_*}{R_\odot}\right)^{-3/4} m. \tag{4.12}$$

If the mass of the hole is $M_h > M_{dif}$ and therefore we have $r_{cr} > R$, then there is not enough time for the interactions of the stars to maintain an isotropic star distribution and the loss cone is empty within the entire system. The rate of flow of stars onto the hole is now governed by the diffusion of their orbits through the boundary of the empty loss cone. This occurs with the rate

$$\dot{N}_{dif} \approx \frac{1}{\lambda_1}\frac{N}{T_0}, \tag{4.13}$$

where $\lambda_1 = \ln(R/r_t)$, as a result of the change in the angular momenta of the stars in the interactions.[109,110] The diffusion rate (4.13) depends only logarithmically on the size of the loss cone (4.4). For $M_h \gg Nm$ the dependence of the diffusion rate on the mass of the black hole enters in Eq. (4.13) through the relaxation time T_0, since in this case the black hole controls the motion of the stars within the entire star system, while if $M_h \ll Nm$ the diffusion flow depends on the mass of the black hole only logarithmically.

The case $r_{cr} \ll r_h$, when there is enough time for the star distribution to be isotropized within part of the region of gravitational domination of the black hole $r_{cr} < r < r_h$, is less trivial. It turns out that finite stars with an isotropic but nonequilibrium distribution function $f(E) \propto |E|^{1/4}$, where $E = (mv^2/2) - GM_h m/r < 0$ is the total energy of a separate finite star, accumulate in this region.[111,112] This distri-

bution function is the nonlinear isotropic solution of the stationary kinetic equation in the Fokker-Planck approximation and describes diffusion of stars which are gravitationally bound to the black hole toward the black hole. These stars form in the region $r_{cr} < r < r_h$ a self-consistent density distribution

$$n(r) \propto \int_{-GM_h m/r}^{0} f(E)\left[E+\left(GM_h m/r_h\right)\right]^{1/2} dE \propto r^{-7/4}. \quad (4.14)$$

The formation of such a peak, which is of nonlinear origin, in the star density can be considered to be the result of the collective interaction of stars in the gravitational field of the black hole. The finite stars make the main contribution to the total local star density in the region $r < r_h$, since the density of infinite stars in this region varies according to a weaker law $n_{inf}(r) \propto r^{-1/2}$. The density distribution $n(r) \propto r^{-7/4}$ can be easily obtained from the qualitative relation

$$4\pi r^2 \frac{GM_h m}{2r} n(r) \frac{r}{T(r)} \sim \text{const} \quad (4.15)$$

for the steady-state energy flow arising from the black hole when stars which interact with one another accumulate and diffuse in the gravitational field of the black hole.

The typical orbit of a finite star in the region $r_{cr} \ll r \ll r_h$ has the form of trajectory which spirals slowly towards the black hole and along which the binding energy of the star with the hole gradually increases. The energy released in the process by means of the interactions is redistributed among the other stars in the system and ultimately contributes to heating up of the system. However, after the energy of the star drops to the value $E_{cr} = -GM_h m/2r_{cr}$, when the semimajor axis of the orbit of the star is of the order of the critical radius r_{cr}, the character of the diffusion changes. For $E < E_{cr}$, because the loss cone is not completely filled the distribution function of finite stars becomes anisotropic: $f(E,J) \propto |E|^{1/4} \ln(J/J_{min})$, where $J_{min} = m(2GM_h r_t)^{1/2}$. Such stars now approach the region of tidal breakup $r = r_t$ as a result of diffusion of their orbits along not the energy but rather the angular momenta, as a result of which the velocity vectors of the finite stars turn in the direction of the loss cone.[107] As a result, the rate of flow of finite stars onto the black hole, calculated by solving the two-dimensional steady-state Fokker-Planck equation for an anisotropic distribution function $f = f(E, J)$, is equal to[113,114]

$$\dot{N}_{fin} \approx \frac{s}{\lambda_2}\left(\frac{M_h}{Nm}\right)^3 \frac{r_{cr}}{r_h} \cdot \frac{N}{T_0}, \quad (4.16)$$

where $\lambda_2 = \ln(r_{cr}/r_t)$, and the numerical constant is $s \approx 10^3$. Under the condition $r_{cr} \ll r_h$, when there exists a peak of the density $n(r) \propto r^{-7/4}$, the critical radius is given by Eq. (4.9) with the exponent $q = 4/9$. From Eqs. (4.10) and (4.16) we find that $\dot{N}_{fin}/\dot{N}_{inf} \sim (r_h/r_{cr})^{5/4} \gg 1$, and, therefore, when the density has a peak the finite stars predominate in the flow of stars onto the black hole. The expression for the rate of flow of finite stars (4.16) can be represented in the form

$$\dot{N}_{fin} \sim \frac{N(r_{cr})}{\lambda_2 T(r_{cr})}, \quad (4.17)$$

where $N(r_{cr})$ is the total number of stars within the radius r_{cr} and $T(r_{cr})$ is the relaxation time near r_{cr}. This means that over the time $\tau_J \sim \lambda_2 T(r_{cr})$ all stars with energies $E \leqslant E_{cr}$ fall within the region of tidal breakup $r < r_t$ as a result of the diffusion of their orbits along angular momenta and they cease to exist. During this time interval the energy of a separate star can decrease, as a result of the diffusion of its orbit, to $E_{min} \sim \lambda_2 E_{cr}$, differing from E_{cr} by a factor of only several units. The fact that the energy of finite stars is bounded by the quantity E_{min} results in the region $r_t < r \ll r_{cr}$ in a smoother increase of the density: $n(r) \propto r^{-1/2}$.

In a star system with $M_h \gg Nm$ all stars are finite with respect to the black hole and the contribution of self-gravitation of the stars to the integral parameters of the star system can be neglected. In particular, the effective radius of such a star system is $R = GM_h/2v_0^2$, and the local velocities are completely determined by the parameters E and J of the orbits of the stars in the gravitational field of the black hole. As before, however, the interactions of the stars determine the rates of different processes, for example, the rate of tidal breakup of the stars or the rate at which the stars are swallowed up by the hole. Extension of the solution of the two-dimensional Fokker-Planck equation to the case $M_h \gg Nm$ (Refs. 113 and 114) leads for $r_{cr} \ll R$ to the rate of tidal breakup of stars

$$\dot{N}_{fin} \approx \frac{10}{\lambda_2} \cdot \frac{r_{cr}}{R} \cdot \frac{N}{T_0}. \quad (4.18)$$

The critical radius r_{cr} for $r_{cr} \ll R$ and $M_h \gg Nm$ is given by Eq. (4.9), in which r_h must be replaced by R and q must be set equal to 4/9. The star density varies as $n(r) \propto r^{-7/4}$ in the region $r_{cr} \ll r \ll R$ and as $n(r) \propto r^{-1/2}$ in the region $r_t \ll r \ll r_{cr}$.

4.3. The conditions of formation of a peak in the star density around a massive black hole

A peak in the star density $n(r) \propto r^{-7/4}$ is formed around a black hole by stars that are gravitationally bound with the hole, if in the region of the gravitational influence of the hole on the motion of the stars there is enough time for the loss cone to be filled with new finite stars which replace the stars entering the region of tidal breakup $r < r_t$. The necessary condition for this is $r_{cr} \ll r_h$ in the case $M_h \ll Nm$ or $r_{cr} \ll R$ in the case $M_h \gg Nm$. The critical radius r_{cr}, outside which the loss cone is filled with stars, is given by Eq. (4.9) in which the exponent is $q = 4/9$, and the obvious substitution of R for r_h must be made if $M_h \gg Nm$. The density of stars continues to increase toward the hole as $n(r) \propto r^{-7/4}$ only up to the radius $r \sim r_{cr}$. In the region $r \ll r_{cr}$, where the diffusion of the orbits of stars over angular momenta predominates over the diffusion of orbits over energies, the density varies more slowly: $n(r) \propto r^{-1/2}$. It should be noted that the search for massive black holes in nearby galaxies and globular clusters on the basis of the distribution of the density of stars in their central parts, when their surface density is actually recorded, is problematic. The reason lies in the fact that an appreciable peak of the surface density $rn(r) \propto r^{-3/4}$ exists only in the case $r_{cr} \ll r_h$. At the same time, for $r_{cr} > r_h$ the small peak in the density of infinite stars $n(r) \propto r^{-1/2}$ in the surface density distribution is smoothed over and it is virtually unrecordable.

Apart from the condition $r_{cr} \ll \min(r_h, R)$, which is

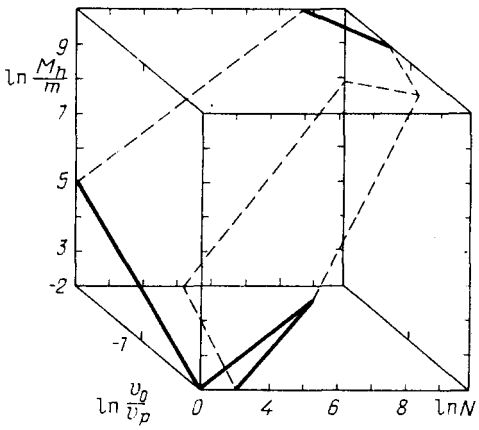

FIG. 2. A peak exists in the star density $n(r) \propto r^{-7/4}$ in systems with $N \gg 1$ stars and velocity dispersion of stars $v_0 \ll v_p = (2Gm/r_*)^{1/2}$ around a massive black hole whose mass falls into the range $M_{fin}(N, v_0) < M_h < N_{dif}(N, v_0)$ with the limits given by Eqs. (4.12) and (4.19).

necessary for the existence of a peak in the density of stars $n(r) \propto r^{-7/4}$, there is an additional limit associated with the effect of disintegrative collisions. Comparing the rate of star collisions (2.35) with the rates of tidal breakup (4.10), (4.13), (4.16), and (4.18) we arrive at the conclusion that for $v_0 > v_p$ collisions become the dominant factor in the change in the number of stars in the system and feeding of accretion onto the central black hole. In particular, they destroy the peak in the star density $n(r) \propto r^{-7/4}$, thereby preventing stars from accumulating in the region of the gravitational influence of the hole $r < r_h$. Therefore the two conditions $r_{cr} \ll \min(r_h, R)$ and $v_0 < v_p$ must be satisfied in order for a peak to exist in the star density $n(r) \propto r^{-7/4}$ around a massive black hole. With the help of the virial theorem and Eq. (4.9) these conditions reduce for $N \gg 1$ to the relations $M_{fin} \ll M_h \ll N_{dif}$, where $M_{fin} \ll Nm$ and is equal to

$$M_{fin} \approx 0{,}1\left[\Lambda^{-1}N^2(v_0/v_p)^2\right]^{3/5} m, \tag{4.19}$$

while the quantity M_{dif} is given by Eq. (4.11). The region of values of the parameters of star systems with a peak in the star density around a central black hole is shown in Fig. 2.

5. JOINT EVOLUTION OF THE NUCLEUS OF A GALAXY AND OF A CENTRAL BLACK HOLE

5.1. Basic evolutionary processes

The formation of a massive black hole at the center of a star cluster complicates and makes more diverse the paths of evolution of the combined system. Additional evolutionary factors appear, such as, a decrease in the number of stars in the system and simultaneous increase in the mass of the black hole as the hole absorbs these stars or part of their matter. In addition, the total energy of the stars remaining in the system changes as stars from the density peak diffuse toward the hole. The mass of the black hole ultimately stops increasing when the gravitational influence of the hole becomes dominant within the entire star system.

The physically distinguishable stages of the joint evolution of a star system and a central massive black hole are characterized by domination of one of the processes described by Eqs. (2.3), (2.35), (4.10), (4.13), and (4.16) or (4.18): (1) evaporation of stars out of the system, (2) tidal breakup of stars which are not gravitationally bound with the hole (infinite stars) with a filled loss cone, (3) diffusion of stars toward the hole under the conditions of an empty loss cone, (4) diffusion toward the hole of stars which are gravitationally bound with the hole in the presence of a peak in the density and, finally, (5) star collisions. The ratios of the rates of these processes depend on the mass of the black hole and one or two virial parameters, for example, the velocity dispersion v_0 of the stars in the system and the total number of stars N. A significant fraction of the gas freed by tidal breakup of the stars by the hole and disruptive collisions of stars at the later stages of evolution is accreted onto the black hole. For this reason it is natural to suppose that the rate of accretion on the black hole is close to the rate of accumulation of gas in the system, if the latter rate does not exceed the value necessary for maintaining the luminosity of the black hole nearing the Eddington limit $L_E = M_h c^2/t_E$, where $t_E = \sigma_T c/4\pi G m_p \sim 5\cdot 10^8$ yr. In the opposite case the luminosity of the black hole will be maintained at the maximum level $L \simeq L_E$ and the mass of the black hole will grow exponentially with the characteristic time $\tau = t_T/\eta$, where $\eta \sim 0.1$ is the accretion efficiency. A strong stellar wind could also be an additional source of accreted gas. Such a wind arises as a result of the heating of the outer layers of the stars by the intense radiation from the central source[115-120] and partially blows out of the system. It is natural to study the possible stages of the joint evolution of the composite system starting with relatively small masses of the black hole, in the limit when $M_h \ll Nm$ and the velocity dispersion of the stars $v_0 \ll v_p$, i.e., stellar collisions play a small role. The general sequence of change in the regimes of joint evolution of a star system and a central black hole is shown in Fig. 3 and will be justified in subsequent sections.

5.2. Increase of the mass of the black hole at the stage of evaporation of stars from the system ($r_h \ll r_{cr} \ll R$, $v_0 \ll v_p$, $N_{inf} \ll \dot{N}_{ev}$)

Let the mass of the central black hole be so small that the conditions for filling of the loss cone with stars and absence of a peak in the density of finite stars are satisfied: $v_0 \ll v_p$ and $r_h \ll r_{cr} \ll R$; according to Eqs. (4.11) and (4.19), these conditions have the form $M_h \ll \min(M_{dif}, M_{fin})$. In such a star system the mass of the black hole increases as a result of accretion of gas freed in the process of tidal breakup of stars which are not gravitationally bound with the hole. Let the condition $\dot{N}_{inf} \ll \dot{N}_{ev}$ also be satisfied. With the help of Eqs. (2.3) and (4.10) this condition can be written in the form $M_h \ll \alpha^{3/4} M_{dif}$, where $\alpha \simeq 7\cdot 10^{-3}$, so that the decrease in the total mass of the stars is associated primarily with the evaporation of stars from the system. Under such conditions the star system evolves independently of the black hole according to the classical evaporation scenario (2.4), when $R \propto N^2$ and $v_0 \propto N^{-1/2}$. Integrating Eq. (4.10) with the help of Eq. (2.4), we obtain the law of growth of a black hole at the evaporation stage of evolution[121]

$$\mu = \left\{1 + \mu_1\left[1 - \left(1 - t\tau_{ev}^{-1}\right)^{-2/7}\right]\right\}^{-3}, \tag{5.1}$$

where $\mu = M_h(t)/M_h(0)$, $\mu_1 = (7/6)\tau_{ev}\dot{\mu}(0)$. The duration of this stage is determined by the moment at which rapid

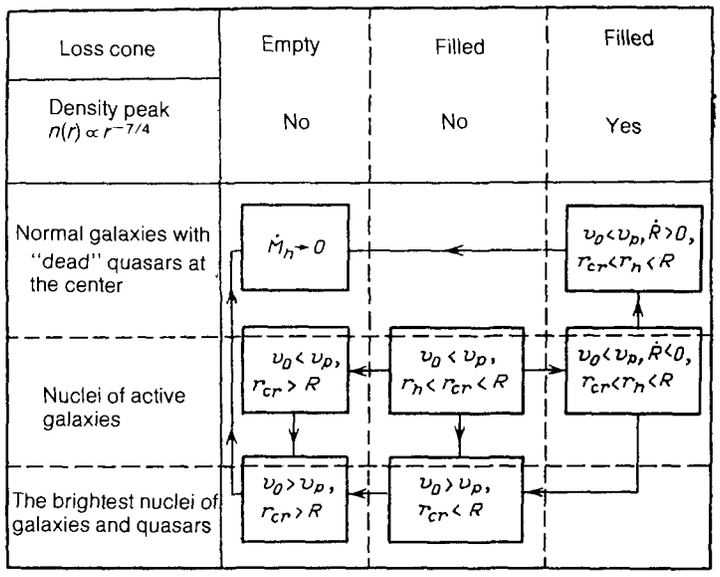

FIG. 3. The sequence of change in the physically different stages of evolution of galactic nuclei with a central massive black hole, depending on the relative values of the velocity dispersion v_0 of the stars and the escape speed v_p from the surface of the star and the relative values of the radius of influence of the hole $r_h = GM_h/v_0^2$, the critical radius r_{cr}, within which there is not enough time for the loss cone to be filled with stars, and the radius of the star system $R = GmN/2v_0^2$. The dashed lines show the conventional boundaries separating the nuclei of galaxies exhibiting different degrees of activity. A detailed analysis of each stage of evolution is given in Secs. 5.2–5.6.

growth of the mass of the black hole starts, which according to Eq. (5.1) is equal to

$$\tau_h = \left[1 - \left(\frac{\mu_1}{1+\mu_1}\right)^{7/2}\right]\tau_{ev}. \tag{5.2}$$

By this time the number of stars remaining in the system $N_h = \mu_1(1+\mu_1)^{-1}N(0)$ and therefore the maximum mass which the black hole can acquire at subsequent stages is equal to $N_h m$. According to Eq. (5.1), over the contraction time of the system τ_{ev} the mass of the black hole increases weakly if $\mu_1 < 1$. This condition can be represented in the form $M_h < M_{min}$, where

$$M_{min} < \left[\alpha\Lambda\left(\frac{v_p}{v_0}\right)^2\right]^3 m. \tag{5.3}$$

If, however, $M_h > M_{min}$, then as the system contracts the mass of the black hole increases significantly and after reaching the value $M_h = \alpha^{3/4}M_{dif}$ the rate of tidal breakup of infinite stars $\dot{N}_{inf}$ starts to exceed the rate of evaporation of stars $\dot{N}_{ev}$ from the system.

5.3. Rapid growth of a black hole at the stage without a peak in the density ($r_h \ll r_{cr} \ll R$, $v_0 \ll v_p$, $\dot{N}_{inf} \gg \dot{N}_{ev}$)

At this stage there is no peak in the density of finite stars around a black hole $n(r) \propto r^{-7/4}$.[122] The mass of the black hole satisfies the relations $\alpha^{3/4}M_{dif} \ll M_h \ll \min(M_{dif}, M_{fin})$, where M_{dif} and M_{fin} are given by Eqs. (4.11) and (4.19). The rate of flow of infinite stars toward the black hole (4.10) determines the growth in the mass of the black hole and the change in the number of stars in the system $\dot{N} = -\dot{N}_{inf}$. If $\dot{N}_{inf} \gg \dot{N}_{ev}$, the characteristic evolution time of the system is $\tau_{inf} = N(0)/\dot{N}_{inf}(0) \ll \tau_{ev}$. In this limit the decrease in the total mass of the system $M = Nm + M_h$ over the time τ_{inf} is relatively small and it can be neglected in the evolutionary equation (2.2). In the meantime, the tidal breakup and absorption of infinite stars by the hole is not accompanied by a change in the total energy of the system E and therefore neglecting star evaporation, $M = Nm + M_h = \text{const}$ and $\dot{R} = \dot{v}_0 = 0$. The only variable parameters are the number of stars in the system N and the mass of the hole M_h. As a result, the system evolves according to the law

$$\mu = 1 + \chi(1-\nu), \tag{5.4}$$

$$t = \tau_{inf}\int_\nu^1 \frac{dx}{x[1+\chi(1-x)]^{4/3}}, \tag{5.5}$$

where $\mu = M_h(t)/M_h(0)$, $\nu = N(t)/N(0)$, $\chi = N(0)m/M_h(0)$. Under conditions of rapid growth of the black hole, which is equivalent to relatively slow change in the parameters of the star system, we obtain[108] the following approximate law of growth of the mass of the black hole with time:

$$\mu \approx \left[1 - \left(\frac{\chi}{3}\right)\frac{t}{\tau_{inf}}\right]^{-3}, \tag{5.6}$$

which also follows from Eq. (5.1) with $t \ll \tau_{ev}$. Over the lifetime of the nucleus of a galaxy, comparable to the age of the universe, the mass of the central black hole will grow significantly if $(3/\chi)\tau_{inf} \leqslant 10^{10}$ yr. This condition can be written in the form $M_h \gg M_{int}$, where

$$M_{inf} \approx \left[\left(\frac{\pi}{6}\right)^{1/2}\left(\frac{R}{r_*}\right)^2\frac{v_p}{v_0}\cdot\frac{r_*}{v_p}\cdot(10^{10}\text{a})^{-1}\right]^3 m. \tag{5.7}$$

Here and everywhere below, unless otherwise specified, the numerical value for clusters consisting of stars of solar type is given. Significant growth in the mass of the black hole at this stage of evolution is accompanied by an increase in the velocity dispersion of the stars v_0 and a decrease of the ratios r_{cr}/r_h and R/r_{cr}. Ultimately the system passes to the stage of evolution when in the system either a peak forms in the density $n(r) \propto r^{-7/4}$ for $r_{cr} \ll r_h$ or the loss cone is empty everywhere in the system for $r_{cr} > R$ or, finally, collisions of stars with $v_0 > v_p$ begin to dominate.

5.4. Heating of the star system by the peak in the star density around a black hole ($r_{cr} \ll \min(r_h, R)$, $v_0 \ll v_p$)

In order for a peak to form in the density of finite stars $n(r) \propto r^{-7/4}$ in the case $M_h \ll NM$, apart from the condition $r_{cr} \ll r_h$, which is satisfied when $M_h \gg M_{fin}$ from Eq. (4.19),

the additional requirement $v_0 \ll v_p$, when collisions between stars are relatively rare, must also be satisfied. The accumulation of stars which are gravitationally bound with the black hole can qualitatively change the fate of the star system, converting the evolutionary contraction of the system into expansion. This happens when the system is heated quite strongly by the stars forming the peak in the density.[123] As they diffuse toward the black hole their total energy decreases and, correspondingly, the energy of the stars remaining in the system increases. Each star precipitating from the peak of the density onto the hole transfers energy $|E_{\min}| \simeq \lambda_2 GM_h m/2r_{cr}$, where $\lambda_2 = \ln(r_{cr}/r_t)$, to heating of the star system. The rate of heating of the system by the flow of finite stars (4.16) does not depend on r_{cr} and is equal to

$$\dot{E}_{\text{fin}} = \frac{s}{2}\cdot\left(\frac{M_h}{Nm}\right)^3 \frac{|E|}{T_0}, \tag{5.8}$$

where the numerical constant is $s \simeq 10^3$ and the total energy is $E = -Nmv_0^2/2$. From a comparison of the rates (2.3), (4.10), (4.16), and (5.8) it follows that for $M_h \gg M_{\text{fin}}$, when $r_{cr} \ll r_h$, the rate of heating $\dot{E}_{\text{fin}}$ and the rate of change of mass of the system owing to star evaporation $\dot{N}_{ev}$ make the dominant contributions to the right-hand side of the evolutionary equation (2.2). The remaining contributions can be neglected. Substituting Eqs. (2.3) and (5.8) into Eq. (2.2) we obtain the condition of expansion of the star system $\dot{R} > 0$ in the form $M_h > M_{\exp}$, where

$$M_{\exp} = \left(\frac{2\alpha}{s}\right)^{1/3} Nm. \tag{5.9}$$

Thus a black hole whose mass is equal to several percent of the total mass of the star system is capable of supporting secular expansion of the system. At the same time, expansion will be accompanied by slowing down or even complete stopping of the growth of the black hole. The simultaneous solution of Eqs. (22), (23), and (5.8) with a frozen black hole for $N_{ev} \gg N_{\inf}$ has the form[121]

$$\rho = \nu^2 \exp\left[\frac{2}{3}\cdot\nu_0^3(\nu^{-3} - 1)\right], \tag{5.10}$$

$$t = \alpha^{-1}T(0)\int_\nu^1 x^{5/2}\exp\left[\nu_0^3(x^{-3} - 1)\right]dx, \tag{5.11}$$

where $\rho = R(t)/R(0)$, $\nu = N(t)/N(0)$, and

$$\nu_0 = \left(\frac{s}{2\alpha}\right)^{1/3} \frac{M_h(0)}{mN(0)}. \tag{5.12}$$

For $M_h = 0$ this solution reduces to the classical evaporation evolution of the system (2.4). The condition for the system to expand initially will be $\nu_0 > 1$, which is identical to the condition $M_h > M_{\exp}$. For $\nu_0 \gg 1$ the expansion proceeds according to the simpler law

$$\rho = \left[1 + \left(\frac{t}{\tau_{\exp}}\right)\right]^{2/3}, \tag{5.13}$$

where $\tau_{\exp} = (3\nu_0^3\alpha)^{-1}T_0(0) < 45T_0(0)$. Then the number N of stars in the system decreases with time logarithmically and all dynamical processes slow down. When the number of stars increases slowly, the mass of the black hole grows as

$$\mu = \left[1 - \mu_2\left(1 - \rho\right)^{-4/9}\right]^{-27/47}, \tag{5.14}$$

where $\mu = M_h(t)/M(0)$, $\mu_2 = (47/4)\dot{\mu}(0)\tau_{\exp}$. From here we find the self-consistent condition of weak growth of a black hole in the form $\mu_2 \ll 1$ or $M_h \gg M_{fr}$, where

$$M_{fr} \approx 0{,}1\lambda_2^{-27/47}\Lambda^{-12/47}N^{51/47}(v_0/v_p)^{24/47}m. \tag{5.15}$$

For $t \gg \tau_{\exp}$ the accretional luminosity of a frozen black hole decreases as $L \propto t^{-35/27}$.

The initial stage of expansion of the system can be realized under conditions when $\dot{N}_{ev} \ll \dot{N}_{\inf}$. In this limit the dimensionless mass of the black hole is $\mu = 1 + \chi(1 - \nu)$ and the system expands according to the law

$$\rho = \left\{1 - \xi\int_\nu^1 \frac{dx}{x^{5/9}}\left[1 + \chi\left(1 - x\right)\right]^{20/27}\right\}^{-9/4}, \tag{5.16}$$

$$t = \tau_{\text{fin}}\int_\nu^1 \frac{dx\, x^{-14/9}}{[1 + \chi(1 - x)]^{61/27}} \times \left\{1 - \xi\int_x^1 \frac{dz}{z^{5/9}}\left[1 + \chi(1 - x)\right]^{20/27}\right\}^{-35/8}, \tag{5.17}$$

where $\xi = (2/9)\lambda_2[r_h(0)/r_{cr}(0)] \gg 1$ and $\tau_{\text{fin}} = N(0)/\dot{N}_{\text{fin}}(0)$. The radius of the star system as well as the total evolution time, as follows from Eqs. (5.16) and (5.17), formally approach infinity, when a finite number of stars remains in the system. Correspondingly the growth of the black hole stops. After the system freezes it expands according to Eq. (5.13), and the number of stars decreases as

$$\nu \simeq 1 - \xi^{-1}\left(1 - \rho^{-4/9}\right), \tag{5.18}$$

and in addition the characteristic decay time of the number of stars in the system at this stage is $\tau_{\text{fin}} = (27/8)\xi\tau_{\exp} \gg \tau_{\exp}$. The condition of self-consistency of Eqs. (5.16) and (5.17) with the requirement that the black hole be frozen is identical to the condition $M_h \gg M_{fr}$ from Eq. (5.15).

In a star system with a gravitationally dominant black hole with mass $M_h \gg Nm$ a peak forms in the star density $n(r) \propto r^{-7/4}$ if $v_0 \ll v_p$ and $r_{cr} \ll R$. The last inequality is satisfied when $M_h \ll M_{\text{dif}}$ from Eq. (4.11). The rate of flow of stars out of the density peak toward the black hole (4.18) in this limit is accompanied by heating of the star system at the rate $\dot{E} = \lambda_2(R/r_{cr})|E|\dot{N}_{\text{fin}}/N$ and an unavoidable expansion of the system according to the law[124]

$$\rho = \left[1 + \left(\frac{t}{\tau^*_{\exp}}\right)\right]^{2/3}, \tag{5.19}$$

where $\tau^*_{\exp} = T_0(0)/15$. For $\dot{N}_{ev} \ll \dot{N}_{\inf}$ the number of stars in the expanding star system changes as

$$\nu = \left[1 - \zeta^{-1}\left(1 - \rho^{-4/9}\right)\right]^{-9/4}, \tag{5.20}$$

where $\zeta = \lambda_2[R(0)/r_{cr}(0)] \gg 1$, while for $\dot{N}_{ev} \gg \dot{N}_{\inf}$ the number of stars changes as $\nu = \rho^{-\alpha/10}$, where $\alpha \approx 7\cdot 10^{-3}$. For $t \gg \tau^*_{\exp}$ the luminosity of a gravitationally dominant black hole decays, analogously to the case of a frozen black hole with mass $M_h \ll Nm$, $L \propto t^{-35/27}$ or even more rapidly, if the quite compact constituent stars of the system are swal-

lowed whole by the hole without tidal breakup outside the event horizon.

As the star system expands the ratios v_0/v_p, r_{cr}/r_h, and r_{cr}/R decrease. For this reason, the conditions of existence of a peak in the star density, which governs the heating and expansion of the system, are still satisfied. However the ratio $\dot{N}_{ev}/\dot{N}_{inf}$ increases and ultimately there occurs a transition to the state of expansion with $\dot{N}_{ev} \gg \dot{N}_{inf}$. Such evolution will end in a dead quasar at the center of a frozen star system (Fig. 3).

5.5. Evolution with an unfilled loss cone ($r_{cr} \gg R$, $v_0 \ll v_p$)

A star system passes into this regime when $M_h \gg M_{dif}$ from Eq. (4.11) in the process of contraction at the stage without a density peak (see Secs. 5.2 and 5.3) or with a density peak (see Sec. 5.4) incapable of stopping the contraction. Tidal breakup or capture of stars by the hole now occurs as a result of diffusion of the orbits of the stars into an empty (within the entire system) loss cone at the rate (4.13), which is significantly higher than the rate (2.3) at which stars evaporate from the system. For this reason, neglecting star evaporation, just as in the analysis of the stage of rapid growth of a black hole in the regime without a density peak (see Sec. 5.3), we have $M = M_h + Nm = \text{const}$, $\dot{R} = \dot{v}_0 = 0$, $\mu = 1 + \chi(1 - \nu)$ with $\chi = N(0)/M_h(0)$. In this case the integration of Eq. (4.13) with $\dot{N} = -\dot{N}_{dif}$, irrespective of the ratio of M_h and Nm, gives a single law governing the temporal evolution of the system

$$\nu \simeq \left[1 + \left(\frac{t}{\tau_{dif}}\right)\right]^{-1}, \tag{5.21}$$

where $\tau_{dif} = \lambda_1 T_0(0)$. Star evaporation is nonetheless accompanied by relatively slow contraction $\rho = \nu^k$ in the case $M_h \ll Nm$ and $\rho = \nu^{k/2}$ in the case $M_h \ll Nm$, where $k = 2\alpha\lambda_1/3 \sim 10^{-1}$, and by slow growth $v_0 \propto \nu^{-k/4}$. As the number of stars in the system decreases the ratio r_{cr}/R increases and therefore the conditions of formation of a peak in the star density which is capable of suppressing contraction are not satisfied. Conversely, as the velocity dispersion v_0 of the stars increases the role of disruptive collisions of stars increases; the rate of such collisions starts to exceed the rate of diffusion of stars toward the hole when $v_0 > v_p = (2\ Gm/r_*)^{1/2}$.

5.6. Stage of collisions between stars ($v_0 > v_p$)

Only star systems that have a peak in the star density $n(r) \propto r^{-7/4}$, and in which heating by stars gravitationally bound to the black hole gives rise to secular expansion of the system (see Sec. 5.4), as well as star systems which consist entirely of neutron stars and black holes of star type masses for which $v_p \sim c$, escape this stage of evolution. Comparing the rates of the evolutionary processes (2.3), (2.35), (4.10), (4.13), (4.16), and (4.18) shows that collisions between stars become the dominant factor in the evolution when $v_0 > v_p$. The peak in the star density $n(r) \propto r^{-7/4}$, if it is incapable of stopping the increase in the velocity dispersion v_0 of the stars, also breaks up at the collisional stage of evolution. When collisions predominate the contraction of the star system with a central black hole is controlled by the same laws as in the case, studied in Sec. 2.7, without a black hole, since collisions do not change the total mass of the system $M = M_h + Nm$ and they predominate over all other processes. Irrespective of the ratio of the masses M_h and Nm the effective radius R of the system decreases according to the approximate law (2.36) as a result of energy dissipation, and the duration of the collisional stage of evolution $\tau_{coll} \sim (v_p/v_0)^4 \Lambda T_0(0)$ depends on the mass of the black hole only in the case $M_h > Nm$.

At the collisional stage of evolution the characteristics of the growth of the mass of a black hole are determined only by the ratio of the critical radius r_{cr}, within which there is not enough time for stars to fill the loss cone, and the radius of the system R. When $r_{cr} \ll R$, infinite stars, filling the loss cone, dominate in the flow of stars toward the hole, and the mass of the black hole, taking into account Eq. (4.10), grows as

$$\mu = \left[1 - \mu_{inf}(1 - \rho)\right]^{-3}, \tag{5.22}$$

where $\mu_{inf} = (\chi/3)\dot{N}_{inf}(0)/\dot{N}_{coll}(0)$, $\chi = mN(0)/M_h(0)$. However this regime of evolution is accompanied by an increase of the ratio r_{cr}/R as the star system contracts and the black hole grows and the regime unavoidably ends with emptying of the loss cone everywhere in the system when $r_{cr} \gg R$. Further growth of the black hole (see Fig. 3) now proceeds in the stable evolutionary regime when $r_{cr} > R$, and the rate of flow of stars toward the hole is determined by the rate of diffusion (4.13) into the empty loss cone. This results in growth of the black hole according to the law[424]

$$\mu = 1 + \mu_{dif}(1 - \rho^2), \tag{5.23}$$

where $\mu_{dif} = (\chi/2)\dot{N}_{dif}(0)/\dot{N}_{coll}(0)$.

The collisional stage of the evolution of a star system surrounding a massive black hole is the most striking phase of the activity of galactic nuclei: in this phase the luminosity of the black hole can be close to the Eddington luminosity. The physical conditions at the collisional stage correspond most closely to regions in which wide emission lines are formed in the nuclei of active galaxies. These regions can be interpreted as a collection of high-velocity clouds of gas in the intense field of radiation from the central source.[125,126] In the most powerful nuclei of active galaxies and quasars with a massive central source $M_h > M_{max}$ from Eq. (4.3), when stars are swallowed whole by the black hole without tidal breakup, it is precisely the star collisions that are capable of supplying excess accreted material to the black hole and maintaining the luminosity of the hole at a high level. In addition, the total intensity of energy release accompanying collisional self-breakup of a star system with a sufficiently high velocity dispersion v_0 of the stars can exceed the Eddington limit $L_E = Mc^2/t_E$, where $t_E = \sigma_T c/4\pi G m_p \sim 5 \cdot 10^8$ yr, for the total mass of the system $M = M_h + Nm$ during the time interval $t \ll (v_0/c)^2 t_E$. Most of the released energy comes from the kinetic energy of the clouds of gas of the disrupted stars. As the clouds expand and interact with one another they form a single spherical shock wave emanating from the nucleus of the galaxy. Some of the energy of this shock wave should be transformed into radiation emitted from an extended region of formation of narrow emission lines[126,127] and some of the energy should go into acceleration of cosmic rays.[128,129]

The stage of collisional self-disruption of dense star

clusters in galactic nuclei, which is accompanied by the generation of a gigantic spherical shock wave, can occur at least twice—prior to the formation of a massive black hole (see Sec. 2.6) and in the process of contraction of the remnant of the star system around the central black hole.

A star system ceases to exist as a cluster of stars when most of the stars break up as a result of collisions. By this time the mass of the black hole grows, according to Eq. (5.23), up to the size $M_h \approx (1/2)N(0)m\dot{N}_{dif}(0)/\dot{N}_{coll}(0) < N(0)m$ or an even smaller size, if the luminosity of the black hole is maintained at the Eddington limit and there is not enough time for all of the gas supplied to the hole to accrete onto the hole. Therefore, at the collisional stage of evolution, there is not enough time for the black hole to absorb the main fraction of the mass of the expanding star system and the evolution of the star system stops with the formation of a gas shell, containing an admixture of surviving stars, around the black hole. At the subsequent gas-dynamic stage of evolution,[83-90] depending on the angular momentum of the gas remnant of the expanding star system, there forms a massive accretion disk which is capable of spinning up the black hole[130] or there is realized a regime of spherical accretion with the rate[2]

$$\dot{M}_h = \alpha(\gamma) r_g^2 c \rho_\infty (a_\infty/c)^{-3}, \tag{5.24}$$

where the numerical constant $\alpha(\gamma)$, which depends on the ratio of the heat capacities, is equal to ~ 1, and the density ρ_∞ and velocity of sound in the gas a_∞ at large distances from the black hole correspond to their values near the radius of influence of the hole $r_h = GM_h/a_\infty^2$. The fragmentation, accompanied by star formation in the gas cloud surrounding the black hole after some of the mass is ejected in the form of a spherical shock wave, partially regenerates the star system, but only under extreme conditions of a deep potential well and high accretional luminosity of the black hole.

6. CENTRAL SOURCE

6.1. Basic principles of operation

From the qualitative point of view there is no doubt that a massive black hole can function as a central "machine" in the nuclei of active galaxies and quasars. This is attributable to the uniqueness of the properties of the gravitational field of the black hole: This field is capable of converting into radiation a significant fraction of the mass of the matter which falls on the black hole. But the detailed picture of the interaction of a black hole with its environment is extremely complicated and confused. At the present time the development of this picture is far from complete, in spite of the significant efforts made by numerous investigators. The fundamental features of the operation of the central machine have, apparently, nonetheless already been determined.[1-7,31-34,43,131-138] These features are as follows:

a) *Accretion*. The external source of energy is undoubtedly the external gas accreted onto the black hole. Since a black hole does not have a solid surface, energy release is most efficient under conditions of disk accretion of gas having quite high angular momentum.

b) *Relativism*. In the strong gravitational field of a black hole the accreted-mass-to-radiation conversion efficiency reaches $\eta = L/\dot{M}_h c^2 \sim 0.1$, which provides high radiation power for a long period of time.

c) *Electrodynamics*. A magnetosphere with a large-scale regular magnetic field, external with respect to the black hole, is formed in the accreted plasma by means of the dynamo mechanism. As a result the entire system operates as an electromagnetic generator.

d) *Rotation*. Rarefied funnels (nozzles) form in the nonstationary region along the common axis of rotation of the black hole and the accretion disk. Oppositely directed relativistic jets are formed in these funnels under the action of the gradients of the gas and radiation pressures as well as the action of the magnetospheric electromagnetic field.

Finding a detailed solution of the problem of accretion onto a black hole is a challenge for modern astrophysics, since this requires a self-consistent analysis of gravitational, gas-dynamic, and electromagnetic factors in a wide range of conditions, encompassing the real diversity of phenomena occurring in the nuclei of active galaxies. At the same time, the diversity of possible conditions of accretion makes it possible to single out limiting regimes that can be described with only a small number of parameters. In the Newtonian region, at large distances from the region of main energy release, the balance of the gradient of the pressure and the forces of gravitation in the rotating accretion flow results in the formation of a geometrically thin disk.[34,131-134] Owing to viscosity, the differential rotation of the layers in turn gives rise to transfer of the excess angular momentum from the interior regions of the disk to the periphery. The energy that is released in the process is radiated out locally from the surface of the disk. When the luminosity is high the outer layers forming the atmosphere of the disc are subjected to additional strong heating owing to absorption of some of the energy emitted from the central relativistic region of the disk. In the process of disk accretion the black hole is spun up by the matter incident on it to a state close to extremely rapid rotation; this is accompanied by an additional increase of the accretion efficiency up to $\eta \simeq 0.3$.[140,141] In the relativistic region, where the main energy release occurs, under conditions of a high rate of accretion the thin disk expands as a result of the development of thermal instability[142-149] and becomes geometrically thick. In a thick accretion disk surfaces of equal density (isodenses) have the shape of coaxial tori;[97-99] this, in particular, gives rise to additional collimation of the radiation along the rotation axis of the black hole. Steady-state accretion on a black hole is possible, generally speaking, only if the total luminosity of the central source is $L \leqslant L_E$, when the gravitational attraction of the black hole can withstand the radiation pressure exerted on the plasma. For luminosity close to the Eddington limit, an accretion disk will be surrounded by an extended hot chromosphere, formed by plasma flowing out of the disk.[150-152] In the case of a quite dense plasma a super-Eddington radiation flux could lead not to the expulsion of plasma out of the vicinity of the central source, but rather to the development of strong turbulence in this region and transition to convective energy transfer, as happens in the deep interior of massive stars.

6.2. Characteristics of the magnetosphere of a black hole

Because of contraction of the plasma and the dynamo mechanism, intensification of the starting magnetic field frozen into the plasma is unavoidable in the accretion flow onto the black hole.[130,153-156] Quite close to the black hole the

electromagnetic forces determine, on an equal footing with gravitation, the character of the distribution and motion of charged particles.[43,157–159] Because of the Bardeen-Peterson effect,[160] in the region of the strong gravitational field of a rapidly rotating black hole an accretion disk with an arbitrary orientation of the angular momentum lies in the plane of the equator of the hole and correspondingly the magnetosphere of a rapidly rotating black hole has an axisymmetric structure in the relativistic region. The event horizon of a black hole is electromagnetically qualitatively similar to a sphere with finite conductivity,[43,161,162] so that the magnetospheres of black holes are in many respects similar to the magnetospheres of rotating neutron stars—pulsars.[163,164] At the same time there is a significant difference owing to the fact that the seed source of the magnetosphere of a neutron star is the star's characteristic internal magnetic field, to which the accreted plasma or the plasma generated by cascade electron-positron pair production becomes adjusted. Conversely, in the more complicated case of the magnetosphere of a black hole all seed sources of electromagnetic energy lie in the accreted plasma itself and are external relative to the black hole.[165] The magnetosphere, in turn, exerts a back electromagnetic effect on the black hole, causing the hole to become polarized or leading to accumulation of electric charge inside the hole. In particular,[166] in the vacuum limit the equilibrium electric charge of a black hole with angular momentum $J = aM_h G/c$ in a collinear magnetic field of intensity B_0 is equal to $q_0 = 2B_0 JG/c^3$. This charge, taking into account the back effect of the magnetic field on the metric,[167] is conserved by the hole, even for an extreme value of the rotation parameter of the hole $a = M_h$.

In the general case, when a large-scale magnetic field of intensity B_0 is generated in the magnetosphere owing to the polarization of the hole and the accreted plasma, there is induced in addition near the event horizon of the rotating black hole an electric field of intensity $\mathscr{E} \sim (a/M_h)B_0$, where the rotation parameter of a Kerr black hole is $|a| \leqslant M_h$. The electromagnetic equilibrium of the rotating magnetosphere and the black hole is maintained by the polarization of the hole and the plasma, as well as by the currents flowing in the plasma, and in addition the black hole itself is part of the electric circuit. Such an electric battery operates owing to the emission of electromagnetic waves and current losses.[43] The latter apparently are the dominant mechanism of energy release, since owing to the polarization the electromagnetic radiation can be suppressed in a dense plasma, as happens in the magnetosphere of pulsars.[164] In order of magnitude, the total power of energy release in the magnetosphere of a black hole reaches

$$W \sim 10^{-2}\left(\frac{a}{M}\right)^2 B_0^2 r_g^2 c$$

$$\sim 10^{45}\left(\frac{a}{M}\right)^2\left(\frac{B_0}{10^4\,\mathrm{G}}\right)^2\left(\frac{M_h}{10^9 M_\odot}\right)^2 \mathrm{erg\cdot s^{-1}}, \qquad (6.1)$$

with the maximum emf in such a battery

$$U \sim \frac{a}{M} B_0 r_g \sim 4{,}4\cdot 10^{20}\frac{a}{M}\cdot\frac{B_0}{10^4\,\mathrm{G}}\left(\frac{M_h}{10^9 M_\odot}\right) V. \qquad (6.2)$$

The hopes for acceleration of particles in the vacuum funnels along the rotation axis of the black hole are pinned on this electric potential. The maximum intensity of the magnetic field in the magnetosphere of a black hole can be estimated, for example, from the condition that the energy densities of the accreted plasma and the electromagnetic field are equally distributed; this leads to the limit

$$B_0 \leq \left(\frac{m_p c^2}{\eta r_g \sigma_T}\right)^{1/2} \simeq 4\cdot 10^4\left(\frac{\eta}{0{,}1}\right)^{-1/2}\left(\frac{M_h}{10^8 M_\odot}\right)^{-1/2} \mathrm{G}. \qquad (6.3)$$

6.3. Generation of cosmic rays

The gravitational field of a black hole is by itself, without the magnetospheric electromagnetic field (which is of secondary origin) a powerful accelerator, since the velocities of all particles falling on the hole approach the velocity of light relative to a locally stationary observer as the event horizon of the black hole is approached. This gravitational accelerator will operate efficiently not only with respect to itself, but also with respect to an external observer, if conditions for emission of the products of interaction of the accelerated particles are satisfied, which in turn requires appropriate targets in the region of the steady-state orbits. Under real conditions the targets are the particles of the accreted gas itself with relativistic relative velocities of the particles of the gas. This possibility is realized at least under conditions of quasispherical accretion of rarefied gas or near the inner boundary of a thick disk. In a thin disk, however, in spite of the relativistic character of the volume flow of gas, the particles in the gas, even near the inner boundary of the disc, move only with thermal, generally speaking, nonrelativistic velocities. A rapidly rotating black hole has additional advantages for generation of fast particles. In the region adjoining the event horizon of a rotating black hole and called the ergosphere it is possible for Penrose's process,[168] qualitatively analogous to the loss of energy by a rotating body interacting with an external field, to occur.[169,170] Aside from the usual precession, strong precession of the instantaneous planes of nonequatorial orbits around the rotation axis of the hole—Wilkins effect[171]—occurs in the gravitational field of a rotating black hole near the event horizon of the hole. As a result, in the relativistic region the close-to-circular nonequatorial trajectories rapidly intermix and near the plane of the equator of the black hole the particles of accreted gas move along opposite trajectories with low radial but relativistic relative velocities. Such trajectories of particles of the accreted gas are typical for the inner boundary of a thick disk. The mechanism of generation of fast particles in collisions and decays of particles of the accreted gas operates most efficiently near the boundary of stable motion,[140] and the particles which escape most rapidly from the black hole are produced in Penrose processes. For example,[172] the elastic interaction of two protons on converging nonequatorial orbits, which initially were almost at rest far from the black hole, in the ergosphere of an extremely rapidly rotating black hole can result in one of the particles escaping to infinity with maximum kinetic energy $2(2/3)^{1/2} m_p c^2 \approx 1.5$ GeV. In similar inelastic interactions accompanied by the production of γ-rays the upper limit on the energy of a γ-ray at infinity is $2m_p c^2 \approx 1.9$ GeV. These energies correspond to

the optimal geometry of interaction and escape. Such a geometry is unlikely under conditions of statistical averaging. The thickness of the matter traversed by the fast particle in the quasispherical flow or in the thick disk near its inner boundary is $x \sim (L/L_E)(m_p/\eta\sigma_T) \sim 25(L/L_E)(\eta/0.1)^{-1}$ $g \cdot cm^{-2}$, where η is the matter-to-radiation conversion efficiency and σ_T is the Thompson cross section. For luminosity near the Eddington limit $L \sim L_E$ and $\eta \sim 0.1$ the accumulated mass of matter is sufficient for efficient nuclear interactions, whose typical cross sections in the subrelativistic regime are $\sigma_{pp} \sim 10^{-26}\,cm^2$. The final products of p–p interactions near a black hole will be γ-rays and neutrinos with characteristic energies at infinity of the order of 10–100 MeV. The total intensity of their radiation depends on the character of the proton trajectories in the relativistic region of accretion near a black hole.[173,174] The idealized model calculation[175] for the case of a nonrotating black hole with total luminosity $L < 10^{-3}L_E$, when protons of the accreted plasma move along regular orbits, gives γ-ray luminosities $L_\gamma \approx 10^{-2}L$.

The possibility that massive black holes participate in the generation of cosmic rays with high and superhigh energies is connected with the electromagnetic fields and shock waves[176–180] initiated by the accretional activity of the black hole, similarly to the acceleration of particles on non-steady-state stars, in explosions of supernova, and in the cosmic medium. In these processes black holes play the role of a general powerhouse—a primary mover, determining only the scale of the phenomena, while the specific features of their strong gravitational field, in particular, the extreme compactness of the source and the presence of an event horizon, are unimportant. The energy of the accelerated particles, which is in principle reached by means of unipolar induction in the magnetosphere of the black hole, according to Eqs. (6.2) and (6.3) is $E_{max} \sim eBr_g \leqslant 5 \cdot 10^{20}(M/10^9 M_\odot)$ eV. The real energies of particles accelerated near massive black holes by the induced electromagnetic field are apparently significantly lower than this maximum estimate, taking into account screening in the plasma and energy lost by the accelerated protons when they interact with the intense light field of the accretion disk.[181,182]

The observed powerful relativistic jets from many nuclei of active galaxies and quasars indicate that particles are efficiently accelerated in the rarefied funnels of a thick accretion disk near the rotation axis of a massive black hole. Even in the absence of a strong induced electromagnetic field, the plasma penetrating into these funnels is accelerated up to subrelativistic velocities by the intense flux of radiation emanating from the central source and is collimated by the walls of the funnels.[183–185] In order to maintain the collimation of the plasma jets at large distances from the black hole a characteristic focusing magnetic field must be generated in them.[4–7,186,188] The energy in the large-scale plasma jets could be transported only by protons and relativistic neutrons,[189,190] since the electronic component losses energy extremely rapidly owing to unavoidable synchrotron losses and this energy must be continuously replenished. An alternative variant is conversion, occurring in the source itself, of accelerated particles into a beam of high-energy γ-rays[191,192] that is capable of propagating without losses over large distances. When such a beam enters a cloud of dense plasma the energy of the beam is once again converted into the energy of charged particles.

The development of instabilities in powerful jets and their interaction with the surrounding galactic and intergalactic matter is accompanied by the formation of large-scale shock waves, which are observed[6,193–197] in the form of hot spots in jets or at the tips of the jets. Formation of standing shock waves is also possible in the accretion flow itself[198–201] in direct proximity to the black hole. The acceleration of the particles on the shockfronts near the black hole and in hot spots in the jets is apparently more effective for generation of cosmic rays than the direct acceleration of particles in the magnetosphere of the black hole by the induced electromagnetic field. In this case the most energetic cosmic rays are generated in the large-scale shock waves far from the central source, where the density of the radiation from the source is low and the energy losses of the accelerated protons owing to photodisintegration are correspondingly small.

7. OBSERVATIONAL STATUS OF MASSIVE BLACK HOLES

7.1. The key facts regarding the existence of massive black holes in the nuclei of active galaxies

The nucleus of a galaxy is the natural environment for the existence of elusive black holes. In spite of the fact that a black hole is an entity in itself, its interaction with the surrounding world is strictly determined. For this reason, the reality of black holes can be confirmed or rejected experimentally. The gigantic difference between the spatial and temporal scales of the observer and a black hole makes it impossible to perform a direct (contact) experiment with the object of study that would prove the reality of the existence of the object. In this respect, black holes, as generally all distant astrophysical objects and the universe as a whole, are similar to the objects of the microuniverse. The fact of their objective reality as well as their properties are determined only by a systematic analysis (theory) of a collection of indirect indications (experiment). The observational problems of detection of black holes are, in turn, similar to the difficulties of any experiment and are attributable to the low signal-to-noise ratio owing to the low intensity of the signal or attenuation of the signal on passage through matter between the black hole and the observer.

There are many detailed reviews of the observational status of the phenomenon of activity in galactic nuclei (see, for example, Refs. 4–7, 127, and 203 and the literature cited there). The enormous mass of accumulated observational data agrees completely with the scheme of accretion on massive black holes and does not contain even one fact that contradicts this picture. The heuristic advantage of the scheme with a massive black hole is its universality and evolutionary unavoidability,[1–4] in contrast to the alternative models (supermassive star, compact cluster of exploding stars, etc.), which explain only part of the observed picture and can be realized only at a relatively short intermediate stage along the path to a massive black hole. We shall list the basic experimental facts pertaining to the problems of massive black holes.

The nuclei of an entire class of galaxies manifest high activity and contain a powerful spatially unresolvable central source of radiation. Galaxies with active nuclei have, as a rule, a spiral structure and they constitute of the order of 10% of the total number of galaxies.[202] The gigantic elliptical cD galaxies, likewise lying at the centers of rich clusters of galaxies, also exhibit intense activity. They apparently

form when several galaxies merge. The brightest objects of this class are quasars, round which it is possible to observe, in spite of the enormous distances to the galaxies, diffuse nebulae having spectra which are characteristic for galaxies. Radiation from the nuclei of active galaxies is recorded at all wavelengths, ranging from radio waves to hard γ-rays, and corresponds qualitatively to an equal distribution of radiation energy over the frequency decades. Under the assumption that the radiation is isotropic, the luminosity of the brightest objects reaches 10^{47}–10^{48} ergs·s^{-1}, which, on the basis of the Eddington limit of stationary luminosity, corresponds to masses of the central sources $M_h \sim 10^9$–$10^{10} M_\odot$ or even larger. The central sources are often variable. Their brightness can change by tens of percent in one day, in a few hours, or even more quickly; this indicates that the size of the region in which the radiation is formed is of the order of the gravitational radius of the source $r_g = GM_h/c^2 \approx 1.5 \cdot 10^{14} (M_h/10^9 M_\odot)$ cm. In the radio range and in some cases also in the optical range (the galaxy M87 and the quasar 3C 273), gigantic collimated jets emanating from the nuclei of active galaxies are observed.[4-7,203] These ejections, whose linear dimensions are significantly greater than those of the galaxy, have diverse shapes and a complicated structure. Near the source they usually extend in two opposite directions, but unilateral ejections are also observed. From the polarization variations it follows that a regular magnetic field exists in the jets. In many cases the jets terminate in enormous regions of radio emission with linear dimensions of hundreds of kpc and even larger [about 5 Mpc for 3C 236 (Ref. 204)]. Such extended binary structures, whose brightness usually increases toward the outer edge, have synchrotron emission spectra and the total energy of the relativistic electrons in them reaches 10^{60}–10^{61} ergs. The extended binary radio structures are genetically related (through the jets) with the central source, so that they should exhibit the properties of a massive gyroscope and maintain continuous pumping of energy into the jet for long time intervals, up to 10^8 yr. Material mass exceeding 10^7–10^8 $M_\odot$ should be "converted" only into the generation of jets. Only a very massive relativistic object could cope with such a problem.

Analysis of the spatial distribution of the nuclei of active galaxies and radio sources associated with them shows that there exists a strong evolution of their brightness on time scales of the order 10^8–10^9 yr. The strong evolution agrees with the scheme of growth of the mass of a central source, whose rate, if the reservoir of the accreted matter is sufficiently large, is related with the Eddington luminosity limit and corresponds to exponential growth with characteristic time $\tau = \eta t_E$, where $t_E = \sigma_T c/4\pi G m_p \approx 5 \cdot 10^8$ yr, and accretion efficiency $\eta \sim 0.1$. When the accreted material is exhausted, the central source is extinguished and the parent galaxy returns into the state of a normal galaxy. Radical restructuring of galactic nuclei occurs in tidal interactions of close galaxies, resulting in merging of the galaxies. Traces of merging occurring in the recent past are observed in the morphological structure of some active galaxies.[46,47] Dissipative tidal interactions and merging of galaxies are most likely to occur in the central, densest parts of clusters of galaxies. Here separate galaxies can survive several mergings, resulting in recurrent activity of their nuclei. Examples of such galaxies are the gigantic cD galaxies which predominate in the central parts of rich clusters. The mass of these same galaxies and correspondingly the mass of their nuclei increase as the hot intergalactic gas contained in the clusters accretes on them. The contraction of the accreted gas is accompanied by rapid emission of its thermal energy and formation of the so-called cooling flows,[44,45] in which intense formation of stars occurs.

The spectral characteristics and temporal variations of the radiation fluxes give information about accretion in the immediate vicinity of the central source. In the x-ray range the spectra of the nuclei of active galaxies usually have a power-law form[205-207] with an exponent $\alpha \sim 0.6$–0.8. Such spectra are characteristic for accretion disks, which in the most extreme cases are surrounded by a hot outflowing corona. In the energy range $\leqslant 1$ keV there is observed a small excess of radiation (usually varying on the scale of several hours and less), which can be interpreted as thermal radiation from the interior of the disk.[208,209] Sharp changes in the luminosity of the central source are usually accompanied, after several days, by a change in the intensity of the emission lines. The observation of delay in the growth of the red wing of the lines relative to the blue wing can be regarded as direct evidence of accretion accompanying such a jump in the intensity of radiation from the nucleus of an active galaxy.[24-29] The Doppler effect and the red shift cause the spectral lines emitted by a rotating disk to have an asymmetric, double-hump shape. The brightest Balmer line H_α of several quite close nuclei of active galaxies is observed to have such a profile.[29] Tracking of the change in the line profiles also gives information about the inner and outer sizes of the disk and the mass of the central source.

The natural striving to correlate diverse observational data has led to an attempt to construct for the activity of galactic nuclei a universal scheme[210-213] which reduces the differences of their qualitative characteristics to a single dependence on the orientation of the symmetry axis of the central source relative to the line of sight. This beautiful idea has been tested for only several examples, and its viability, at least in the initial form, has not been reliably established. Nonetheless, justification is being developed for the hypothesis that the central source is usually surrounded by a massive toroidal absorbing cloud of dust and gas with outer size ranging from ~ 1 pc to ~ 1 kpc, while the main energy release occurs from a much smaller region whose size is of the order of the gravitational radius of the black hole. The axis of symmetry of the toroidal cloud is close to or coincides with the rotation axis of the hole and ultimately is related with the rotation of the parent galaxy. The anisotropy of absorption of the radiation emitted from the central source in the toroidal cloud is equivalent to collimation of the radiation along the axis of the torus and strong dependence of the spectrum and apparent shape of the source on the mutual orientation of the axis of the torus and the line of sight. For this reason the study of the nucleus of only one active galaxy gives only partial and far from complete information about the essence of the "monster" at its center, analogously to determining the shape of an elephant by feel in the ancient parable about the blind wanderers. If the angle between the axis of symmetry of the torus and the line of observation is sufficiently small, then the observer can see the characteristic radiation emitted from the accretion disk and the relativistic jet. In this case, depending on the ratio of the intensities of the radiation of the disk and the jet, the source is a classical quasar or

an object of the type BL Lac with rapid variability and weak emission lines or no emission lines at all. Finally, in the case of more moderate luminosity a nucleus of a Seyfert galaxy of type 1 is observed with wide emission lines, forming in the high-velocity clouds of gas directly adjacent to the central source. If, however, the angle between the symmetry axis of the torus and the line of sight is quite large, then the spectrum of the central source is strongly distorted due to absorption. In this case, sufficiently close sources look like the nucleus of a type-2 Seyfert, whose emission lines are narrower than those of Seyfert galaxies of type 1. More distant sources with powerful jets are, however, now represented in the form of radio galaxies with weak central component or no central component at all and extended binary regions of radio emission, into which energy is pumped by jets from the central source. A universal source also solves the problem of the observed x-ray background, since it predicts for most nuclei of active galaxies anisotropic x-ray emission with average intensity at a level typical for Seyfert galaxies of type 1. In this case the x-ray background can be attributed to the activity of their central sources.[213]

7.2. The problem of a massive black hole in the nucleus of our galaxy

The search for "dead" quasars in the nuclei of normal galaxies is based on the study of the distribution and motion of matter in their central parts as well as the detection of possible sporadic outbursts of activity. The most intriguing possibility here is the center of our own galaxy, the distance to which is "only" 8 kpc. This special region has a complicated spatial and dynamical structure,[20,21] which is extremely difficult to interpret because of the dense molecular clouds and dust obscuring it. Most of the matter within the central parsec is concentrated in stars and constitutes only $10^7 M_\odot$. Because of the enormous attenuation of stellar radiation by dust (up to 30 stellar magnitudes) the parameters of the central region of this cluster and the star composition are known discouragingly poorly. The radius of the dynamically separated core of the star cluster is apparently of the order of 0.1 pc, but it could be even smaller. The velocity dispersion of the stars in this region is equal to about 100 $km \cdot s^{-1}$; the velocities of the gaseous condensations are significantly higher. The massive black hole proposed at the center of the galaxy is associated with the unique radio source Sgr A*, which has an unusual (for galactic objects) growing spectrum with exponent $a \sim 0.25$ and a linear size of about 10 AU.[214] The radio source lies within a system of infrared radiation of IRS 16, but it does not coincide with any of its bright parts.[215,216] The bolometric luminosity of Sgr A*, determined from the thermal emission of the dust surrounding it,[217] does not exceed $\sim 2 \cdot 10^7 L_\odot$. In the x-ray range this is an ordinary source with a luminosity of not greater than 10^{36}–10^{37} $ergs \cdot s^{-1}$. Variable emission of a narrow 0.511 MeV annihilation line is recorded sporadically from the center of the galaxy,[218–220] but it is difficult to identify the source of this emission because of the low angular resolution of γ-ray telescopes. It has not been excluded that a significant fraction of the radiation from the source Sgr A* escapes observation and is channeled into narrow jets of high-energy γ-rays, pulsating along the rotation axis of the galaxy.[191] Such reduced-scale analogs of jets from the nuclei of active galaxies can be observed only when a sufficiently dense cloud of gas crosses them.

A compact structure of ionized gas with the dimensions 1×3 pc^2 and the shape of a triple "spiral," part of which most likely forms a ring, is also associated with the radio source Sgr A*.[221,222] Some fifteen dense clouds with star type masses, dimensions of $\sim$0.1–0.2 pc, and lifetimes of $\sim 10^4$ yr, move along the spirals. Their individual radial velocities have been measured from the Doppler shift of the infrared line of Ne II recorded from the clouds. If it is assumed that the compact clouds of gas move in a steady-state fashion, then a "point-like" mass $M_h \approx (2\text{–}4) \cdot 10^6 M_\odot$, which could be a black hole, should also be located at the center of the galaxy in addition to the distributed mass of the stars.[20,21] In addition, the existence of invisible mass $M_h \approx 2 \cdot 10^6 M_\odot$ follows from data on the velocities of individual stars,[223,224] though a structure of the central cluster without a massive black hole at all is also admissible.[215,225–230] The stars, whose motion is controlled only by gravitational forces, move with lower velocities than the gas. This indicates that the gaseous component is not in a steady state and, in part, weakens arguments in favor of a massive black hole which are based on measurements of the cloud velocities.

For a star system consisting of solar type stars and close to the predicted mass of the black hole the radius R of the core of the central star system, the radius of influence of the hole r_h, and the critical radius r_{cr} of the region with an unfilled loss cone differ by not more than an order of magnitude. Under these conditions an appreciable peak is not formed in the density of stars $n(r) \propto r^{-7/4}$, and the loss cone is practically completely filled with stars. The rate of flow of stars onto the black hole is equal to $\dot{M}_h \sim 10^{-4}$–$10^{-3} M_\odot$ yr^{-1} and should lead to an average luminosity of the galactic center exceeding the observed luminosity.[231] This difficulty can be overcome, if the black hole "digests" sufficiently rapidly the matter of each star that is destroyed by its tidal forces. Then the activity of the black hole is manifested in the form of short, of the order of several years, bright outbursts,[103,104] which repeat with intervals of $\sim 10^3$–10^4 yr. Another variant could be dispersal of a significant mass of gas, freed in the process of tidal breakup of stars, from the immediate vicinity of the hole.[104,105] Finally, absorption of whole stars by the hole without the formation of an accretion disk is possible when the stars fall inside the loss cone and their trajectories do not have a turning point.

Even when the problem of the luminosity of the galactic center is solved, however, there remains the more difficult problem of the unavoidable strong growth of a sufficiently massive central black hole over the lifetime of the galaxy.[122] Over a time of the order of 10^{10} yr the predicted massive black hole should "eat away" the entire star cluster surrounding it and by now it would be completely dominant at the center of the galaxy. Hope of slowing down the growth of the black hole by dispersal of the remnants of disintegrated stars can potentially be justified only under conditions of an empty loss cone, when the stars approach the sphere of tidal breakup along tangential trajectories and therefore have a turning point. The conditions of a galactic center with $M_h \approx (2\text{–}4) \cdot 10^6 M_\odot$ and a solar type star composition, however, correspond to a filled loss cone. For this reason, a significant fraction of the stars penetrating the sphere of tidal

breakup have trajectories which dive into the hole and are absorbed by the hole practically completely without appreciable release of energy (i.e., without a burst of activity). The growth of a massive black hole in a star cluster can stop, as shown in Sec. 5.4, only if there forms around the hole a peak in the star density $n(r) \propto r^{-7/4}$, that is capable of supporting heating and secular expansion of the central part of the cluster. The necessary condition of formation of such a density peak $r_{cr} \ll r_h$ is easily satisfied, if the central cluster in the nucleus of the galaxy consists primarily of stars which are more compact than the sun, for example, white dwarfs. This possibility is quite realistic for a very old cluster of stars in the nucleus of the galaxy. The main growth of the black hole in it should occur primarily as a result of absorption of the most massive and extended stars, concentrated closer to the center of the cluster and having a large cross section for tidal breakup. The growth of a massive black hole at the center of the galaxy with a large abundance of white dwarfs there will stop[121] already for $M_h \geqslant 10^5 M_\odot$. The activity problem is also solved at the same time, since white dwarfs are absorbed whole by a black hole of such a mass without first being broken up by tidal forces.

7.3. The nuclei of nearby galaxies

The significant distance of the nuclei of neighboring galaxies, as compared with the center of our galaxy observed through a thick layer of absorbing matter, is partially compensated by the possibility of performing detailed spectroscopic observations in the optical range. The nucleus of the closest spiral galaxy M31 (the Andromeda nebula) and its satellite, the dwarf galaxy M32, is observed to have a strong radial dependence of the mass—luminosity ratio as well as a sharp growth of the rotational velocity and velocity dispersion of the stars, within several central angular seconds, toward the center.[13-18] An analogous kinematic picture was found in the nucleus of the more distant spiral galaxy M104 ("Sombrero").[19] Modeling of the distribution and motion of stars in the nuclei of these galaxies as well as in our galaxy shows that they must contain "point-like" masses of size $(6\text{–}7)\cdot 10^7 M_\odot$ (in M31), $(5\text{–}8)\cdot 10^6 M_\odot$ (in M32), and $\sim 10^8 M_\odot$ (in M104). We also note that a compact formation with a mass of the order of $10^3 M_\odot$ is observed at the center of the rich globular star cluster M15.[232] All these point-like (at the current level of observations) concentrations of matter could easily be massive black holes with weak activity owing to weak accretion flow on them. Under the conditions of the nuclei of normal galaxies the accreted gas is supplied mainly by stars that are broken up by the tidal forces of massive black holes once in 10^3–10^5 yr. For this reason, in such normal galaxies bursts of activity, which convert their nuclei into Seyfert nuclei for a short time, should unavoidably occur. Tt has not been excluded, however, that the "point-like" mass observed in the nuclei of normal galaxies, including our own galaxy, consists merely of a very compact star cluster or a massive bridge—a bar.[225-230] The dilemma—a massive black hole or a compact star cluster—also cannot be solved, as yet, in the case of the closest active elliptical galaxy M87, lying at the center of the cluster of galaxies in Virgo. The nucleus of this gigantic galaxy, which is also famous for its optical jet, could contain a black hole with mass $M_h \approx (3\text{–}4)\cdot 10^9 M_\odot$.[22,23,233-235]

Many uncertainties, due to the inadequate angular resolution of present instruments, regarding the presence of massive black holes in the nuclei of nearby galaxies could be resolved in the next few years with the help of the Hubble space telescope, which is already in orbit, or with the help of ground-based optical interferometers that are being designed.

8. CONCLUSIONS

The evolutionary chain "star cluster—supermassive star—massive black hole at the center of a compact star cluster—isolated supermassive black hole" is the most direct and natural path of the evolution of galactic nuclei. Such an evolutionary scheme is realized with the minimum requirements on initial conditions, which reduce, generally speaking, only to the fact itself that a massive compact star cluster is formed at the center of the galaxy and the further fate of the cluster is predetermined. The dynamical evolution of star clusters with parameters typical for galactic nuclei is accompanied by an increase in the velocity dispersion of the stars in their central parts and correspondingly by an increase in the depth of the potential well. This process proceeds quickly, as compared with the evolution of the galaxies themselves and the expansion of the universe. Dissipative processes accelerate the contraction of the central parts of the galactic nuclei and increase the mass of the core, which transforms, as a result of disruptive collisions of stars, into a self-gravitating cloud of gas—a short-lived supermassive star. A collapsing gas cloud with mass $M > 10^8 M_\odot$ has, as the gravitational radius is approached, an average density of less than that of water and a temperature at the center that is too low for ignition of nuclear reactions. For this reason the collapse proceeds under conditions such that the standard laboratory physics is applicable locally. In principle only extremely rapid rotation, resulting in fragmentation, is capable of preventing at this stage the collapse of a large part of a supermassive star into a black hole. The diversity of astrophysical manifestations of massive black holes is connected with the uniqueness of the distribution and motion of matter in separate galaxies, which are also subjected to significant structural changes which occur in interactions with neighboring galaxies. Physically distinguishable stages of the dynamical evolution of central star clusters, which end with the formation of a massive black hole, encompass the entire wide range of observed galactic nuclei, extending from dwarf galaxies with weak central contraction through different types of galaxies with active nuclei up to gigantic elliptical galaxies in clusters with a powerful central source and their extreme manifestation—quasars.

In the large ocean of accumulated observational facts and theoretical innvestigations three "whales" of the astrophysics of massive black holes stand out:

1. *The evolutionary unavoidability of condensation in the centers of galaxies of large masses of gas and stars within a region of the order of the overall gravitational radius.*

2. *The absence of interference and difficulties of a fundamental character, preventing the collapse of large masses of matter into a black hole.*

3. *The consistency of the observed picture of activity of the nuclei of galaxies and quasars (source compactness, lifetime, radiation intensity, overall energetics, jets) with the expected properties of accreting black holes.*

These three whales also serve as a support for black holes of star type masses, which can be manifested most actively in binary star systems. Nonetheless confidence in the physical reality of black holes is currently based on the results of a comprehensive analysis of only indirect indications. For this reason, it is natural to have a feeling of dissatisfaction. This stimulates further research in this field, which is fundamental for all of physics. The last step in the problem of black holes can be only direct evidence for the existence on a massive object of a steady-state light surface—the event horizon.

In conclusion I wish to express my deep appreciation to L. M. Ozernoĭ for fruitful collaboration and my gratitude to V. S. Berezinskiĭ, V. L. Ginzburg, I. D. Novikov, A. G. Polnarev, and V. P. Frolov for discussions and valuable remarks.

[1] Ya. B. Zel'dovich and I. D. Novikov, *Relativistic Astrophysics*, Univ. Chicago Press, 1971 [Russ. original, Nauka, M., 1967].
[2] Ya. B. Zel'dovich and I. D. Novikov, *The Structure and Evolution of the Universe*, Univ. Chicago Press, 1983 [Russ. original, Nauka, M.], *Theory of Gravitation and the Evolution of Stars* [in Russian], Nauka, M., 1971.
[3] C. W. Misner, K. S. Thorne, and J. A. Wheeler, *Gravitation*, Freeman and Co., San Francisco, 1973. [Russ. transl., Mir, M., 1977].
[4] V. L. Ginzburg and L. M. Ozernoy, Astrophys. Space Sci. **50**, 23 (1977).
[5] M. J. Rees, Phys. Scr. **17**, 193 (1978).
[6] M. C. Begelman, R. D. Blandford, and M. C. Rees, Rev. Mod. Phys. **56**, 255 (1984).
[7] P. J. Wiita, Phys. Rep. **123**, 117 (1985).
[8] F. Hoyle and W. A. Fowler, *Nature* **197**, 533 (1963).
[9] L. M. Ozernoĭ, Astron. Zh. **43**, 300 (1966) [Sov. Astron. **10**, 241 (1966)].
[10] G. S. Bisnovatyĭ-Kogan, Ya. B. Zel'dovich, and I. D. Novikov, Astron. Zh. **44**, 525 (1967) [Sov. Astron. **11**, 419 (1967)].
[11] M. Morrisin, Astrophys. J. Lett. **157**, L73 (1969).
[12] S. A. Colgate, Astrophys. J. **150**, 163 (1967).
[13] J. L. Tonry, Astrophys. J. Lett. **283**, L27 (1984).
[14] A. Dressler, Astrophys. J. **286**, 97 (1984).
[15] J. L. Tonry, Astrophys. J. **322**, 632 (1987).
[16] J. Kormendi, Astrophys. J. **325**, 128 (1988).
[17] A. Dressler and D. O. Richstone, Astrophys. J. **324**, 701 (1988).
[18] J. Kormendi, Astrophys. J. **335**, 40 (1988).
[19] B. J. Jarvis and P. Dubeth, Astron. Astrophys. **201**, L33 (1988).
[20] R. L. Brown and H. S. Liszt, Ann. Rev. Astron. Astrophys. **22**, 223 (1984).
[21] R. Gensel and C. H. Townes, Ann. Rev. Astron. Astrophys. **25**, 377 (1987).
[22] P. J. Young, J. A. Westphal, J. Kristian, C. P. Wilson, and C. P. Landauer, Astrophys. J. **221**, 721 (1978).
[23] W. L. W. Sargent, P. J. Young, A. Boksenberg, K. Shortindge, C. R. Lynds, and F. D. A. Hartwick, Astrophys. J. **221**, 731 (1978).
[24] S. M. Fabrika, Astron. Tsirkulyar, No. 1109, 1 (1980).
[25] R. D. Blandford and C. F. McKee, Astrophys. J. **255**, 419 (1982).
[26] E. Perez, M. V. Penston, C. Tadhunter, E. Mediavilla, and M. Males, Mon. Not. R. Astron. Soc. **230**, 353 (1988).
[27] K. Chen, P. Helpern, and A. V. Fillipenko, Astrophys. J. **339**, 742 (1988).
[28] C. M. Gaskell, Astrophys. J. **325**, 114 (1988).
[29] L. Stella, *Nature* **344**, 747 (1990).
[30] M. Schmidt, *Nature* **197**, 1040 (1963).
[31] Ya. B. Zel'dovich and I. D. Novikov, Dokl. Akad. Nauk SSSR **155**, 1033 (1964) [Sov. Phys. Dokl. **9**, 246 (1964)].
[32] Ya. B Zel'dovich and M. A. Podurets, Dokl. Akad. Nauk SSSR **156**, 57 (1964) [Sov. Phys. Dokl. **9**, 373 (1964)].
[33] E. E. Salpeter, Astrophys. J. **140**, 796 (1964).
[34] D. Lynden-Bell, *Nature* **223**, 690 (1969).
[35] Ya. B. Zel'dovich and I. D. Novikov, Astron. Zh. **43**, 758 (1966) [Sov. Astron. **10**, 602 (1966)].
[36] S. W. Hawking, Mon. Not. R. Astron. Soc. **152**, 75 (1971).
[37] A. Polnarev and M. Yu. Khlopov, Usp. Fiz. Nauk **145**, 369 (1985) [Sov. Phys. Usp. **28**, 213 (1985)].
[38] S. W. Hawking, Phys. Lett. B **231**, 237 (1989).
[39] I. D. Novikov and V. P. Frolov, *The Physics of Black Holes* [in Russian], Nauka, Moscow (1986).
[40] S. L. Shapiro and S. A. Teukolsky, *Black Holes, White Dwarfs and Neutron Stars*, Wiley, N. Y., 1983. [Russ. transl., Mir, M., 1985].
[41] S. Chandrasekhar, *Mathematical Theory of Black Holes*, Oxford Univ. Press, N.Y., 1983 [Russ. transl., Mir, M., 1986], Vols. 1 and 2.
[42] D. V. Gal'tsov, *Particles and Fields in the Vicinity of a Black Hole* [in Russian], Moscow Univ. Press, M., 1986.
[43] K. S. Thorne, R. H. Price, and D. A. Macdonald [Eds.], *Black Holes: The Membrane Paradigm*, Yale University Press, Hew Haven, 1986 [Russ. transl. Mir, M., 1988].
[44] A. C. Fabian, P. E. J. Nulsen, and C. R. Canizares, *Nature* **310**, 733 (1984).
[45] C. L. Sarazin, Rev. Mod. Phys. **58**, 1 (1986).
[46] A. Toomre and J. Toomre, Astrophys. J. **178**, 623 (1972).
[47] G. B. Byrd and S. V. Valtaoja, Astron. Astrophys. **166**, 75 (1986).
[48] V. A. Ambartsumyan, Uch. zap. LGU **22**, 19 (1938).
[49] L. Spitzer, Mon. Not. R. Astron. Soc. **100**, 396 (1940).
[50] V. A. Antonov, Vestn. LGU. Ser. Matematika, Mekhanika, Astronomiya **7**, 135 (1962).
[51] D. Lynden-Bell and R. Wood, Mon. Not. R. Astron. Soc. **138**, 495 (1968).
[52] A. V. Gurevich and K. P. Zybin, Zh. Eksp. Teor. Fiz. **97**, 20 (1990) [Sov. Phys. JETP **70**, 10 (1990)].
[53] Ya. B. Zel'dovich and M. A. Podurets, Astron. Zh. **42**, 963 (1965) [Sov. Astron. **9**, 742 (1966)].
[54] K. S. Bisnovatyĭ-Kogan and Ya. B. Zel'dovich, Astrofizika **5**, 223 (1969). [Astrophysics **5**, 105 (1969)].
[55] S. L. Shapiro and S. A. Teukolsky, Astrophys. J. **298**, 34 (1985).
[56] S. L. Shapiro and S. A. Teukolsky, Astrophys. J. **307**, 575 (1986).
[57] G. D. Quinlan and S. L. Shapiro, Astrophys. J. **343**, 725 (1989).
[58] Lyman Spitzer, Jr., *Dynamic Evolution of Globular Clusters*, Princeton Univ. Press, 1987 [Russ. transl., Mir, M., 1990].
[59] W. C. Saslaw, *Gravitational Physics of Stellar and Galactic Systems*, Cambridge Univ. Press, N.Y., 1985 [Russ. transl., Mir, M., 1989].
[60] L. D. Landau and E. M. Lifshitz, *The Classical Theory of Fields*, 4th ed., Pergamon Press, Oxford, 1975 [Russ. original, Nauka, M., 1973].
[61] S. Chandrasekhar, *Principles of Stellar Dynamics*, Univ. of Chicago Press, Chicago, 1942. [Russ. transl., IL, M., 1948].
[62] L. Spitzer and R. Härm, Astrophys. J. **127**, 544 (1958).
[63] L. E. Gurevich and B. Yu. Levin, Dokl. Akad. Nauk SSSR **70**, 781 (1950) (not translated).
[64] M. Milgrom and S. L. Shapiro, Astrophys. J. **223**, 991 (1978).
[65] M. Henon, Astron. Astrophys. **2**, 251 (1969).
[66] S. J. Aarseth, in *Gravitational N Body Problem*, edited by M. Lecar, D. Reidel, Dordrecht, 1972, p. 88.
[67] D. C. Heggie, Mont. Not. R. Astron. Soc. **173**, 729 (1975).
[68] V. I. Dokuchaev and L. M. Ozernoĭ, Astron. Zh. **55**, 27 (1978) [Sov. Astron. **22**, 15 (1978)].
[69] L. M. Ozernoy and V. I. Dokuchaev, Astron. Astrophys. **111**, 1 (1982).
[70] V. I. Dokuchaev and L. M. Ozernoy, Astron. Astrophys. **111**, 16 (1982).
[71] A. C. Fabian, J. E. Pringle, and M. J. Rees, Mon. Not. R. Astron. Soc. **172**, 15P (1975).
[72] A. P. Lightman and S. L. Shapiro, Rev. Mod. Phys. **50**, 437 (1978).
[73] L. Spitzer and R. D. Mathieu, Astrophys. J. **241**, 618 (1980).
[74] J. Stodolkiewicz, Acta Astron. **32**, 63 (1982).
[75] S. Inagaki, Mon. Not. R. Astron. Soc. **206**, 149 (1984).
[76] E. Bettwieser and D. Sugimoto, Mon. Not. R. Astron. Soc. **208**, 493 (1984).
[77] H. Cohn, P. Hut, and M. Wise, Astrophys. J. **342**, 814 (1989).
[78] W. H. Press and S. A. Teukolsky, Astrophys. J. **213**, 183 (1977).
[79] A. Nduka, Astrophys. J. **170**, 131 (1971).
[80] J. G. Hills, Astron. J. **80**, 1075 (1975).
[81] V. I. Dokuchaev and L. M. Ozernoĭ, Pis'ma Astron. Zh. **7**, 280 (1981) [Sov. Astron. Lett. **7**, 155 (1981)].
[82] L. Spitzer and W. Saslaw, Astrophys. J. **143**, 400 (1966).
[83] L. Spitzer, *Galactic Nuclei*, (Ed.) D. O. O'Connel, North-Holland, Amsterdam, 1971, p. 443.
[84] G. S. Bisnovatyĭ-Kogan, Pis'ma Astron. Zh. **4**, 130 (1978) [Sov. Astron. Lett. **4**, 69 (1978)].
[85] R. H. Sanders, Astrophys. J. **162**, 791 (1970).
[86] I. Shlosman and M. C. Begelman, Astrophys. J. **341**, 685 (1989).
[87] T. Langbein, R. Spurzem, K. J. Fricke, and H. W. Yorke, Astron. Astrophys. **227**, 333 (1990).
[88] J. P. Ostriker and P. J. E. Peebles, Astrophys. J. **186**, 467 (1973).
[89] M. C. Begelman and M. J. Rees, Mon. Not. R. Astron. Soc. **188**, 847 (1978).
[90] A. F. Illarionov and M. M. Romanova, Astron. Zh. **65**, 535 (1988) [Sov. Astron. **32**, 274 (1988)].
[91] S. Chandrasekhar, *Stellar Structure*, University of Chicago Press, Chicago, 1939. [Russ. transl., IL, M., 1950].

[92] M. Schwarzschild, *Structure and Evolution of the Stars*, Princeton Univ. Press, 1958 [Russ. transl., Inostr. lit., M., 1961].
[93] W. A. Fowler, Astrophys. J. **144**, 180 (1966).
[94] A. G. Cavaliere, P. Morrison, and P. Pacini, Astrophys. J. Lett. **162**, L133 (1970).
[95] L. M. Ozernoy and V. V. Usov, Astrophys. Space Sci. **13**, 209 (1973).
[96] L. M. Ozernoĭ and V. E. Chertoprud, Astron. Zh. **43**, 20 (1966) [Sov. Astron. **10**, 15 (1966)].
[97] M. A. Abramowicz, M. Jaroczynski, and M. Sikora, Astron. Astrophys. **63**, 221 (1978).
[98] M. Jaroczynski, M. A. Abramowicz, and B. Paczynski, Acta Astron. **30**, 1 (1980).
[99] M. A. Abramowicz, M. Calvani, and L. Nobili, Astrophys. J. **242**, 772 (1980).
[100] M. J. Rees, M. C. Begelman, R. D. Blandford, and E. S. Phinney, *Nature* **295**, 17 (1982).
[101] J. G. Hills, *Nature* **254**, 295 (1975).
[102] C. R. Evans and C. S. Kochanek, Astrophys. J. Lett. **346**, L13 (1989).
[103] V. G. Gurzadyan and L. M. Ozernoy, Astron. Astrophys. **95**, 39 (1981).
[104] M. J. Rees, *The Galactic Center*, edited by G. R. Riegler and R. D. Blandford, American Institute of Physics, New York (1982), p. 166 (AIP Conf. Proc. No. 83).
[105] M. J. Rees, *Nature* **333**, 523 (1988).
[106] L. D. Landau, Zh. Eksp. Teor. Fiz. 7, 203 (1937), transl. in *Collected Papers of Landau*, edited by D. terHaar, Gordon and Breach, N. Y., 1967, p. 163.
[107] J. Frank and M. J. Rees, Mon. Not. R. Astron. Soc. **176**, 633 (1976).
[108] V. I. Dokuchaev and L. M. Ozernoĭ, Pis'ma Astron. Zh. **3**, 391 (1977) [Sov. Astron. Lett. **3**, 209 (1977)].
[109] V. I. Dokuchaev and L. M. Ozernoĭ, Pis'ma Astron. Zh. **3**, 295 (1977) [Sov. Astron. Lett. **3**, 157 (1977)].
[110] M. J. Duncan and S. L. Shapiro, Astrophys. J. **268**, 565 (1983).
[111] A. V. Gurevich, Geomagn. Aeron. **4**, 247 (1964) [Geomagn. Aeron. (USSR) **4**, 192 (1964)].
[112] J. N. Bahcall and R. A. Wolf, Astrophys. J. **209**, 214 (1976).
[113] A. P. Lightman and S. L. Shapiro, Astrophys. J. **211**, 244 (1977).
[114] V. I. Dokuchaev and L. M. Ozernoĭ, Zh. Eksp. Teor. Fiz. **73**, 1587 (1977) [Sov. Phys. JETP **46**, 834 (1977)].
[115] A. C. Fabian, Proc. R. Soc. London A **366**, 449 (1979).
[116] A. C. Edwards, Mon. Not. R. Astron. Soc. **190**, 757 (1980).
[117] P. Gondhalekar, P. O'Brian, and R. Wilson, Mon. Not. R. Astron. Soc. **222**, 71 (1986).
[118] A. J. Allen and P. A. Hughes, Astrophys. J. **313**, 152 (1987).
[119] G. M. Voit and J. M. Shull, Astrophys. J. **331**, 197 (1988).
[120] R. C. Pueter, *Active Galactic Nuclei*, (Eds.) D. E. Osterbock and J. S. Miller, Kluwer Acad. Publ., Dordrecht (1989), p. 137.
[121] V. I. Dokuchaev, Pis'ma Astron. Zh. **15**, 387 (1989) [Sov. Astron. Lett. **15**, 167 (1989)].
[122] V. I. Dokuchaev and L. M. Ozernoĭ, Pis'ma Astron. Zh. **3**, 212 (1977) [Sov. Astron. Lett. **3**, 112 (1977)].
[123] S. L. Shapiro, Astrophys. J. **217**, 281 (1977).
[124] V. I. Dokuchaev, Pis'ma Astron. Zh. **16**, 970 (1990) [Sov. Astron. Lett. **16**, 416 (1990)].
[125] W. G. Mathews and E. R. Capriotti, *Astrophysics of Active Galaxies*, edited by J. S. Miller, Mill Valley Univ. Science Books, Oxford (1985), p. 185.
[126] D. J. Raine, Vistas Astron. **32**, 321 (1988).
[127] J. N. Bregman and J. R. Boisseau, Astrophys. J. **347**, 118 (1989).
[128] M. Contini and S. M. Viegas-Aldrovandi, Astrophys. J. **350**, 12 (1990).
[129] V. S. Berezinskiĭ and L. M. Ozernoĭ, Astron. Zh. **58**, 505 (1981) [Sov. Astron. **25**, 286 (1981)].
[130] A. Beloborodov, R. Ivanov, A. Illarionov, and A. Polnarev, to be published.
[131] V. G. Gorbatskiĭ, Soobshch. LGU **22**, 16 (1965).
[132] V. F. Shvartsman, Astron. Zh. **48**, 479 (1971) [Sov. Astron. **15**, 377 (1971)].
[133] J. Pringle and M. Rees, Astron. Astrophys. **21**, 1 (1972).
[134] N. I. Shakura and R. A. Sunyaev, Astron. Astrophys. **24**, 337 (1973).
[135] I. D. Novikov and K. S. Thorne, *Black Holes*, (Eds.) C. de Witt and B. S. de Witt, Gordon and Breach, N.Y., 1973, p. 343.
[136] M. J. Rees, Ann. Rev. Astron. Astrophys. **22**, 471 (1984).
[137] D. Lynden-Bell and M. J. Rees, Mon. Not. R. Astron. Soc. **152**, 461 (1971).
[138] R. D. Blandford and R. L. Znajek, Mon. Not. R. Astron. Soc. **179**, 433 (1977).
[139] P. V. E. Lovelace, *Nature* **262**, 649 (1976).
[140] J. N. Bardeen, W. H. Press, and S. A. Teukolsky, Astrophys. J. **178**, 347 (1972).
[141] K. S. Thorne, Astrophys. J. **191**, 507 (1974).
[142] J. E. Pringle, M. J. Rees, and A. G. Pacholczyk, Astron. Astrophys. **29**, 179 (1973).
[143] N. I. Shakura and R. A. Sunyaev, Mon. Not. R. Astron. Soc. **175**, 613 (1976).
[144] J. E. Pringle, Ann. Rev. Astron. Astrophys. **19**, 137 (1981).
[145] J. F. Hawley, L. L. Smarr, and J. R. Wilson, Astrophys. J. **277**, 296 (1984).
[146] R. D. Blandford, M. Jaroszynski, and S. Kumar, Mon. Not. R. Astron. Soc. **215**, 667 (1985).
[147] J. C. B. Papaloisou and J. E. Pringle, Mon. Not. R. Astron. Soc. **225**, 267 (1987).
[148] M. Camenzind, F. Demole, and N. Straumann, Astron. Astrophys. **158**, 212 (1986).
[149] Yu. E. Lyubarskiĭ and N. I. Shakura, Pis'ma Astron. Zh. **13**, 917 (1987) [Sov. Astron. Lett. **13**, 386 (1987)].
[150] G. S. Bisnovatyi-Kogan and S. I. Blinnikov, Astron. Astrophys. **59**, 111 (1977).
[151] F. Takahara, R. Rosner, and M. Kusanose, Astrophys. J. **346**, 122 (1989).
[152] M. Czerny and A. R. King, Mon. Not. R. Astron. Soc. **241**, 839 (1989).
[153] G. S. Bisnovatyi-Kogan and A. A. Kuzmaikin, Astrophys. Space Sci. **42**, 401 (1976).
[154] G. D. Lominadze and Q. Q. Chagdeshvili, in *Problems of Nonlinear and Turbulent Processes in Physics* [in Russian], Kiev, 1985, Part 2, p. 311.
[155] R. E. Pudritz, Mon. Not. R. Astron. Soc. **195**, 897 (1979).
[156] W. Kluzniak, M. Ruderman, J. Shaham, and M. Tavani, *Nature* **336**, 558 (1988).
[157] T. Damour, R. S. Hanni, R. Ruffini, and J. R. Wilson, Phys. Rev. D **18**, 1518 (1978).
[158] D. V. Gal'tsov and V. P. Petukhov, Zh. Eksp. Teor. Fiz. **74**, 801 (1978) [Sov. Phys. JETP **47**, 419 (1978)].
[159] A. N. Aliev and D. V. Gal'tsov, Usp. Fiz. Nauk **157**, 129 (1989) [Sov. Phys. Usp. **32**, 75 (1989)].
[160] J. M. Bardeen and J. A. Peterson, Astrophys. J. Lett. **195**, L65 (1975).
[161] T. Damour, Phys. Rev. D **18**, 3598 (1978).
[162] R. L. Znajek, Mon. Not. R. Astron. Soc. **185**, 833 (1978).
[163] F. C. Michel, Rev. Mod. Phys. **54**, 1 (1982).
[164] V. S. Beskin, A. V. Furevich, and Ya. N. Istomin, Usp. Fiz. Nauk **150**, 257 (1986) [Sov. Phys. Usp. **29**, 946 (1986)].
[165] V. L. Ginzburg and L. M. Ozernoi, Zh. Eksp. Teor. Fiz. **47**, 1030 (1964) [Sov. Phys. JETP **20**, 689 (1964)].
[166] R. M. Wald, Phys. Rev. D **10**, 1680 (1974).
[167] V. I. Dokuchaev, Zh. Eksp. Teor. Fiz. **92**, 1921 (1987) [Sov. Phys. JETP **65**, 1079 (1987)].
[168] R. Penrose Rev. Nuovo Cimento **1**, 252 (1969).
[169] Ya. B. Zel'dovich, Pis'ma Zh. Eksp. Teor. Fiz. **14**, 270 (1971) [JETP Lett. **14**, 180 (1971)].
[170] C. W. Misner, Phys. Rev. Lett. **28**, 999 (1972).
[171] D. S. Wilkins, Phys. Rev. D **5**, 814 (1972).
[172] V. I. Dokuchaev, Pis'ma Astron Zh. **12**, 770 (1986) [Sov. Astron. Lett. **12**, 322 (1986)].
[173] R. I. Kolykhalov and R. A. Syunyaev, Astron. Zh. **56**, 338 (1979) [Sov. Astron. **23**, 189 (1979)].
[174] B. V. Vainer and V. N. Panov, Astrophys. Space Sci. **113**, 1 (1985).
[175] V. S. Berezinskii and V. I. Dokuchaev, Zh. Eksp. Teor. Fiz. **96**, 1537 (1989) [Sov. Phys. JETP **69**, 871 (1989)].
[176] V. S. Berezinsky and V. L. Ginzburg, Mon. Not. R. Astron. Soc. **194**, 3 (1981).
[177] R. J. Protheroe and D. Kazanas, Astrophys. J. **265**, 620 (1983).
[178] V. S. Berezinskiĭ, S. V. Bulanov, V. L. Ginzburg, V. A. Dogel', and V. S. Ptuskin, *Astrophysics of Cosmic Rays* [in Russian], Nauka, M., 1984.
[179] A. M. Hillas, Ann. Rev. Astron. Astrophys. **22**, 425 (1984).
[180] D. Kazanas and D. C. Ellison, Astrophys. J. **304**, 178 (1986).
[181] M. Sikora, J. G. Kirk, M. C. Begelman, and P. Schneider, Astrophys. J. Lett. **320**, L81 (1987).
[182] V. S. Berezinsky, S. I. Grigor'eva, and V. A. Dogel, Astron. Astrophys. **232**, 582 (1990).
[183] M. Sikora and D. B. Wilson, Mon. Not. R. Astron. Soc. **197**, 529 (1981).
[184] J. Fukue, Publ. Astron Soc. Japan **34**, 163 (1982).
[185] A. Ferrari, S. Habbal, R. Rosner, and K. Tsinganos, Astrophys. J. Lett. **277**, L35 (1984).
[186] J. Bekenstein and D. Eichler, Astrophys. J. **298**, 493 (1985).
[187] A. P. Bell, *Nature* **346**, 136 (1990).
[188] J. C. L. Wang, M. E. Sulkanen, and R. V. E. Lovelace, Astrophys. J. **355**, 38 (1990).
[189] J. G. Kirk and A. Mastichiadis, Astron. Astrophys. **213**, 75 (1989).
[190] M. Sikora, M. C. Begelman, and B. Rudak, Astrophys. J. Lett. **341**,

L33 (1989).
[191] N. S. Kardashev, I. D. Novikov, A. G. Polnarev, and B. U. Shtern, Astron. Zh. **60**, 209 (1983) [Sov. Astron. **27**, 119 (1983)].
[192] R. V. E. Lovelace, Astron. Astrophys. **173**, 237 (1987).
[193] C. A. Norman, *Unstable Current Systems and Plasma Instabilities in Astrophysics*, (Eds.) M. R. Kundu and G. D. Holman, D. Reidel, Dordrecht, 1983, p. 85 (IAU Symposium No. 107).
[194] K. Meisenheimer and A. F. Heavens, *Nature* **323**, 419 (1986).
[195] J. J. Quenby and R. Lien, *Nature* **342**, 654 (1989).
[196] M. Sikora and I. Shlosman, Astrophys. J. **336**, 593 (1989).
[197] M. C. Begelman and J. G. Kirk, Astrophys. J. **353**, 66 (1990).
[198] P. Meszaros and J. P. Ostriker, Astrophys. J. Lett. **273**, L59 (1983).
[199] A. Babul, J. Ostriker, and P. Meszaros, Astrophys. J. **347**, 59 (1989).
[200] S. K. Chakrabarti, Astrophys. J. **350**, 275 (1990).
[201] M. A. Abramowicz and S. K. Chakrabarti, Astrophys. J. **350**, 281 (1990).
[202] M. Schmidt, *Highlights of Astrophysics*, (Ed.) D. McNally, IAU (1989), Vol. 8, p. 33.
[203] A. H. Bridle and R. A. Perlay, Ann. Rev. Astron. Astrophys. **22**, 319 (1984).
[204] A. G. Willis, R. G. Strom, and A. S. Wilson, *Nature* **250**, 625 (1974).
[205] A. M. Atoyan and A. Nahapetian, Astron. Astrophys. **219**, 53 (1989).
[206] L. L. Cowie, *Two Topics in X-Ray Astronomy, Proceedings of the 23rd ESLAB Symposium*, ESA (SP-296), (1989), p. 707.
[207] T. J. Turner and K. A. Pounds, Mon. Not. R. Astron. Soc. **240**, 833 (1989).
[208] P. Madau, Astrophys. J. **327**, 116 (1988).
[209] M. C. Begelman, *Active Galactic Nuclei*, (Eds.) D. E. Osterbrock and J. S. Miller, Kluwer Acad. Publ., Dordrecht, 1989, p. 141.
[210] A. Lawrence and M. Elvis, Astrophys. J. **256**, 410 (1982).
[211] P. K. J. Antonucci and J. S. Miller, Astrophys. J. **297**, 621 (1985).
[212] P. D. Barthel, Astrophys. J. **336**, 606 (1989).
[213] G. Setti and L. Woltjer, Astron. Astrophys. **224**, L21 (1989).
[214] K. U. Lo, *The Galactic Center*, (Ed.) D. C. Backer, AIP, N.Y., 1987, p. 30 (AIP Conf. Proc. No. 155).
[215] D. A. Allen and R. M. Sanders, *Nature* **319**, 191 (1986).
[216] E. E. Becklin, H. Dinerstein, J. Gatley, and M. W. Werner, in Ref. 214, p. 162.
[217] E. E. Becklin, I. Gatley, and M. W. Werner, Astrophys. J. **258**, 135 (1982).
[218] N. S. Kardashev, in *Progress in Science and Technology*. Series on Astronomy, VINITI, Akad, Nauk SSSR, Moscow, 1983, Vol. 24, p. 183.
[219] M. Leventhal, C. J. MacCallum, S. D. Barthelmy, N. Genrels, B. J. Teegarden, and J. Tueller, *Nature* **339**, 36 (1989).
[220] R. E. Lingenfelter and R. Ramaty, Astrophys. J. **343**, 686 (1989).
[221] R. D. E. Ekers, J. van Gorcom, U. J. Schwartz, and W. M. Goss, Astron. Astrophys. **122**, 143 (1983).
[222] K. Y. Lo and M. J. Claussen, *Nature* **306**, 647 (1983).
[223] G. H. Rieke and M. J. Rieke, Astrophys. J. **330**, L33 (1988).
[224] M. T. Ginn, K. Sellgren, E. E. Becklin, and D. N. B. Hall, Astrophys. J. **338**, 824 (1989).
[225] M. E. Bailey, Mon. Not. R. Astron. Soc. **190**, 217 (1980).
[226] D. A. Allen, in Ref. 214, p. 1.
[227] L. M. Ozernoy, in Ref. 214, p. 181.
[228] O. E. Gerhard, Mon. Not. R. Astron. Soc. **232**, 13P (1988).
[229] J. Goodman and H. M. Lee, Astrophys. J. **337**, 84 (1989).
[230] G. H. Rieke, M. J. Rieke, and A. E. Paul, Astrophys. J. **336**, 752 (1989).
[231] L. M. Ozernoy, Observatory **96**, 67 (1976).
[232] R. C. Peterson, P. Seitzer, and K. M. Cudworth, Astrophys. J. **347**, 251 (1989).
[233] M. J. Duncan and J. C. Wheeler, Astrophys. J. Lett. **237**, L27 (1980).
[234] J. Binney and G. A. Mamon, Mon. Not. R. Astron. Soc. **200**, 361 (1982).
[235] A. D. Dressler and D. O. Richstone, Astrophys. J. **348**, 120 (1990).

Translated by M. E. Alferieff