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On the formation of elliptical galaxies*

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It is suggested that elliptical galaxies are formed in an expansion from a steady-state situation in which the mean density is of order 10^{-8} g/cm^3. Such an expansion is possible because the inhomogeneous steady-state theory has an instability in which the creation process is essentially cut-off, and in which expansion proceeds according to the Einstein–de Sitter law. The characteristic mass of the 'observable universe' at the onset of such an instability is *ca.* $10^{13} M_{\odot}$, and this is taken to set the upper limit to the masses of galaxies.

We suggest that the condensation of elliptical galaxies depends on the presence of inhomogeneities, in particular that a galaxy is formed around a central mass concentration. Because the Einstein–de Sitter expansion law is the limiting case between expansion to infinity at finite velocity and a full-back situation, in which expansion stops at some minimum but finite density, a central condensation with mass appreciably less than that of the associated galaxies suffices to prevent continuing expansion. A mass of $10^9 M_{\odot}$, for example, will restrain a total mass of *ca.* $10^{12} M_{\odot}$ from expanding beyond normal galactic dimensions.

The emissivity, assumed to follow the mass distribution, can readily be calculated and is found to follow an $r^{-\frac{8}{3}}$ law, which results in a dependence on $r^{-\frac{5}{3}}$ when projection against the sky is considered. This law applies outside a central region with radius of order 30 parsecs. It is close to Hubble's law $(1+r/a)^{-2}$, with a an adjustable parameter, but actually seems to fit the observed $I(r)$ curves better than Hubble's law.

These considerations apply to a spherically symmetric case. Slight deviations from a spherically symmetric expansion with respect to the central object can be discussed by introducing a local rate of strain tensor, ϵ_{ij}. The resulting galaxy then has the shape of an ellipsoid with principal axes in the directions of the principal axes of ϵ_{ij}. In the special case where two axes are equal, the spheroid can be either prolate or oblate. Unless rotation is set up by a subsequent accretion of gas, the elliptical galaxies are not rotating. The ellipticities of the isophotal contours of the projected image should be closely constant, with a very slight increase outwards. It is shown that cases of high ellipticity must be comparatively rare.

The origin of spiral galaxies, and the possibility of there being mixed spiral and elliptical forms, is briefly discussed.

The mass distribution

The cosmological basis of the present paper has been discussed in the preceding paper (Hoyle & Narlikar 1966). We shall not therefore be concerned here with the cosmological aspects of the theory. We shall accept that the Universe, or a portion of it, expands from an initially steady-state situation with $\rho \simeq 10^{-8}$ g/cm^3, $H^{-1} \simeq 10^{18}$ cm, that creation is effectively zero during this expansion, and that the Einstein–de Sitter expansion law holds in first approximation.

The Newtonian analogue of the Einstein–de Sitter law is given by

$$\dot{r}^2 = 2GM/r, \tag{1}$$

in which r is the radial coordinate of an element of material defined by the condition that, in a spherically symmetric situation about $r = 0$, the mass interior to r is M. For a given sample of material M remains constant and $\dot{r} \to 0$ only as $r \to \infty$. Equation (1) is an integral of the second order Newtonian equations, and the fact

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that no constant of integration appears represents the analogue of the Einstein–de Sitter law.

Next, consider the Newtonian problem of an object of mass μ placed at the origin $r = 0$, all conditions for a particular element of the cloud being the same as before at a particular moment. Denote the value of r at this moment by r_0. Then $\dot{r}$ at this moment is $(2GM/r_0)^{\frac{1}{2}}$, as before, and the subsequent motion of the element in question is determined by

$$\dot{r}^2 = \frac{2G(M+\mu)}{r} - \frac{2G\mu}{r_0}. \tag{2}$$

The outward velocity drops to zero, and the element subsequently falls back toward $r = 0$. The maximum radial distance $r_{\text{max.}}$ reached by the element is given by

$$r_{\text{max.}} = \{1 + (M/\mu)\}\, r_0, \tag{3}$$

and for sufficiently large M/μ, $r_{\text{max.}} \simeq Mr_0/\mu$, so that the fractional increase, $r_{\text{max.}}/r_0$, above the radius r_0 at which the element had the same radial motion as in the Einstein–de Sitter case, is just M/μ. This factor is larger for elements more distant from μ than for the inner parts of the cloud, so the outer parts recede proportionately further than the inner parts.

What determines the particular moment at which the Einstein–de Sitter condition, $\dot{r} = (2GM/r)^{\frac{1}{2}}$, holds for any particular sample of material? To come to grips with this important question we must consider the relativistic formulation of the problem. Again assume spherical symmetry, and use T, R as time and radial coordinates. The line element of the Einstein–de Sitter cosmology is

$$\mathrm{d}s^2 = \mathrm{d}T^2 - S^2(T)\,(\mathrm{d}R^2 + R^2\,\mathrm{d}\Omega^2), \tag{4}$$

with $S \propto T^{\frac{2}{3}}$. Transform to locally flat space-time for an observer at $R = 0$, and let t, r be the new time and radial coordinates.* Locally, the line element is

$$\mathrm{d}s^2 = \mathrm{d}t^2 - \mathrm{d}r^2 - r^2\,\mathrm{d}\Omega^2 + o_4. \tag{5}$$

Since the angular coordinates used need not be changed in the transformation, $r = RS$, so that the observer at $R = 0$ is the same as the observer at $r = 0$. Local gravitational effects are contained in the o_4 terms of (5). Their main effect can be written explicitly as

$$\mathrm{d}s^2 = \mathrm{d}t^2(1 - 2GM/r) - \mathrm{d}r^2 - r^2\,\mathrm{d}\Omega^2, \tag{6}$$

in which M is the total mass interior to r. For a uniform cloud, M is proportional to r^3, and the term in M is of order $r^2\,\mathrm{d}t^2$, so that for r of order $\mathrm{d}t$ this is o_4. The line element (6) leads in first approximation back to the Newtonian situation. And if we wish to display the effect of a mass μ at $r = 0$ we can write

$$\mathrm{d}s^2 = \mathrm{d}t^2[1 - 2G(M+\mu)/r] - \mathrm{d}r^2 - r^2\,\mathrm{d}\Omega^2, \tag{7}$$

in which M is now, not the total mass interior to r, but the cloud mass interior to r.

* We are here using coordinates T, R for the cosmological situation and t, r for the local situation. This is the opposite way round to our usage in other papers. We prefer t, r as working coordinates. Previously we have been working in the cosmological situation, now we are working in the local situation.

ıplete solution of a local gravitational problem can be represented as a ›ries in the dimensionless parameter $2G(M+\mu)/r$, which must be $\ll 1$, this ıat we mean by a 'local problem'. The Newtonian solution is of course the n in this series. However, it is clear that we cannot use the Newtonian solution for the effect of μ if the second order term in $2GM/r$ exceeds the first order term in $2G\mu/r$, as is possible when $\mu/M \ll 1$. Hence the Newtonian equations for the effect of μ, (2) and (3), cannot be used unless the moment for which we use $r \equiv r_0$, $\dot{r} = (2GM/r_0)^{\frac{1}{2}}$, is such that

$$\frac{2G\mu}{r_0} \geqslant \left(\frac{2GM}{r_0}\right)^2. \tag{8}$$

By taking the equality sign in (8) we do indeed define a particular value of r, corresponding to a specified M, namely,

$$r_0 = (M/\mu)\, 2GM. \tag{9}$$

The situation is that the Newtonian calculation for the effect of μ can be applied to the subsequent motion of an element of material such that the specified M lies interior to it. But can we use $(2GM/r_0)^{\frac{1}{2}}$ as the starting velocity in this calculation? Not in general, because in general the cloud will have at least small fluctuations from the Einstein–de Sitter expansion. We shall deal with such fluctuations in the following section, confining ourselves here to the case in which the conditions $r \equiv r_0$, $\dot{r} = (2GM/r_0)^{\frac{1}{2}}$, with r_0 given by (9), hold for all M.

From (3) and (9) we have

$$r_{\text{max.}} \simeq \frac{M}{\mu} r_0 \simeq 2GM\left(\frac{M}{\mu}\right)^2. \tag{10}$$

This result has a number of interesting consequences. Set $r_{\text{max.}}$ equal to a typical galactic radius, $r_{\text{max.}} = 3 \times 10^{22}$ cm.

Then (10) leads to

$$\frac{M}{M_\odot} \simeq 5 \times 10^5 \left(\frac{\mu}{M_\odot}\right)^{\frac{2}{3}}. \tag{11}$$

A central object of mass $\mu = 10^9 M_\odot$ gives $M = 5 \times 10^{11} M_\odot$, while $\mu = 10^7 M_\odot$ gives $M = 2 \times 10^{10} M_\odot$. It is of interest that the central condensations present in massive elliptical galaxies are known to be of order $10^9 M_\odot$ and that the total masses are believed to be $\sim 10^{12} M_\odot$.

Suppose that during expansion stars are formed from the gas. The stars will continue to occupy the full volume corresponding to their maximum extension from the centre, so that the mass of the stars interior to r is given by setting $r_{\text{max.}} = r$ in (10). Numerically, we have

$$\frac{M(r)}{M_\odot} \simeq 2 \times 10^5 \left(\frac{\mu}{M_\odot}\right)^{\frac{2}{3}} r^{\frac{1}{3}}, \tag{12}$$

in which r is in kiloparsecs. Evidently, the mean star density at distance r from the centre is proportional to M/r^3, i.e. to $r^{-\frac{8}{3}}$. So long as the stars have everywhere the same luminosity function the emissivity per unit volume at distance r is proportional to $r^{-\frac{8}{3}}$. This determines the light distribution in a spherical elliptical galaxy.

To obtain the projected intensity distribution we first note that the above considerations can be applied to values of r beyond normal galactic dimensions. There is no upper limit to r so long as we are dealing with a single condensation. This agrees with observation, in that no ultimate maximum radius has yet been found, the conventional radii are simply those set by the sensitivity of particular observing techniques. This being so, the intensity distribution, $I(r)$, of the projected image is obtained by multiplying the volume emissivity by the factor r, and is $I(r) \propto r^{-\frac{5}{3}}$. This proportionality is slightly less steep than Hubble's law for $r \gg a$,

$$I(r) \propto (r/a+1)^{-2} \approx r^{-2}.$$

The measurements of Liller (1960) for early ellipticals E 1, E 2, E 3 give very good agreement with $r^{-\frac{5}{3}}$, better than with r^{-2}.

The $r^{-\frac{5}{3}}$ proportionality must not be used for too small r, because r_0 given by (9) becomes invalid as M is reduced towards μ. The reason is simply that if M is set too small the mean density corresponding to (9), $M/\frac{4}{3}\pi r_0^3 \propto M^{-5}$, becomes larger than the steady-state value of $\sim 10^{-8}\,\text{g/cm}^3$ from which the expansion started. Instead of (9), we then have an initial radius r_i given by

$$\tfrac{4}{3}\pi r_i^3 \rho_i = M, \quad \rho_i \simeq 10^{-8}\,\text{g/cm}^3, \tag{13}$$

and instead of (10)

$$r_{\text{max.}} \simeq \frac{M}{\mu} r_i, \quad \frac{M}{\mu} \gg 1. \tag{14}$$

In place of (12), with r now in parsecs, we have

$$r \simeq 10^{-5} \frac{M}{\mu} \left(\frac{M}{M_\odot}\right)^{\frac{1}{3}}. \tag{15}$$

As an example, for $M = 10^{11} M_\odot$, $\mu = 10^8 M_\odot$, (15) gives $r \simeq 30$ parsecs. This result is very satisfactory in that it predicts a highly concentrated point of light at the centres of elliptical galaxies.

Deviations from spherical symmetry

We shall continue, for the moment, to regard the central object as being at rest relative to the cosmological substratum, i.e. to be at rest relative to $R = 0$, with R as the cosmological coordinate in (4). Velocities relative to the cosmological substratum are known to decrease as the universe expands, according to S^{-1}. Much of the critical aspects of the formation of ellipticals occurs at the stage where the mean density is comparable to that actually found in ellipticals, say $10^{-20}\,\text{g/cm}^3$ or less. Starting from $10^{-8}\,\text{g/cm}^3$, an increase of S by 10^4 is required to reduce the initial density to this value. During this measure of expansion, any initial random motion possessed by the object μ is therefore reduced by a factor 10^{-4}. Our assumption must be a good approximation during the relevant stages of galaxy formation. We shall return to this point at the end of the present section.

The motion of the gas cloud surrounding the object will not in general correspond precisely to the Einstein–de Sitter situation assumed in the preceding section, namely $\dot{r}_0 = (2GM/r_0)^{\frac{1}{2}}$ at the appropriate r_0. There must be some velocity fluctuations, however slight. In this section we proceed to investigate the effect of such fluctuations. Once again, we shall consider the situation at the stage of expansion when the Newtonian approximation for the effect of μ becomes valid, i.e. for $r \geqslant 2GM(M/\mu)$, M being large enough for $M/\frac{4}{3}\pi r^3$ to be appreciably less than the initial density, $\sim 10^{-8}\,\mathrm{g/cm^3}$.

The deviation from the Einstein–de Sitter expansion can be represented as a superposition of a rate of strain tensor ϵ_{ij} and of a rotation about μ. Choose coordinate axes at $r = 0$ so that ϵ_{ij} is referred to principal axes, $\epsilon_{ij} = \mathrm{diag}\,(\epsilon_1, \epsilon_2, \epsilon_3)$, and consider the situation in the direction of one of the principal axes. At distance r in this direction the radial velocity is of the form

$$(2GM/r)^{\frac{1}{2}} + \epsilon r, \tag{16}$$

in which ϵ is the appropriate choice from $\epsilon_1, \epsilon_2, \epsilon_3$. Again choose a particular element of material, defined by a specified M, and let r_0 given by (9) define the moment at which $\epsilon = \epsilon_0$; the components of ϵ_{ij} are function of time, as also is the angular velocity about μ. For the moment we omit rotation. The equation determining the subsequent radial motion of our element is then

$$\dot{r}^2 = \frac{2G(M+\mu)}{r} - \frac{2G(M+\mu)}{r_0} + \left(\sqrt{\frac{2GM}{r_0}} + \epsilon_0 r_0\right)^2. \tag{17}$$

With $M/\mu \gg 1$, and with the term in ϵ_0 small, (17) can be approximately written as

$$\dot{r}^2 = \frac{2GM}{r} - \frac{2G\mu}{r_0} + 2\epsilon_0(2GMr_0)^{\frac{1}{2}}. \tag{18}$$

Radial motion ceases when r increases to

$$r_{\mathrm{max.}} = r_0 \frac{M}{\mu}\left[1 - 2\left(\frac{M}{\mu}\right)^{\frac{3}{2}} \epsilon_0 r_0\right]^{-1}, \tag{19}$$

from which a factor $r_0^{\frac{1}{2}}$ has been removed by using (9).

Going from one element of material to another, but still in the direction of the same principal axis of ϵ_{ij}, M varies and so the factor $1 - 2\epsilon_0 r_0 (M/\mu)^{\frac{3}{2}}$ varies, not only because of $(M/\mu)^{\frac{3}{2}}$, but also because both ϵ_0 and r_0 vary, r_0 in accordance with (9). Because $\epsilon_0 r_0$ is a fluctuation velocity, there must be a decrease by S^{-1}, S being the cosmological scale factor. Writing ϵ_i for the initial value of ϵ at the beginning of the expansion from the steady-state situation, and with r_i as the initial value of r, given by (13) in terms of M, we get

$$S = \frac{r_0}{r_i}, \quad \epsilon_0 r_0 = \frac{r_i}{r_0}\epsilon_i r_i. \tag{20}$$

Now ϵ_i is indeed independent of M, and with $r_i \propto M^{\frac{1}{3}}$, $r_0 \propto M^2$ from (9), $\epsilon_0 r_0 \propto M^{-\frac{4}{3}}$. We can therefore write (19) in the form

$$r_{\mathrm{max.}} = r_0(M/\mu)\,(1 - \lambda M^{\frac{1}{6}})^{-1}, \tag{21}$$

in which λ is independent of M. For the three principal directions of ϵ_{ij} we accordingly have

$$(r_{\mathrm{max.}})_i = 2GM(M/\mu)^2\,(1 - \lambda_i M^{\frac{1}{6}})^{-1}, \tag{22}$$

in which $\lambda_1, \lambda_2, \lambda_3$ are constants. Apart from the factor $(1-\lambda_i M^{\frac{1}{6}})^{-1}$, this is the same as (10).

Turning now to rotation about μ, rotation gives a velocity component perpendicular to the radial direction. The effect is to introduce a quadratic term in the rotational velocity into the equation for $\dot{r}^2$. Hence for small velocities, of the same order as the velocities introduced by e_{ij}, the effect of rotation is *second order*, whereas the rate of strain terms appear in first order. Clearly, for comparable velocity fluctuation, rotation about μ has negligible effect compared to the rate of strain tensor. We therefore omit rotation in the following discussion.

Three cases may be distinguished: (i) the λ_i terms negligible, in which case we have the previous spherically symmetric case, (ii) $\lambda_i M^{\frac{1}{6}}$ of order unity, (iii) $\lambda_i M^{\frac{1}{6}} \gg 1$ for the values of M of interest, $M \simeq 10^{11} M_{\odot}$. In the last case the strain terms dominate, and except under special circumstances, with λ_1, λ_2, λ_3 all negative, the restraining influence of the object μ is too weak to prevent a dissipation of the cloud in at least one direction. The interesting intermediate case leads to stable galaxies with a variety of forms. Negative values of the λ terms have comparatively little effect on the lengths of principal axes, but positive values can lead to a very sensitive situation, since as $\lambda_i M^{\frac{1}{6}} \to 1$, $(r_{\text{max.}})_i \to \infty$. Markedly prolate forms are perhaps more likely than markedly oblate forms, at least if the three values λ_1, λ_2, λ_3 are uncorrelated. This follows because $\lambda_i M^{\frac{1}{6}}$ will be closer to unity for one value of i than for the other two, and because of the sensitivity the corresponding principal axis will tend to be grossly exaggerated as the term in question approaches unity. Oblate forms are likely to be confined to cases where none of the λ_i terms approaches unity very closely, and where two of the terms are positive and one negative; for example $\lambda_1 M^{\frac{1}{6}} = \lambda_2 M^{\frac{1}{6}} = -\lambda_3 M^{\frac{1}{6}} = 0{\cdot}5$ gives an oblate object with axial ratio 1:3.

In general, we expect an 'ellipsoidal' form with unequal principal axes. The axial ratios depend on M through the $M^{\frac{1}{6}}$ factor. But this is an extremely weak dependence, unless one of the $\lambda_i M^{\frac{1}{6}}$ factors is close to unity. This case apart, we have an extension that varies with M closely as M^3, and this is so for each of the principal axes, giving $M^{\frac{1}{6}} \propto \sim r^{\frac{1}{18}}$. Hence we expect the isophotes of an elliptical galaxy to have very nearly the same ellipticity, the ellipticity increasing very slightly with distance from the centre.

At this stage we can readily deal with the case when the object μ possesses a small velocity relative to the cosmological substratum. Consider this effect alone, and let the velocity be in the direction x_1. For the Newtonian problem it is irrelevant whether we consider the object to be moving relative to the cloud or the cloud moving relative to the object. Taking the latter case, we see that in the x_2 and x_3 directions there are small velocity terms at right angles to the main radial motion. This is analogous to the case of rotation and gives only a second order modification to the equations of radial motion. In the x_1 direction, on the other hand, there is a first order modification, but the signs of the new terms are opposite for the positive and negative x_1 directions. The effect is similar to the rate of strain effect, except that λ must be taken to have opposite signs for the positive and negative x_1 directions. In one of these directions the radius is therefore increased by a term of the type $(1-|\lambda| M^{\frac{1}{6}})^{-1}$ while in the opposite direction the radius is reduced by

$(1+|\lambda| M^{\frac{1}{3}})^{-1}$. The net effect is to alter the total extension in the x_1 direction by the factor $(1-\lambda^2 M^{\frac{1}{3}})^{-1}$. Evidently there is a *prolate* extension in the x_1 direction, which becomes appreciable as $|\lambda| M^{\frac{1}{6}}$ approaches unity. However, as $|\lambda| M^{\frac{1}{6}}$ decreases below unity the modification becomes of second order. The rate of strain terms remain of first order as $\lambda_i M^{\frac{1}{6}}$ falls below unity, so that, unless the motion of μ relative to the cosmological substratum appreciably exceeds the velocity deviations of the cloud itself, the latter are the more important.

To conclude this section we note that in the case just considered the object μ does not lie at the geometrical centre of the resulting galaxy. Consequently the bright central concentration of light associated with the object would be expected to be displaced from the centre. In fact, one would expect the object to oscillate about the centre and instead of a point of light there could be a line distribution.

Galaxies

An immediate consequence of the above investigation is that the shapes of elliptical galaxies are not due to rotation. Any initial rotary forces would become very small after an expansion from a high density of $\sim 10^{-8}$ g/cm^3 down to galactic densities. Prolate forms should be more common than is usually supposed. Galaxies with an extreme inequality of the principal axes should be rare, because such cases depend on one of the numbers $\lambda_i M^{\frac{1}{6}}$ being only slightly less than unity.

Spiral galaxies raise a different problem. The distribution of angular momentum in the disks of spirals appears to be the same as that of a uniformly rotating cloud (Crampin & Hoyle 1964). The well-known condition for rotary forces not to exceed gravitational forces, $\omega^2 < 2\pi G\rho$, demands condensation rather than expansion, since ω^2 depends on the inverse fourth power of the linear scale, whereas ρ depends on the inverse cube. The necessary angular momentum could not be stored in a much denser cloud, whereas it can be in a more diffuse cloud.

We arrive therefore at two radically different modes of formation of galaxies, the spirals by condensation, the ellipticals by expansion. Our picture concerning the ellipticals has some features in common with that of Ambartzumian, in that we are concerned with a progression from high density to low. An important quantitative difference is that elliptical galaxies are not formed in our picture by current violent events—we are concerned essentially with a cosmological situation rather than with an astrophysical one. There is also a sense in which our point of view agrees with that proposed by evolutionary cosmologists, in that we are now associating the formation of the ellipticals with an earlier phase in the expansion of the universe, or at any rate with an earlier phase of the portion of the universe we observe, the 'bubble' in the terminology of a previous paper. However, our association with past history is radically different in that we take the process of formation as having commenced to operate at very much higher density than has hitherto been proposed. It is of course this feature of the present theory that leads to the high central concentration of the ellipticals, and to the existence of massive objects at their centres.

So far as the condensation of spirals is concerned, it is possible that, as suggested by Burbidge, Burbidge & Hoyle (1963), high energy particles play an important part in producing velocity fluctuations and local vorticity. The mean space density of matter is $\sim 3\mathscr{H}^2/8\pi G$ for the Einstein–de Sitter expansion, and with

$$\mathscr{H}^{-1} \simeq 4 \times 10^{17}\,\mathrm{s},$$

this gives a density somewhat above $10^{-29}\,\mathrm{g/cm^3}$, more than ten times the mean spatial density of luminous material. It seems likely, therefore, that ample material exists for further condensation. Why should condensation not have taken place at densities higher than $10^{-25}\,\mathrm{g/cm^3}$? Perhaps because collisional losses prevent cosmic rays from existing for a time $S/\dot{S}$ if the density is higher than this.

It is clear that condensation will proceed more readily on already existing galaxies than it will in the absence of an initial concentration of matter, unless the cosmic rays generated by an already existing galaxy produces such a large local pressure that condensation is inhibited. It is possible that this is always the situation for the very large elliptical systems that comprise the radio sources. However, elliptical-spiral mixtures could arise for less massive ellipticals and it is possible that an understanding of the rich variety of galactic forms should be sought in terms of such mixtures. Several parameters must affect the resulting forms. First, the mass ratio of the elliptical component to the disk component, the latter condensed with angular momentum from diffuse gas. Secondly, the shape of the elliptical and the relation of the shape to the direction of the angular momentum vector of the disk component. Thirdly, the length of time for which the system has evolved since the mixture was made.

The simplest case is that of condensation on to a spherical galaxy. This would be expected to yield the *Sa–Sb–Sc* series, the type being determined mainly by the mass ratio parameter. The next most frequent case is likely to be that of condensation on to a prolate elliptical galaxy—it will be recalled that on the present theory it seems more likely that 'non-spherical ellipticals' are prolate rather than oblate. When the angular momentum vector of the disk is parallel to the longest axis of the ellipsoid there must be a reduction in the length of this axis, because the gravitational field of the disk always acts to reduce stellar motions in a direction perpendicular to itself. When the disk component surrounds the elliptical the effect on stars of the elliptical that are moving parallel to the disk is comparatively less than on stars that move perpendicular to the disk. The shape of the elliptical component would therefore be changed, towards a spherical form. Indeed for a sufficiently massive disk the prolate character could even be changed to oblateness. We might expect this to happen for *Sb* and *Sc* galaxies, but perhaps not for type *Sa*, since in the latter the disk component is not so dominant. The next possibility is that the longest axis of an elliptical is perpendicular, or approximately so, to the angular momentum vector of the disk—on a random basis this case is more likely than the previous one. A surrounding disk would not destroy the prolate character of the elliptical, and it is attractive to take the view that the non-axisymmetric field of the elliptical leads to the barred sequence of spirals. Capture of gas by the gravitational field of a prolate elliptical could set the elliptical in rotation. Indeed the ultimate development of a barred spiral could require the elliptical component to develop a considerable measure of rotation.

These ideas can be put to a test by considering the special case of our own galaxy. Is our galaxy a two-component system? It is attractive to suggest that the answer may be affirmative, with the system of nucleus and halo as one component, and the disk as the other. The halo is known to be rotating slowly compared to the disk, if indeed the halo is rotating at all. Eggen (1964) in a recent catalogue of high velocity stars separates those of large ultraviolet excess (low metal content) from those of smaller excess. His plot of the stars on a Bottlinger diagram shows that the former group has little or no net angular momentum, whereas the latter group has essentially the normal rotation of the galactic disk. The eccentricities of the orbits of stars of the former group are large, showing that they dip close in to the galactic centre. Indeed, many of the stars listed by Eggen, while moving out to distances of from 15 to 40 kiloparsecs, dip to within 0·5 kiloparsecs of the centre, strongly suggesting that they have been expelled essentially from the very centre of the galaxy.

On the present picture all ellipticals must be of essentially the same age, none can have formed recently, as can be the case for some spirals. Granted that stars form by a similar process in all ellipticals, the present-day integrated colours should evidently be identical. The ages should be $\sim \frac{2}{3}\mathscr{H}^{-1}$, with $\mathscr{H}^{-1} \simeq 4 \times 10^{17}$ s, set by the present observed expansion rate. This is about 9×10^9 years. Considerations of nuclear chronology, as well as the direct calculation of stellar ages, suggest that a value $\sim 12 \times 10^9$ years would be more appropriate. A change of this amount would require the magnitudes of the brightest ellipticals in clusters to be about $\frac{3}{4}$ magnitude more luminous than is usually supposed.

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