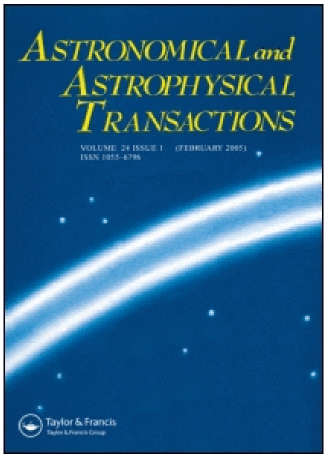

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INTERACTION OF GALAXIES AND ACTIVITY PROBLEM

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(16 August 1991)

In this brief review we regard the activity of galaxies as a result of their interaction.† One of the main mechanisms of galaxy activity seems to be compensation of momentum under the coalescence of spiral galaxies. As a result, some part of the disk gets an opportunity to fall to the nucleus.

On this occasion the mass and momentum distribution is obtained as a solution of a generalized Smoluchowsky kinetic equation (KE) describing the coalescence of galaxies. The luminosity function of active galaxy nuclei is calculated on the assumption that activity is governed by accretion during the merging of spiral galaxies.

KEY WORDS Galaxies, interaction, activity, kinetic equation, angular momentum, luninosity function.

1. INTRODUCTION

"The past decade has brought much new evidence that galaxies do not evolve in splendid isolation, but instead interact with each other. Collisions among them are more frequent than one would expect. . ." This is a quotation from the first of two invited talks on the theme "Interacting Galaxies" on the 173rd meeting of the American Astronomical Society in January 1989 (Schweizer, 1988; Toomre, 1988). The galactic nuclei activity leading to the phenomena of Seifert galaxies, radiogalaxies and quasars depending on their activity scale is closely connected with their interactions (see reviews by Balik and Heckman, 1982; Heckman, 1989, for example, see also the list of phenomena and references in Fuentes-Williams and Stoke, 1988).

In recent years definite confirmation has been received (Toomre, 1977) of the point of view that the activity is connected with accretion due to angular momentum compensation when spiral systems merge. Several arguments were given by Komberg (1984, 1989) in favour of the fact that quasars are objects of the second generation formed due to coagulation of less massive objects. In accordance with data from Hutching's (1983) review, nearly 30% of galaxies belonging to quasars (i.e. to their nuclei) are in a state of interaction (collision)

† In this text in contrast to the report at the conference we omit the arguments (based on observations) in favour of the significant role of galaxy interactions in their activity, addressing readers to the cited reviews.

with a neighbouring galaxy. Pictures of the brightest infrared sources in the IRAS catalogue illustrate a phase of merging of all objects (Sanders, Soifer *et al.*, 1988).

The dependence of galaxy morphology on density of their surroundings which affects luminosity function can be an extra argument there. The necessity or possibility of coagulation processes appears in theoretical models of galaxy formation (Shandarin, Doroschkevich and Zeldovich, 1983; Quinn, 1990).

Further we shall restrict oursleves to some statistic consequences of the merging process which gives a joint distribution function in masses and momenta of galaxies when the suggestion that collision of galaxies is inelastic but mass and rotational momentum (spin) are conserved in the system is made. We will disregard probability dependence on spin moment. Also we will from the present ignore the role of orbital momentum of pair impacting galaxies taking into account the relatively small value of impact parameter (arm moment) in the case of merging. Of course here we make not only qualitative errors but also overlook effects such as the appearance of rotational momentum in off-center collisions of elliptical galaxies, see Barausov, Ushakov, Chernin (1988).

But in this assumption we can assert our problem in the closed form for kinetic equation (KE) in describing coagulation with regard for momentum and mass conservation and we have succeeded in solving it under certain limitations (Kats and Kontorovich, 1989; 1990).

2. THE KINETIC EQUATION OF COAGULATION WITH MASS AND MOMENTUM CONSERVATION

Let us consider the mass and momentum galaxy distribution function $f(M, \mathbf{S}; t)$ to be governed by the generalized Smoluchowsky KE which describes particles (here we have galaxies) coalescence with mass M and rotational momentum $\mathbf{S}$ (classical variant of spin) being conserved. This KE has the form:

$$\frac{\partial f}{\partial t} = \int \mathrm{d}M_1\,\mathrm{d}M_2\,\mathrm{d}\mathbf{S}_1\,\mathrm{d}\mathbf{S}_2[U\delta_M\delta_{\mathbf{S}}f_1f_2 - \curvearrowright - \curvearrowright\curvearrowright]. \tag{1}$$

Here the transition probability $U\delta_M\delta_{\mathbf{S}}$ contains two δ-functions which express the conservation laws in the coagulation process $\delta_M \equiv \delta(M - M_1 - M_2)$, $\delta_{\mathbf{S}} \equiv \delta(\mathbf{S} - \mathbf{S}_1 - \mathbf{S}_2)$, indices in coefficient U are omitted: $U \equiv U_{M\mathbf{S}|M_1\mathbf{S}_1M_2\mathbf{S}_2}$, shortened designations are used for: $f_1 \equiv f(M_1, \mathbf{S}_1; t)$ and so on; arrows $\curvearrowright$ and $\curvearrowright\curvearrowright$ as in (1) denote cyclic rearrangements of indices. At first, let us consider the factor of coagulation U as a constant. (Similar investigation is possible for all probabilities in which coefficient U does not depend on momentum and the mass dependence is such that Smoluchowsky KE is integrable.) Making use of Laplace transformation over mass and Fourier over spin (in short it is simply Fourier) $f(M, \mathbf{S}; t) \rightarrow \mathscr{F}(p, \mathbf{q}; t)$ and adding the source-function we will get KE for $\mathscr{F}$ in the form

$$\frac{\partial \mathscr{F}(p, \boldsymbol{q}; t)}{\partial t} = U\mathscr{F}^2(p, \mathbf{q}; t) - 2U\mathscr{F}(p, \mathbf{q}; t)n(t) + \mathscr{D}(p, \mathbf{q}; t), \tag{2}$$

where $n(t) \equiv \mathscr{F}(0, 0; t)$ is a "concentration" of galaxies.

3. THE STATIONARY SOLUTION

An asymptotic solution for large values of M and S is sufficiently universal and is defined by contribution of small p and q in

$$\mathcal{F}_{st}(p, \mathbf{q}) = n_\infty[1 - (M_0 p + i\mathbf{S}_1\mathbf{q} + S_2^2 q^2/2)^{1/2}], \tag{3}$$

where $n_\infty = \sqrt{D_0/U}$, $D_0 \equiv \mathscr{D}(0, 0)$ and M_0, $\mathbf{S}_1$ and S_2^2 are the characteristic source mass, momentum and its square. Note the change in notation; in (3) and corresponding occasions below we will use S_2^2 instead of $2S_2^2$ in terms of our papers (1989, 1990). Conversion in (3) over p is determined by branch point, then

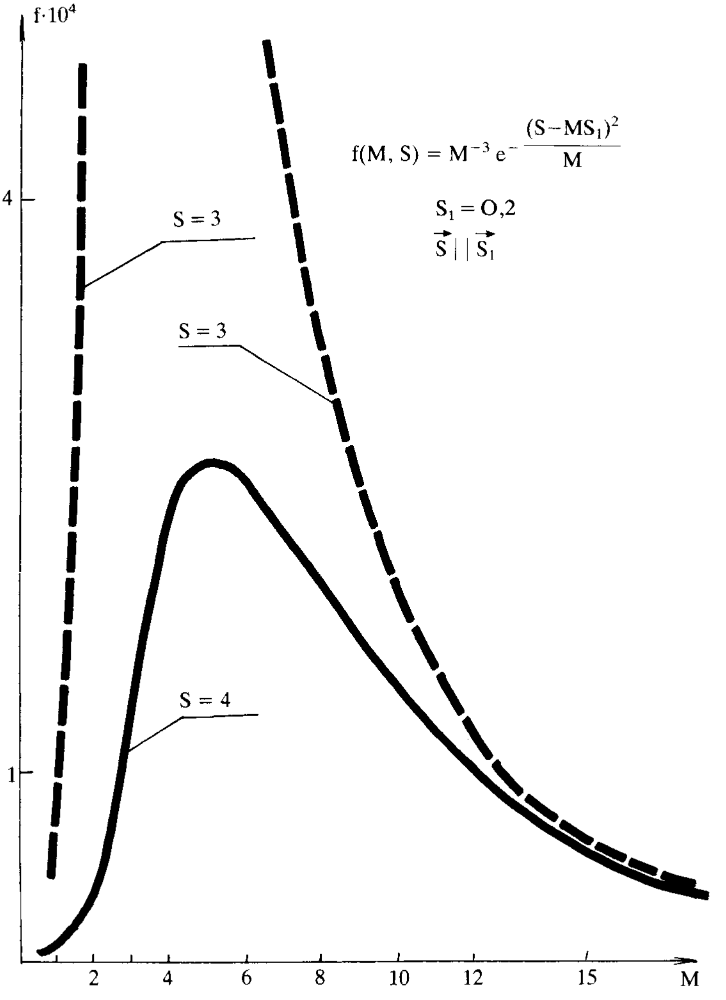

Figure 1 The mass distribution for the fixed momentum in the stationary case. The maximum value depends on S and $M \propto S^2$. The masses are measured in M_0 units and momenta are in $\sqrt{2}S_2$.

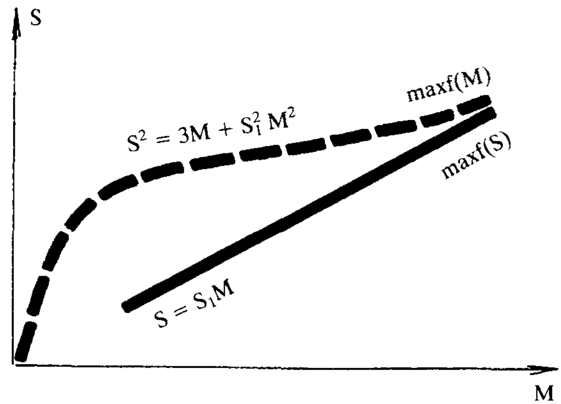

Figure 2 Lines of $f(M, \mathbf{S})$ maxima. For fixed M (solid line) and $\mathbf{S}$ (dotted line), and anisotropic distribution for $\mathbf{S} // \mathbf{S}_1$.

conversion over $\mathbf{q}$ is turned to Gauss integral calculation. So we obtain the result

$$f_{st}(M, \mathbf{S}) = f_{st}(M) g\left(\mathbf{S} - \mathbf{S}_1 \frac{M}{M_0}, \frac{M}{M_0} S_2^2\right), \tag{4}$$

$$g(\boldsymbol{\alpha}, \beta) \equiv (2\pi\beta)^{-3/2} \exp(-\alpha^2/2\beta), \tag{5}$$

$$f_{st}(M) = (J/4\pi U)^{1/2} M^{-3/2}. \tag{6}$$

Here $J = M_0 D_0$ is the mass flux along the spectrum defined by the source, the $g(\boldsymbol{\alpha}, \beta)$ function defines momentum distribution for fixed mass and is normalized to unify. $f_{st}(M)$ is the stationary mass function (see Figures 1, 2). Integration (4) over masses leads to the momentum distribution $f_{st}(S) = n_\infty S_2/(\sqrt{2}\,\pi^2 S^4)$ in the isotropic case.

On the other hand the integration (4) over momenta gives us $f_{st}(M) \propto M^{-3/2}$ what is close to the observed value for cosmic clouds (see review by Elmegreen 1990 and below).

4. EVOLUTION OF INITIAL DISTRIBUTION

Let us examine another assertion of our problem when there is not a permanently acting source but only an initial distribution which we will take in one-modal form with characteristic mass scale $\bar{M}$ some anisotropy $\bar{\mathbf{S}}$ and distribution function in momenta with width $(\overline{S^2})^{1/2}$. In the case of large values of masses and momenta our solution may be found from expansion of initial function at $p, \mathbf{q} \to 0$. Then integral over p is equal to the residue in the pole $p = -(\tau^{-1} + i\mathbf{q}\mathbf{S} + \frac{1}{2} q^2 \overline{S^2})$ and after that Gauss integral over q appears. So we have

$$f(M, \mathbf{S}; t) = f(M, t) g\left(\mathbf{S} - \bar{\mathbf{S}} \frac{M}{\bar{M}}, \frac{M}{\bar{M}} \overline{S^2}\right), \tag{7}$$

$$f(M, t) = n_0 (\bar{M} \tau^2)^{-1} \exp(-M/\bar{M}\tau). \tag{8}$$

Here $\tau = U n_0 t$, n_0 is the initial concentration, g is defined in (5). In the isotropic case ($\bar{\mathbf{S}} = 0$) at the fixed $\mathbf{S} \neq 0$ the mass function has a maximum, which decreases

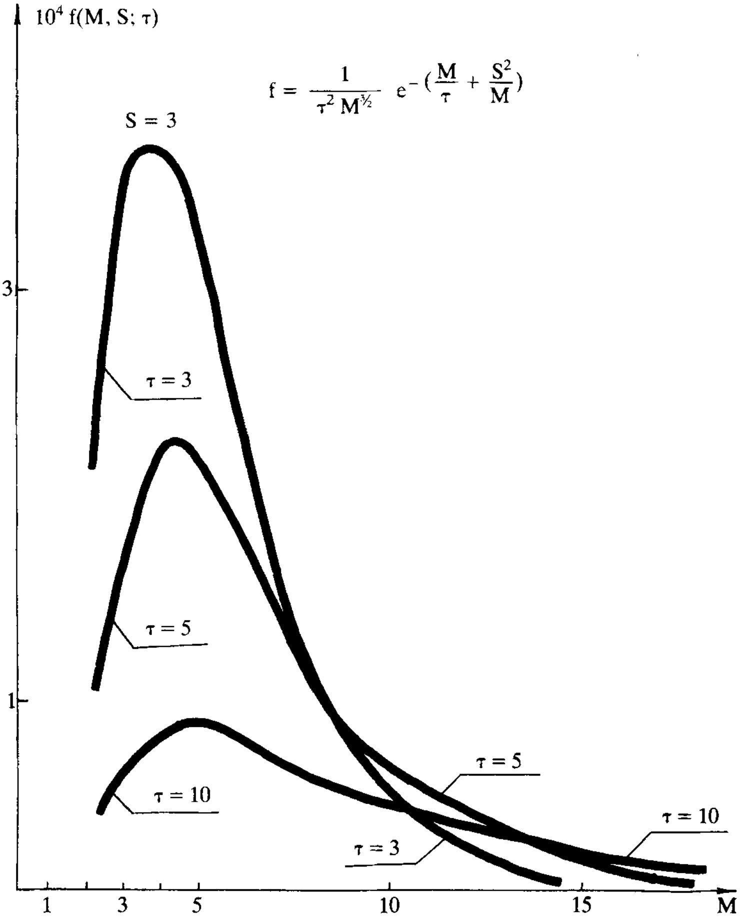

Figure 3 The evolution of initial distribution localized at small mass region in the isotropic case.

and shifts towards the large masses when τ is increasing accordingly to $M \propto \tau$, Figure 3.

5. COSMOLOGICAL EXPANSION

The number of objects in a comoving volume $a^3(t)$ is not changed when we neglect collisions so assuming the time dependence via expansion to be separated in $U: U \to \tilde{U}\chi(t)$, we can be convinced (following Silk and White (1978) that for $fa^3(t) = \tilde{f}(M, \mathbf{S}; \tilde{t})$ where

$$d\tilde{t} = \frac{a^3(t_0)}{a^3(t)} \chi(t)\, dt$$

the KE is represented in the form (1) with time-independent coalescence coefficient $\tilde{U}$. That is why we get the solution of the "initial" problem from (8) by exchanging $n_0 \to n_0 a^3(t_0)/a^3(t)$, $\tau \to \tilde{\tau} - \tilde{\tau}_0$ where $\tilde{\tau} \equiv \tilde{\tau}(t) = n_0 \tilde{U}\tilde{t}(t)$, $\tilde{\tau}_0 = \tilde{\tau}(t_0)$. The $\tilde{\tau}(t)$ dependence is defined by expansion law $a(t)$ and time-dependent probability $U = \overline{\sigma v}$ also. (For example it can be because of the average velocity change during expansion.) The most interesting point is the power function of $a^3(t)/a^3(t_0) \sim (t/t_0)^\lambda$ which for $\lambda = 2$ corresponds to the Einstein–de Sitter universe, and for $\lambda = 3$ corresponds to the "empty" universe. As for $\chi(t)$, it is natural to take it in the power law form $\chi(t) \sim t^\mu$. Then $\tilde{\tau} = n_0 \tilde{U} t^{\mu+1-\lambda}/(\mu + 1 - \lambda)$ and the number of coalescences per unit time increases with t decreasing under realistic conditions of $\lambda > \mu$.

6. EVOLUTION WITH THE SOURCE INCLUSION

The solution of KE (2) may be easily analyzed for the case when the time independent source is switched on at $t = 0$. For large time $\tau = Un_\infty t \gg 1$ and $M, S \to \infty$ only small p and q are essential and using the same $\mathscr{D}(p, \mathbf{q})$ expansion as in sect. 3 (and the same designations) we obtain the solution in the form (7) with

$$f(M, t) = \frac{2\pi^2 n_\infty}{M_0 \tau^3} \sum_{l=0}^{\infty} (l + \tfrac{1}{2})^2 \exp\left[-\frac{\pi^2 M}{\tau^2 M_0}(l + \tfrac{1}{2})^2\right]. \tag{9}$$

The sum in (9) is the derivative of a Weierstrass elliptical function. When $M/M_0\tau^2 > 1$ it is enough to leave only the first term $l = 0$ that gives the exponential tail of the distribution. On the other hand when $M/M_0\tau^2 \ll 1$ ("small masses") about $(M_0/M)^{1/2}\tau$ terms are important in the sum and so it is proportional to $(M_0\tau^2/M)^{3/2}$. In this case (9) gives the stationary distribution (4) which is thus formed behind the front that shifts in the large mass region according to the law $M \sim M_0\tau^2$.

Let us in a brief way consider the problem when the source is switched off after it has been working for a long time so that the distribution (9) and its quasistationary intermediate asymptotic (6) managed to be established (the number of collisions $\tau_* = n_\infty U t_* \gg 1$). Let (9) be the initial condition at $t = t_*$. After switching off the source for a small time (so that $\tau \equiv Un_\infty t \ll 1$) the distribution at $M \gg M_0$ is practically unchanged. At $\tau \gg 1$ a "switching off" front $M \sim M_0\tau^2$ will go over the stationary part of the distribution (9), i.e. the spectrum becomes less steep because of small mass depletion and decreases as τ increases at fixed M. For $M \gg M_0\tau^2$ the distribution remains undisturbed.

Thus, if an undisturbed region of the stationary spectrum $M \gg M_0\tau^2$ is to remain after the course is switched off, it is necessary that $\tau_* \gg \tau$ i.e. the number of collisions that form the stationary distribution (τ_*) should be much larger ($\tau_* \gg \tau$) then the one after the "switching off" (Figure 4).

In the expanding Universe this can easily be the case because of rapid decreasing of the collision probability due to expansion.

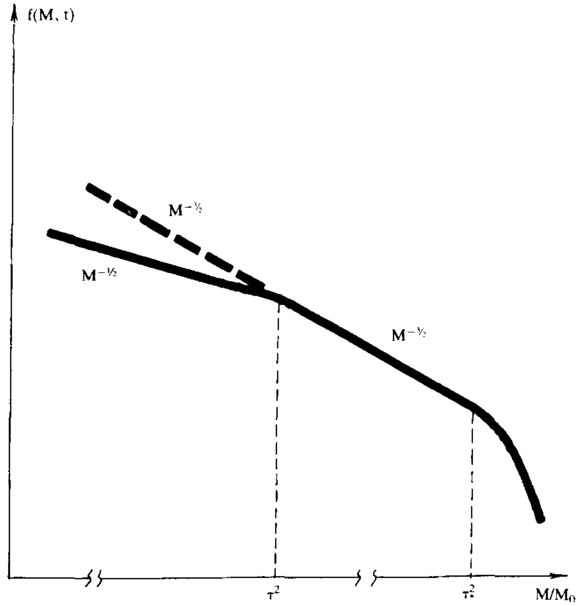

Figure 4 Mass function relaxation after switching off the source acting during finite time t_*.

7. THE SELF-SIMILAR SOLUTIONS AND SCHECHTER FUNCTION

One can see from the distribution in the presence of a source (9) that the mass function is well described asymptotically by an interpolation formula that coincides with the Schechter function used to describe the distribution in luminosities L (and masses $M \sim L$) of the field galaxies and clusters

$$f(M) = \varphi^*(M/M^*)^\alpha \exp(-M/M^*) \tag{10}$$

with $\alpha = -3/2$, $M^* = 4M_0\tau^2/\pi^2$. The Habble morphological types of galaxies can be classified over the values of reduced angular momentum $\zeta = S/m^k$. The above obtained asymptotics define the mass functions for a given ζ $f_\zeta(M, t) \equiv \int d\mathbf{S}\delta(\zeta - \mathbf{S}/M^k)f(M, \mathbf{S}; t)$ and they also have the form of Schechter functions with the parameters: $\alpha_\zeta = 3/2$, $(M_\zeta^*)^{-1} = (\bar{M}\tau)^{-1} + \zeta^2\bar{M}/2\bar{S}^2$; $\alpha = 0$, $M^* = \bar{M}\tau$ and $\alpha_\zeta = 0$, $(M_\zeta^*)^{-1} \simeq \pi^2/4M_0\tau^2 + \zeta^2 M_0/2S_2^2$ for the isotropic variants of spectra (7) or (9) respectively and $k = 1$. Following Genkin and Genkina (1973), Fall (1983), Kats, Kontorovich (1990), the k value equals $(3+\beta)/2$ (β is the power in mass-radius relation, see Appendix) which for constant density gives $k = 5/3$. This leads to $\alpha_\zeta = \alpha_S + 3k$ where α_S is the power in $f(M, S) = M^{\alpha_S}g(S^2/M)$ and the mass distribution does not take Schechter's form but is close to it.

Thus, the power law in (10) may correspond to the stationary spectrum (taking into account the above mentioned averaging over ζ) and the exponential decay to the time dependent front coalescence. It is worth mentioning that this interpretation differs from that of Silk, White, 1978 where a Schechter function with index $-3/2$ corresponds to the evolution of the initial spectrum for $U \propto (M_1 + M_2)$ (solution due to Trubnikov).

It has to be noted that the model with the source inclusion leads to the effective decreasing of the exponent in the power distribution function in comparison with $-3/2$ value (for $U = \text{const}$) due to the source obviously decreasing caused by expansion. (Compare with solutions in sect. 3 and 6.)

When $U \propto VM^u$ $(u \neq 0)$ it is possible to solve the Smoluchowsky equation for $f(M, t)$ only in special cases (see for example the paper of Trubnikov (1971) and reviews of Sofronov and Vityazev, 1983, Voloshchuk, 1984, Elmegreen, 1990). Some important information may be obtained from self-similar solutions (Voloshchuk, 1984). For example, in the source presence the corresponding self-similar mass function has for homogeneous U the following form

$$f(M, t) = f_{\rm st}(M)\psi(M/JUt^2) \tag{11}$$

where $f_{\rm st}$ denotes stationary solution $f_{\rm st}(M) \sim (J/V)^{1/2}M^{-(u+2/3)}$ and $\psi(x \to 0) \to 1$ $(M/JUt^2 \sim M^{1-u}/JVt^2)$. As it can be easily proved for U being independent of $\mathbf{S}$ the KE (1) solution has the factorized form (7) with $f(M, t)$ from Eq. (1). The g-function form in (7) follows from the simple fact that when the number n of coalescent randomly oriented spins of the order $(\overline{S^2})^{1/2}$ with masses $\simeq M_0$ is large $n \gg 1$ the characteristic squared momentum $S^2 \simeq n\overline{S^2}$ and the characteristic mass $M \simeq nM_0$. The second argument in g is independent of n while the first describes the linear growth of mean $\mathbf{S}$ value with increasing n for $\bar{\mathbf{S}} \neq 0$.

Thus we obtain

$$\begin{aligned} \alpha &= -\frac{u+3}{2}, \qquad M^*(t) = (JVt^2)^{1/(1-u)}, \\ \alpha_\zeta &= -u/2, \qquad (M^*_\zeta)^{-1} = (JVt^2)^{-1/(1-u)} + \zeta^2 M_0/S^2_2. \end{aligned} \tag{12}$$

Here we have to note that representation (11) is valid only for homogeneity index values $u < 1$ (the case $U = \text{const}$, $u = 0$ is the simplest example for $u < 1$) when distribution "slowly" shifts to the $M = \infty$ region with the front moving according to the law (11) $M \propto t^{2/(1-u)}$ or $M \propto t^{1/(1-u)}$ for the initial distribution evolution without the source. Instead of this scenario for $u > 1$ evolution has "explosive" character (see for example, Trubnikov's solution for $U \sim M_1M_2$ and other examples in Voloshchuk, 1984), i.e. the relaxation process achieves $M = \infty$ at finite moment $t_{\rm cr}$ depending on initial distribution and/or on source parameters. The shift of characteristic distribution mass in this case may be estimated as $M \propto (t_{\rm cr} - t)^{-1/(u-1)}$ (or $(t_{\rm cr} - t)^{-2/(u-1)}$ for source being included). The "explosive" case and its possible role in galaxy mass function formation are discussed in the author's and Krivitski's article.

Recently the evolution of the initial distribution localized at small mass region (with exploration of galaxy accounting observational data) was analyzed by Khersonskii, Voshchinnikov (1991), which started from coalescence probability decreasing with relative velocity increasing as its square (see Appendix, $\xi = 1$). In this case the U homogeneity power equals 4/3 ($\beta = 1/3$, $U \propto (M_1 + M_2)(M_1^{1/3} + M_2^{1/3})$ both for small and large masses (contact and gravitational collisions), instead of which the authors make use of the exact solution for $U \propto (M_1 + M_2)$ changing the $(M_1^{1/3} + M_2^{1/3})$ factor by its mean value. Some results are illustrated in Figure 5.

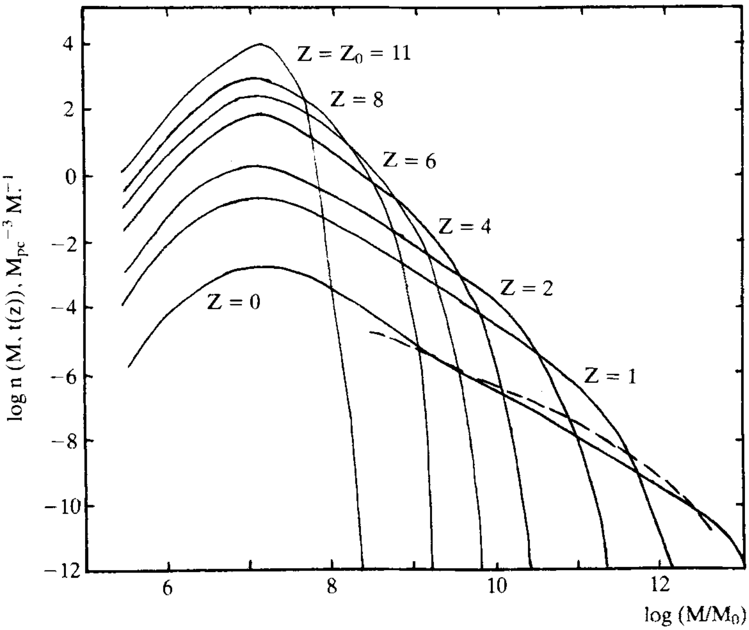

Figure 5 The galaxy mass spectra evolution from initial epoch with the red shift z_0 to present one, $z = 0$ for $\xi = 1$ in accordance with Khersonskii, Voshchinnikov (1991) (solid line). The dotted line corresponds to the observational galaxy mass spectrum.

8. COALESCENCE-INDUCED ACTIVITY

Therefore we will further proceed from the viewpoint of collision nature of activity which makes it possible to apply the kinetic model formulated above for calculating the luminosity function (LF) of quasars for whch reliable data are available (Koo and Kron, 1988; Boyle *et al.*, 1988†).

Whereas LF of galaxies reflects their mass distribution, the luminosity distribution of active objects (Seyfert galaxies, radio galaxies, quasars) is immediately connected with the mechanism of activity: accretion to a compact object in the center of the galaxy. In this case, according to the Toomre hypothesis (1977), the angular momentum compensation during the galaxy merger is responsible for the mass falling at the center. We will calculate the LF of active objects (quasars) in terms of this scheme using the galaxy distribution in masses and angular momenta obtained above, confining ourselves to the simplest, from the viewpoint of calculation, case of ultimately anisotropic momentum distribution of galaxies at large masses.

† For another way of solving the problem see for instance, in Roos (1985), De Robertis (1985), Carlberg (1990). Note also the recent reviews by Stocton (1990), Fricke and Kollatshny (1989), Kennicutt (1991) and Schlosman (1991), which we have not seen yet so references are based on the preprint by Heckman (1990).

9. ACTIVITY MODEL AND LUMINOSITY FUNCTION OF QUASARS

The model under consideration relates the fractions of excessive mass of the disk Δm capable of making its way to the center to masses M and momenta $\mathbf{S}$ of colliding galaxies, see Figure 6. The real problem complexity is so high (see the state of the problem in Hernquist's review article) that we will confine ourselves to the simplest phenomenological model. The latter takes into account only the conservation laws at merging $M = M_1 + M_2$, $\mathbf{S} = \mathbf{S}_1 + \mathbf{S}_2$ and the relation between the exececessive mass and the disk masses: $\Delta m = m_1 + m_2 - m$ where the disk mass $m = (S/M^{2/3})(\rho^{2/3}/\sqrt{G\rho})$ ($m \ll M$, ρ is the density). The luminosity L of an active object (quasar) is proportional to $\Delta m : L = B\Delta m$, $B = \varepsilon\eta c^2/t_{\mathrm{ac}}$ where η is the excessive mass portion actually making its way to the center, t_{ac} is the accretion time ($\Delta\dot{m} \sim \Delta m/t_{\mathrm{ac}}$, ε is the process efficiency).

For the distribution in luminosities or, which is the same, in Δm (with an accuracy to $t_{\mathrm{ac}}(L)$ dependence that can be readily taken into account), we will start from the equation

$$\dot{f}(\Delta m) = I_{\Delta m} - f(\Delta m)/t_{\mathrm{act}},$$

$$I_{\Delta m} = \int \mathrm{d}\tau f(M_1, \mathbf{S}_1) f(M_2, \mathbf{S}_2)\delta_M \delta_{\mathbf{S}} \delta_{\Delta m} U \tag{13}$$

$$\delta_{\Delta m} \equiv \delta[\Delta m - (m_1 + m_2 - m)], \qquad \mathrm{d}\tau = \mathrm{d}M\,\mathrm{d}M_1\,\mathrm{d}M_2\,\mathrm{d}\mathbf{S}\,\mathrm{d}\mathbf{S}_1\,\mathrm{d}\mathbf{S}_2.$$

The integral term $I_{\Delta m}$ in (13) describes the rate of active object formation with the mass excess Δm due to merges, the second term being their rate of decrease due to radiation and so on. The coefficient $U = \overline{\sigma v}$, where σ is the cross-section and v is the relative velocity. The joint functions of mass and momentum distribution $f(M, \mathbf{S})$ will be considered as the known ones (either from observations or as solutions of appropriate kinetic equations).

In the stationary case, it follows from (13) that

$$f(\Delta m) = t_{\mathrm{act}} I_{\Delta m} \tag{14}$$

which will be compared with the stationary part of LF in quasars.

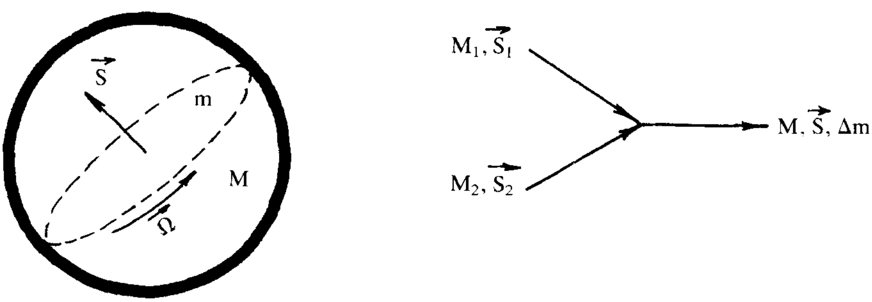

Figure 6 Model representation of the galaxy with mass M, angular momentum $\mathbf{S}$ and disk mass m and the merger process leading to the disk mass excess Δm.

10. LF AT STRONG ANISOTROPY OF MOMENTUM DISTRIBUTION

As it follows from sufficiently general considerations, the momentum distribution $f(M, \mathbf{S})$ for masses that are large compared with the characteristics mass M_0 of the source (or of the initial distribution) takes the Gaussian form named above as g (see Eq. (5) and discussion after (11)).

At sufficiently large masses $(S_0(M/M_0)^{1/2} \gg (\overline{S^2})^{1/2})$ the momentum dispersion is small and the function g becomes so sharp that it can be substituted by the δ-function

$$g\left(\mathbf{S} - \frac{M}{M_0}\mathbf{S}_0, \frac{M}{M_0}\overline{S^2}\right) \to \delta\left(\mathbf{S} - \frac{M}{M_0}\mathbf{S}_0\right), \qquad \overline{S^2} \to 0, \tag{15}$$

The sharpness of distribution (15) makes it possible to integrate (13) over the momenta. At $U = \text{const}$, $\rho = \text{const}$, $\alpha + 1 < 0$

$$I_{\Delta m} \propto \frac{(\Delta m)^{6(1+\alpha)}}{\Delta m}. \tag{16}$$

According to the data reported by Koo and Kron (1988), the exponent of stationary† power-type section slope $\phi(L) \sim L^{\gamma}$ of LF in quasars (at the weak end) is equal to $\gamma = -1$. According to Boyle *et al.* (1988), $\gamma = -1.4 \pm 0.2$. The possible minor deviation from -1 is not very relevant for us.

The known exponent of the Schechter function $\alpha = -1.25$ (see review by Gorbatsky, 1986) is not sufficiently close to the required one by virtue of the rather sharp dependence on α in (16).

It is important, however, that this value of α is determined by the data obtained for clusters. In his recent work Tully (1988) (see also the review article by Bingelli, Sandage and Tammann 1988) gives the value $\alpha = -1.0$ for field

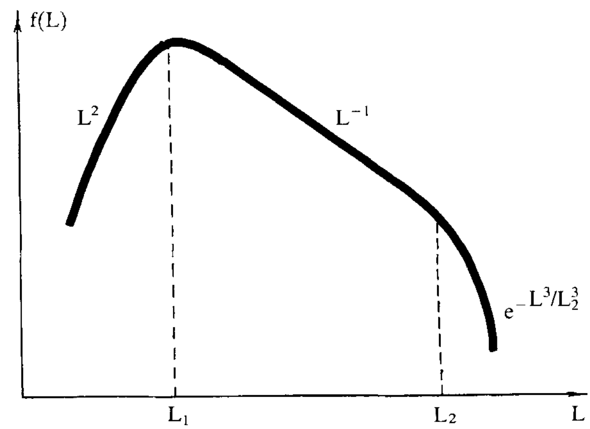

Figure 7 Stationary luminosity function of quasars corresponding to *MF* exponent of galaxies $\alpha \simeq -1$. The exponents in lateral regions are model-dependent.

† If the observed LF evolution is really of a radiation character at the bright end of the previously established stationary LF.

galaxies (in the present epoch). Thus, we obtain the correspondence between the LF index of quasars $\gamma = -1$ and the MF index of field galaxies $\alpha = -1$, see Figure 7.

The other approach to this problem is based on the explosive scenario (see author's article with Krivitski 1992). Note that if the galaxy mass spectrum is also formed as a result of coalescence (as it is for the explosive scenario, for example), then the essential increasing of the number of collisions is needed in comparison with gas-kinetic estimation (for instance, due to (local) increasing of galaxy concentration).

APPENDIX

THE COALESCENCE COEFFICIENT FOR GRAVITATIONAL AND CONTACT COLLISIONS

We will use for the cross-section σ the simplest form

$$\sigma = \pi R^2\left(1 + \frac{2GM}{Rv^2}\right)\varphi, \qquad \varphi = \left(1 + \frac{2GM}{Rv^2}\right)^{-\xi},$$

where $2GM/Rv^2$ is the gravitational focusing factor, $R = R_1 + R_2$, $M = M_1 + M_2$, $\mathbf{v} = \mathbf{v}_1 - \mathbf{v}_2$; $R_{1,2}$, $M_{1,2}$, $\mathbf{v}_{1,2}$ are radii, masses and velocities of colliding objects, and φ represents the conditional probability of coalescence for the head-on collisions. The coalescence coefficient $U = \overline{\sigma v}$ where the bar denotes average according to velocity distribution

$$\overline{\cdots} \equiv \int \mathrm{d}\mathbf{p}_1\, \mathrm{d}\mathbf{p}_2 \cdots \chi(\mathbf{p}_1/M_1^{\alpha})\chi(\mathbf{p}_2/M_2^{\alpha}).$$

The integration is carried out by moments $\mathbf{p}_{1,2}$ of colliding particles, χ is normalized velocity distribution function ($\int \mathrm{d}\mathbf{p}\chi(\mathbf{p}/M^{\alpha}) = 1$), index α equals 1/2 for Maxwellian distribution and unity for Lynden–Bell distribution, see Vinokurov *et al.* (1985). (Do not confuse the α here with Schechter's index).

Both for small (contact collisions) and large (gravitational collisions) average gravitational focusing parameter values $\gamma \equiv \overline{2GM/Rv^2} \equiv v_g^2/\overline{v^2}$ the coefficient U becomes a homogeneity mass function (in symbolic form $U \sim M^u$). The homogeneity power u for $R \propto M^{\beta}$ equals†

$$u = \begin{cases} 2 + \beta - \alpha & \gamma \gg 1 \\ 2\beta - 1 + \xi(3 - \beta) + \alpha(1 - 2\xi) & \gamma \ll 1, \quad \xi < 2 \\ 5 - 3\alpha & \gamma \ll 1, \quad \xi > 2. \end{cases}$$

The explicit U form for $\alpha = 1$ is

$$U \sim \pi R^2 v_g \gamma^{3/2} I(\gamma)$$
$$I(\gamma) = \int_0^{\infty} \mathrm{d}x x\left(1 + \frac{1}{x}\right)\varphi(x)\mathrm{e}^{-\gamma x}$$

† There is a mistake in Appendix of the author's work (1990): the restriction $\xi < 2$ was not indicated so the correct result for $\xi > 2$ was absent.

and leads for φ being a step function (it is equivalent to $\xi \gg 1$) to the following U asymptotics:

$$U = \pi R^2 v_g \gamma^{1/2} \propto (M_1 + M_2)(M_1^\beta + M_2^\beta), \qquad \gamma \gg 1$$

for "large" masses ("cold" galaxy system) and

$$U = \pi R^2 v_g \gamma^{3/2} = C(M_1 + M_2)^2, \qquad C \sim G^2/\bar{v}^3, \qquad \gamma \ll 1$$

for "small" masses ("hot" galaxy system). The radii in the last case vanish in the final result as it has to be for $\xi > 2$. Note that we have omitted numerical factors of the order of unity.

The above expressions for $\xi > 2$ may be obtained in the following (not perfectly exact) way: the coalescence condition means that the total energy of the new system has to be negative and if we assume all members of a collision to be in virial equilibrium then it leads to the $v < v_g$ for the relative velocity, i.e. to the φ-function of the step form: $\varphi \propto \theta(v_g - v)$ (see Khersonskii and Voschchinnikov (1990), for example). The more detailed form leads to the condition which contains not only the masses sum but also the reduced mass and as a consequence of the more complete mass dependence what manifests, for example, in U asymptotics.

The main problem consists in the necessity to take into account the tidal interaction which may transfer the galaxies from hyperbolic to elliptical orbit and also ensure the coalescence from the last. Qualitatively this was clear even for Holtzmark (see Tremain's (1982) review) who was the pioneer of investigations in this area, but the difficulties are not solved just yet (see the review by Alladin and Narasimhan (1982) and Saslaw's (1987) monograph).

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