PHYSICAL CONDITIONS IN THE CENTRAL REGION, AND THE NATURE OF THE ENGINE

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ABSTRACT

The phenomena associated with Active Galactic Nuclei raise three main astrophysical problems: (1) the nature of the primary source of energy; (2) the physical conditions within the radiation source; (3) the nature of the population evolution over cosmological time-scales.

I shall outline the links between (1) and (2), (1) and (3), that briefly go as follows. The Prime Mover is very likely to be a converter of gravitational energy in a very compact mass configuration. The associated radiation source, if it is also very compact, is so efficient and loss-dominated as to require specific conditions for the power transport and supply to the radiating particles: collisionless, effected by electromagnetic fields coherent on scales collective or macroscopic, to the point of producing at times anisotropic bulk motions. Very compact Prime Movers working at high regimes need also a compact mass supply; the output from these compound engines undergoes a characteristic change that accounts for the type of population evolutions of the associated sources.

1. THE PRIMARY ENERGY SOURCE

Active Galactic Nuclei are diverse in morphology (e.g., extended Radiogalaxies vs. pointlike Quasars) and diverse as for the patterns of their continuum radiation (e.g., typically radio-quiet Quasars vs. typical Radiogalaxies optically quiet; the OVV's and BL Lac Objects stand out for evidence of non-thermal emission in the IR-optical band; a class of peculiarly cut-off IR spectra is being observed, Rieke, Lebofsky and Wisniewski 1982).

However, the underlying energetics are similar, $E \simeq 10^{60\pm1}$ erg (Schmidt 1978, Woltjer 1978). And there are widespread, convergent indications of a very compact origin for the primary energy: variability time-scales of days to several hours are common in X-rays (Marshall,

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Warwick and Pounds 1981), common in Blazars (Angel and Stockman 1980), with substantially shorter scales noted in a few cases (Tennant and Mushotzky 1982; Wolstencroft, Gilmore and Williams 1982). At radiofrequencies, super-luminal motions and jets occur on scales $\gtrsim$ 1pc, and are likely to originate at scales $\lesssim 10^{15}$ cm (Rees, Begelman and Blandford 1981).

With scales $R \lesssim 10^{15}$ cm, powers often in excess of $L \simeq 10^{45}$ erg/s, and energies $E = Lt \approx 10^{60}$ erg, a unified gravitational origin of the primary energy is obviously attractive. As for efficiency, already at $R \lesssim 10^{15}$ cm the gravitational energy GM^2/R of a mass $M \approx 10^8 M_{\odot}$ is larger than any possible nuclear yield; the condition is $L_{45}/R_{15} > 0.3\ (\eta_*/0.01)^2/t_8$. The mass involved is $M_8 > (L_{45} t_8\ R_{15})^{1/2}$; it must be close to, if not within, its horizon at $R_H \simeq 3\ 10^{13} M_8$ cm.

Equally larger is of course the energy per particle. But with gravitation, there is also the possibility - so dramatically confirmed by the observations of Pulsars - to concentrate much energy in a single, rotational degree of freedom, and to extract high-grade power when the rotating body is magnetized or is immersed in a magnetic field. A number of specific models for extraction of gravitational energy from massive compact objects are summarized in Table 1.

2. RADIATION

Given that the primary mass converter is compact, how compact can be the radiation source? The link is provided by the quantity L/R that measures the compact ness of the primary source, expresses also the variability of the radiation source $(\Delta L \approx L)\ \Delta L/\Delta t = 3\ 10^{40}\ L_{45}/R_{15}$ erg/s^2, and finally appears in the time-scales for energy transport and for radiation emission.

The simplest source structure: in equilibrium around a Schwarzschild black hole at $R \gtrsim R_H$, emitting near the dynamic (Eddington) limit $L \leq L_E = 4\pi c G m_p M/\sigma_T = 10^{46}\ M_8$ erg/s, yields

$$\frac{L}{R} < \frac{2\pi m_p c^3}{\sigma_T} \quad \text{or} \quad \frac{\Delta L}{\Delta t} < 10^{43} \quad \text{erg/s}^2 \qquad (2.1)$$

A smaller bound is set not only by efficiency limits ($R \gtrsim$ several R_H, $\eta \lesssim 0.1$), but also by the effects of radiative transfer through the same accretion flow that powers the emission: a smaller size implies on one hand shorter crossing times R/c, but on the other hand a larger electron density n interposing to the emerging photons an optical depth $\tau = \sigma_T nR$ that slows down the effective rise-time to $\Delta t \simeq (1+\tau)R/c$. The best compromise ($\Delta t$ minimum) is struck at $\tau = 1$, and correspondingly (Fabian and Rees 1979) at any frequency the limit

$$\frac{\Delta L}{\Delta t} \lesssim \frac{\eta}{0.1}\ 2\ 10^{41} \quad \text{erg/s}^2 \qquad (2.2)$$

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would hold. A number of sources violate the limit, e.g. in the optical-IR AO 0235+164 with $\Delta L/\Delta t=6\ 10^{42}$ erg/s^2 (cf. Angel and Stockman 1980) and OJ287, $\Delta L/\Delta t=2\ 10^{42}$ erg/s (Wolstencroft, Gilmore and Williams 1982); marginally in X-rays QSO 1525+227 with $\Delta L/\Delta t=5\ 10^{41}$ erg/s^2 (Matilsky, Shrader and Tanaubaum 1982). For these sources at least, the structure has to be more complex, with radiation separate in space from power production, or with sharp anisotropy, so that the optical depth is unrelated to the mass accreted.

In the absence of anisotropy, the physical conditions in the intense compact radiation sources are stringent: a high photon density $N/n = 3(10^5/\tau)\ L_{45}/\nu_{15}R_{15}$, with $\tau \lesssim 1$ (to minimize Δt, now in the absence of the mass flow constraint discussed above), and a large probability for electron-photon collisions $\tau_{e\gamma}=\tau\ N/n$. Hence the photons tend to acquire the average electron energy, $h\nu \rightarrow \langle\varepsilon\rangle$, but not to be in a true thermal equilibrium; and the electrons undergo severe energy losses to the photons, likely departing from a Maxwellian distribution. As stressed by Rees (cf. Rees 1981), plasma components that are closer to the thermal limit or closer to non-thermal conditions (at relativistic or transrelativistic energies) may coexist, with either prevailing in each given class of sources and with much superposition.

In a thermal component at $kT \gtrsim 1$ keV, the condition $\tau_{e\gamma} \gg 1$ implies efficient X-ray radiation by Comptonization: low energy ($h\nu_o$) seed photons, on colliding with the hot electrons, undergo a systematic upgrading in energy, $\Delta h\nu/h\nu \simeq kT_e/mc^2$ per scattering; the resulting spectrum is moulded into a power-law from $h\nu_o$ to $\simeq 4\ kT_e$, with a slope ranging from 0 to -1 in energy as the parameter $y=4\tau^2\ kT_e/mc^2$ ranges from 1 to ln $(kT_e/h\nu_o)$ (Takahara 1980; Sunyaev and Titarchuk 1980). In fact, the emission may be so efficient as to drain energy from the electrons on a cooling time $t_r = \tau\ kT/\sigma_T cW \simeq 10^3\tau\ T_9R_{15}^2/L_{45}$ s shorter than the time for energy exchange from protons to electrons $t_{pe} \simeq 2\ (10^2/\tau)\ (kT_e/mc^2)^2\ R/v_{th} \simeq 3\ 10^5\ T_9^{3/2}\ R_{15}/\tau$, and matching the time for redistribution of electron energies $t_{ee} \simeq t_{pe}\ m/m_p$. The trends are apparent (Cavaliere 1981) from the ratios

$$\frac{t_r}{\Delta t} = 3\ 10^{-2}\ \frac{\tau}{1+\tau}\ \frac{T_9R_{15}}{L_{45}} \tag{2.3}$$

meaning fast cooling; and

$$\frac{t_{pe}}{t_r} = \frac{4\ 10^2}{\tau^2}\ \frac{T_9^{1/2}L_{45}}{R_{15}} \qquad \frac{t_{ee}}{t_r} = 5\ 10^{-4}\ \frac{t_{pe}}{t_r} \tag{2.4}$$

meaning slow collisional equipartition, particularly at low n such that $\tau \lesssim 1$ holds, and at high temperatures because $\sigma_{coll} \sim \sigma_T\ (mc^2/kT)^2$.

Two consequences follow. When the cooling is fast, the observed spectra will be time-averaged, resulting in flat power-laws much less dependent on y (Guilbert, Fabian and Ross 1982). If X-ray sources that vary

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fast at low energies like NGC 6814 should extend into the 10^2 keV range, collective processes must be at work to transfer power to electrons from the ions, which hold most of the energy in case of a direct accretion flow (like that contemplated in Eq. 2.2, cf. Guilbert, Fabian and Stepney 1982), or when the energy transport is effected by strong shocks.

Collisionless shocks at moderate Mach numbers across a magnetic field may instead heat preferentially the electrons. But for very compact sources, it may be problematic any way to maintain a truly thermal electron distribution against the enormous energy flux from the electrons by radiation, and to them by fields on a scale larger than microscopic. In the absence of magnetic field, $T_i >> T_e$ is very likely to hold (Rees, Begelman and Blandford 1981) with the electrons in a transrelativistic regime $kT_e \lesssim mc^2$.

Electron cooling may be enhanced still, when $kT_e > mc^2$ holds and $e^+ e^-$ pair production sets in. The $\gamma\gamma \rightarrow e^+ e^-$ process at $h\nu \simeq h\nu^1 \simeq 0.5$ MeV has probability

$$\tau_{\gamma\gamma} = \sigma_\gamma NR \simeq \frac{L_{45}}{R_{15}} \tag{2.5}$$

By itself, it degrades photons emitted at larger energies before they can leave the source, and imposes a spectral cut-off at $h\nu \simeq 0.5$ MeV (Cavallo and Rees 1978). In steady conditions, the newborn pairs share the available power, and the creation rate may easily be catastrophic; but then the energy per particle will decrease. This will not only thermostat the emitting plasma, but also will stabilize it at temperatures < mc^2 (Lightman 1981, Svensson 1981).

Alternatively to the thermal extreme, one may accept that electrons pick up power directly from large scale e.m. fields, and indeed run away in energy before sharing it with their fellows; the radiation rate itself is the ultimate limiting factor. Relativistic energies will result, and in the presence of a magnetic field one has much Synchrotron radiation, a prime candidate for the powerful, linearly polarized radiation of the OVV's and BL-lac Objects. In addition, much inverse Compton radiation will be emitted when the electrons collide with the Synchrotron photons. In a conventional non-thermal scenario, the relativistic electrons lose energy to photons in the amount $d\varepsilon/dt = -\gamma^2\sigma_T Wc$, that is, $\Delta\varepsilon/\varepsilon = \gamma^2 h\nu/\gamma mc^2$ per scattering event; the cumulative loss turns out to be formally large when $L_{45}/R_{15} \gtrsim 1$: $\tau_{e\gamma}\, \Delta\varepsilon/\varepsilon = \gamma L_{45}/R_{15} >> 1$.

From the photon view-point, the corresponding (rare but large) energy gain constitutes the (relativistic) inverse Compton upgrading; $\tau_{e\gamma}\Delta\varepsilon/\varepsilon >> 1$ means that the electron radiative life-time $t_r = mc^2/\gamma\sigma_T Wc$ is much shorter than $t_c = R/v$, the electron crossing time through the source at a velocity $v \lessdot c$:

$$t_r/t_c < t_r c/R = \varepsilon/\Delta\varepsilon\ \tau_{e\gamma} << 1 \tag{2.6}$$

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Catastrophic energy losses thus prevail again. The Synchroton photon production and the associated energy loss can be treated similarly: the process may be viewed as scattering off virtual photons with individual energy $h\nu_B = heB_\perp/mc$, energy density $W_B = B_\perp^2/8\pi > L/4\pi R^2 c$ when S radiation is dominant; so again $t_r/t_c << 1$ obtains, the condition for catastrophic losses.

The difficulty cannot be overcome with continous injection of fresh high energy electrons because the spent electrons would accumulate and easily violate the constraints $\tau < 1$ and Faraday rotation $\psi_B < 1$, erasing first the polarization characteristic of S emission, then affecting the rise time and the spectrum (Blandford and Rees 1978). The conclusion is that replenishment of the electron energy is the key problem with the compact sources, independently even of the radiation mechanism and of the thermal or non-thermal conditions of the emitting plasma.

The conventional S and IC processes are so efficient as to require reformulation. A limiting case has been examined by Pacini and Salvati 1978 and Cavaliere and Morrison 1980. The basic idea is to balance the radiative losses with multiple electron re-acceleration. Each electron is recycled, as it were, a number of times $\Xi >> 1$; the maximum value $\Xi \approx t_c/t_r$ ($\approx 10^2 - 10^3$) corresponds to "continous" reacceleration. To this condition the electrons will tend in the presence of energy gain rates $a(r) \neq 0$ throughout the region were the radiative losses $p = \varepsilon/t_r$ also occur. When the losses grow steeply with γ ($p = \sigma_T c W \gamma^2$ for S and IC, W being the relevant energy density) they soon must strike a balance with the gain such that $d\varepsilon/dt = a - p \to 0$:

$$\gamma^2 \simeq \frac{a}{\sigma_T c W} \tag{2.7}$$

Equivalently, this limiting condition may be written $t_r = t_a$. Note that the protons, with virtually no radiation losses, will be accelerated to $\varepsilon_p \approx \Xi \varepsilon_e$, well into the GeV range (Cavaliere and Morrison 1980).

The source remains very thin to Thomson scattering: $\tau \simeq L/cR^2B^2\gamma^2$. But the Faraday rotation at, or immediately outside, the IR source may be critical:

$$\psi_B \simeq \frac{L}{cR^2B^2} \; \frac{t_r eB}{\gamma mc} \; \left(\frac{\nu_B}{\nu}\right)^2 ; \tag{2.8}$$

and the perturbation of the ambient magnetic field in the source is measured by:

$$\frac{n\, mc^2}{B^2/4\pi} = \frac{L}{cR^2B^2} \; \frac{t_r}{t_c} \; \frac{c}{v} \tag{2.9}$$

Eqs. (2.8) and (2.9) require a source structure such that L/cR^2B^2 is $\lesssim 1$. High polarization requires also, independently of any specific ra-

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diation mechanism, strong order in the ambient $\underline{B}$ geometry. These requirements would be jointly satisfied in the scenario of energy extraction and transport from the primary source by large-scale Poynting flux associated with rotating objects (cf. Table 1); if the flux $\rightarrow cR^2 B^2$, the related efficiency is given by $L/cR^2B^2<1$.

The electrons radiate in the S and IC modes in the relative amounts $(L/cR^2B^2)^{-1}$. To obtain 10^{45} erg/s in the IR-O, and a comparable amount in X-rays, the source parameters needed are: $n \approx 10^{-5}cm^{-3}$, $B \simeq 10^3G$ and $\gamma \lesssim 10^2$; these scale as in the conventional Synchro-Compton theory.

However, the usual equality of the S and IC spectral slopes is broken here, because the emitted spectra are sustained by continuous energy pumping. The integrated spectra are determined by the inhomogeneity of the source: Eq. (2.7) determines local values of $\gamma(r)$ in terms of $W(r) \sim r^{-6}$ or $\sim r^{-2}$, once a(r) is given. The emitted frequencies increase outwards; more extended acceleration (a(r) more nearly constant across the source) will produce more power at higher frequencies, that is, flatter spectra and enhanced IC radiation. Typical S spectra will have slopes around -1; and the IC from -0.2 to -1.

In the relativistic components, pair production is an obvious possibility; when L_{45}/R_{15} > a few holds in X-rays, one expects either a $\gamma\gamma \rightarrow e^+ e^-$ cut off at about 0.5 MeV, or an earlier decline of the spectrum due to reduced efficiency $L/cR^2B^2 < 1$ plus a small range of electron energies lowering the successive scatterings. If continuous acceleration pervades the source (but this depends on geometry and texture), one conjectures a pair production instability similar to the thermal case, eventually lowering the energy per particle to an equilibrium value; in such sources, hard X-ray flares (at variance with IR-O outbursts) could not exceed the intrinsic limit $L/\Delta t \approx 10^{41}$ erg/s^2 by a large margin. To now, no X-ray source has been observed to violate grossly that limit, again at variance with IR-O outbursts; this may still be due to undersampling, throught it is confirmed by the first systematic search published (Tennant and Mushotzky 1982).

X-ray flares with $L/\Delta t >> 10^{41}$ erg/s^2 would constitute strong evidence in favour of sharp anisotropies in the X-ray emitting volume. On the other hand, optical sources highly variable with little or no X-ray counterpart imply on isotropic S-IC models a large magnetic field, so that $L_X/L_O \simeq L_O/cR^2B^2 << 1$; a case in point is OJ 287 (Wolstencroft, Gilmore and Williams 1982) with an output at 1.25 µm of at least 10^{44} erg/s observed to very on scales down to 50 s, which would imply magnetic fields up to 10^5 G within some 10^{12} cm.

The most attractive anisotropy is in the form of a thin jet moving with relativistic bulk velocity (Lorentz factor $\Gamma \simeq 5 - 10$) towards the observer, which connects so naturally with the radio observations. In the jet reference frame, the transverse size and the proper Δt are longer by a factor Γ, while the proper luminosity is smaller by

Γ^{-4}; thus the problems concerning high energy densities and catastrophic losses are apparently eased or eliminated. Jets are amply discussed by Blandford and Rees 1978 and Rees, Begelman and Blandford 1981.

Two comments apply. Technically, a single-jet source combines the simplest inhomogenity and anisotropy; Königl 1981 has produced a complete model of S and IC radiations from such a jet, which however when applied to 3C 273 cannot account for the radio data and also for the X-ray emission. Statistical difficulties in connecting radio and optical data are mentioned by Kellerman in this Conference.

On the other hand, also with jets problems concerning energy transport and replenishment creep in again. In fact, a proton-electron beam at $\Gamma \simeq 10$ requires that the accretion energy (a few 10^{-1} GeV per particle) be concentrated onto a smaller number of particles, either in a thermal form (then to be focussed in a subsonic -supersonic nozzle), or directly forming the jet; coherent accelerating e.m. fields are again first candidates. Much of the kinetic energy stockpiled in protons is to be retrieved down the jet in the form of electron energy for radiation by some kind of in situ electron acceleration, likely in internal shocks. Positron-electron jets, instead, pose lesser problems as for the bulk energy; however, these beams lose rapidly their energy, especially if their collimation is to be effected near the primary energy source: a relatively large γ (in the jet frame) is required to reduce the annihilation rates ($t_{ann} \simeq \gamma^2 R/c\tau$), but then the radiative cooling may be fast, and re-acceleration may be necessary once again. This may be evidence that Poynting flux must be an essential constituent of jets at their outstart. That it can be dominant, is discussed by Blandford 1976, Lovelace 1981, Blandford and Payne 1982.

Thus because of the size of the basic radiation cross-section σ_T, the strong radiation sources fed by a compact engine cannot be too compact after all; in addition, they must be inhomogeneous, sometimes anisotropic and perhaps still with internal reacceleration. Electromagnetic fields coherent on scales larger than microscopic are likely to play an essential role in feeding these sources.

3. THE EVOLUTION

The increase in number and/or luminosity with look-back time appears now to involve not only the classic extended radiosources but, with comparable strength, also the compact ones; it concerns even more strongly the very compact optical Quasars. The corresponding populations $\rho(P,z)$ in the plane monochromatic power-redshift behave similarly in style: they increase sharply from $z \cong 0.2$ to a few by a factor $10^{3\pm1}$, to level off or even bend over beyond $z \simeq 3$. The increase is clearly differential in power at radio frequencies, with a threshold $\simeq 2\,10^{33}$ erg/Hz; less definitely for the optical data, a threshold has been placed at $M_B \simeq -22$ (for recent discussions see: Cheney and Rowan-Robinson 1981, Peacock and Gull 1981, Schmidt and Green 1982, Setti and Woltjer 1982, van der Laan

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and Widhorst 1982, Wall and Benn 1982).

The compact radiosources may constitute just a subset of the extended ones, promoted to higher apparent fluxes by beaming effects (Scheuer and Readhead 1979). Gravitational lensing may affect the bright end of the counts of the very compact QSO s (Peacock 1982). However, a more fundamental driving agency is called for to account for an evolutionary pattern involving both the very extended and the very compact sources.

This may be traced back (Cavaliere et al 1982) to the common nature and working of the driving energy source - gravitational, in a compact configuration - above the specific features obviously required for so diverse radiative out-puts like those of compact, radio-quiet QSO s and of double-lobed, optically-quiet Radiogalaxies. An ispection of Table 1, that summarizes the results of several leading models for the release of gravitational energy, uncovers a persisting trait: the secular change of the primary power output L is of the form

$$L = f(L) \qquad (3.1)$$

with $d|f|/dL > 0$ if not $d(|f|/L)/dL > 0$; often a simple power-law expression $f(L)=\pm A\ L^{1+p}$ $(p>0)$ holds.

These results may be understood heuristically, considering that the primary machine comprises two parts: the active potential well where the energy is actually released, and the mass supply providing gaseous fuel to the former.

As to the latter, a local supply (adding to any slow accretion of diffuse interstellar matter) is required to fuel $L \gtrsim 10^{45}$ erg/s during $\approx 10^{8}$ yr; it must be held up against self-gravitation by a dynamical constraint - angular momentum or random orbital energy - maintaining a virial equilibrium. As the constraint is gradually dissipated or removed, more mass collapses to the inner engine and more energy is liberated.

But by the same token, the supply contracts and/or the central mass grows larger; in either case, the processes governing the constraint disposal speed up, to the effect that L increases secularly. These local processes, once started, will proceed at a rate independent of external, cosmic time-scales, but depending instead on the prime variable of state for the system, L itself; thus $\dot{L}=f(L)$ will obtain.

If most of the mass made available is actually converted at constant efficiency, then the non-linear scale-free trend of gravity is expected to prevail, and $f(L)=AL^{1+p}$ $(p>0)$ to hold. But the detailed working of the inner engine (efficiency, imperfect coupling to the supply, constraints to black hole growth) may intervene here and imprint a scale, for example the scale L/L_E at high regimes. In these cases one expects the more general form $\dot{L}=f(L)$, but still retaining the non-linear trend $d(f/L)/dL > 0$. The key feature is that $\dot{L}$ increases with an effective time-scale $t_s=L/\dot{L}$ that shortens with L increasing.

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On the other hand, when about half of the local supply is exhausted, the rate of constraint dissipation must saturate or decline, either because the central mass cannot grow substantially larger, or because the supply dynamically dominated by the central mass is no longer in virial equilibrium. Thus a dimming phase ($\dot{L}<0$) takes over; its course may be represented as a dynamical time-reversal of a brightening phase (with different initial conditions and a different value of the index p). Again the property $d(|f|/L)/dL \geq 0$ holds, now meaning that $t_s = L/|\dot{L}|$ grows longer as L decreases.

The dynamics of N(L,t), the epoch-dependent luminosity function, is linked to the time-scales $L/\dot{L}$ for the individual objects by a continuity equation in the plane L-t (Cavaliere, Morrison and Wood 1971)

$$\frac{\partial N}{\partial t} + \frac{\partial}{\partial L}(\dot{L}N) = S \qquad (3.2)$$

This contains a birth function S(L,t) describing the appearence of new objects in a region δt, δL; but its direct influence is little outside the original region, when $d(|f|/L)/dL > 0$ holds. For in the remaining part of the L-t plane the distribution of the population over L at successive times is redirected by the geometry of the lines $dL/dt=f(L)$, the characteristics of Eq. (3.2). The key feature is the divergence or the convergence in a log L-t plot; this is directly related to the dependence $t_s=t_s(L)$ seen above, and is described quantitatively by

$$D = p \text{ sign } \dot{L} \quad \text{with} \quad p = -d \ln t_s / d \ln L \qquad (3.3)$$

It is easy to see that density evolution obtains, $N(L,t) \rightarrow g(t)\,N(L)$, in a brightening phase with divergence (i.e., $\dot{L}>0$, $D>0$), in the region of high L defined by : $L > (1/pAt)^{1/p}$ when $\dot{L} \sim L^{1+p}$, or more generally, for L such that $t_s(L) < t$ holds. Uniform or differential luminosity evolutions obtain for a dimming phase ($\dot{L}<0$) with neutral behaviour (D=0), or with convergence (D<0): N(L,t) displaces downward in luminosity, bodily in the formes case, or in the latter with a steepening at the bright edge $L > (1/pAt)^{1/p}$, i.e. for $t_s(L) < t$.

Fig. 1 shows how the simplest scheme works: a pure dimming phase with convergence fed by a birth function extending over substantial times $\delta t \approx H_o^{-1}$. Shorter S(t) would result in a very narrow local N(L); neutral behaviour would imply an ad hoc initial N(L).

The demands on S(L) (to produce objects "ready made" out to the maximum luminosities), and the heavy dependence of the results on the uncertainties of the formation processe underlying S(L,t), are relieved considering also the phase of brightening that ought to precede the dimming, the two being matched by number conservation. So the objects may be produced all at $L \ll L_{Max}$, and the details of S(L,t) (but also much of the associated information) are largely erased when the brightening extends for a range > 10 $L_{initial}$. Fig. 2 shows one example

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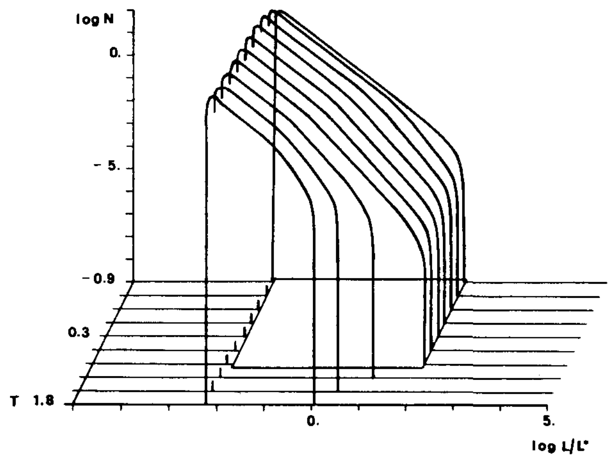

Fig. 1. A population in the L-t plane of nuclei that undergo pure dimming with convergence from a birth function of substantial duration: $\dot{L} \sim -L^{1.5}$, $S(L,t) \sim L^{-1.6}$ • const within the blank rectangle $\delta t\ \delta L$ centered at L^*, t^* and 0 otherwise. $T=2(t-t^*)/\delta t$. Time-scales: $t_s(L^*) = \delta t/2$.

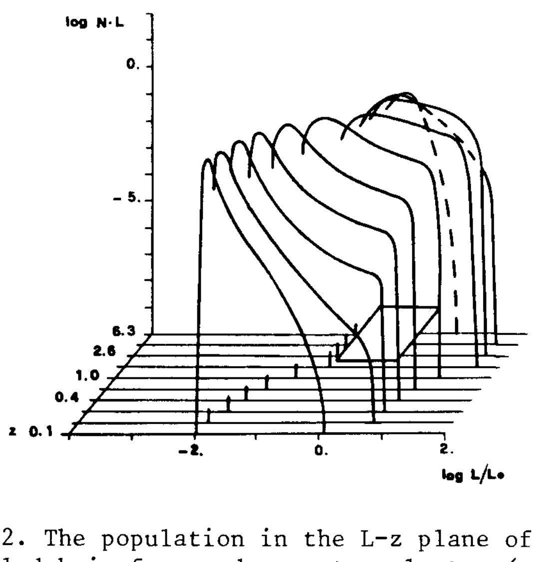

Fig. 2. The population in the L-z plane of black holes accreting collisional debris from a dense star cluster (see text). Brightening with divergence, $\dot{L} \sim -L^{2.2}$, by a factor 20; dimming with convergence, $\dot{L} \sim -L^{1.5}$. The birth function is active in the blank rectangle $\delta t\ \delta L$ centered at t^*, L^*. Time-scales set equal: $\delta t/2=t_{s+}(L^*)=t_{s-}(L^*)$. World model: Einstein-de Sitter; t^* corresponds to $z^* = 3.5$.

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of a full history, early brightening plus late dimming, based on the model discussed by McMillan, Lightman and Cohen 1981: a dense star cluster self-destroying by star collisions, and feeding a black hole at a sub-Eddington rate. We find (see Table 1) for such a model $\dot{L} = \pm AL^{1+p}$ with $p_+ \rightarrow 1.2$ and $p_- = 0.5$; Fig. 2 represents the total population $N = N_+ + N_-$. The model is specific for the optical QSOs, but is also representative for the overall behaviour of the whole class of gravitational sources with a scale-free output.

The following comments apply. The general aspect of the populations of Active Nuclei (cf. Schmidt and Green 1982, Peacock and Gull 1981) is reproduced fairly well without special parameter optimization. The overall course at the (more directly observable) intermediate and late epochs is a luminosity evolution, driven for a while by the rapidly decreasing N_+ that plays the role of an effective birth function at high values of L. In particular, the luminosity functions at early times are displaced to, and peaking at, higher values at L compared with the local one; the slopes of the luminosity functions are $(1+p) \simeq 2$ and steeper; the integral counts corresponding to Fig. 2 have a slope $\simeq 2.2$ at high fluxes, and then converge rapidly. But before attempting detailed fittings of population observables in the optical band, a number of additions shall be made, including specific k-corrections and distributions of spectral indexes, and a realistic statistical distribution of constitutive parameters; these effects will broaded and smoothen the luminosity functions at the bright end.

As for Radiogalaxies, we will try in the same framework the model proposed by Blandford and Znajek 1977 and further discussed by Rees et al 1982. A Kerr hole rapidly spun up to a condition $J \lesssim J_{Max}$ by mass inflow from an accretion disc fueled in turn by gas production in a large star cluster, emits as focussed Poynting flux the stored rotational energy at a rate $L_{45} = 0.2\ B_4^2 M_8^2 (J/J_{Max})^2$ when the accretion flow subsides to the value sufficient to sustain a magnetic field of order 10^4 G at the horizon. From this extreme compact model one expects a prolonged dimming phase with nearly neutral behaviour, modified by the electrodynamical efficiency factor $L/\dot{J}$. Besides local supplies, also the diffuse interstellar medium may contribute to accretion, especially when enhanced by cooling inflows in those galaxies that sit at the center of dense clusters (Cowie and Binney 1977); these contributions, however, add to the low luminosity sections of the populations of Active Nuclei, and the latter in particular sets in only at fairly late epochs.

Thus, beyond the common traits of the evolutionary behaviours that appear to be consistent with the basic properties of gravitational, compact energy sources, engaging prospects arise of indentifying from finer features of the source populations the signature of the specific kind and operation of the primary engine.

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TABLE 1

The secular change of the primary power output for models relying on a compact mass supply

Mode of energy release	Constraint	$\dot{L} = f(L)$	$t_s = L/\dot{L}$ 10^9 yr
Class 1) Accretion onto a black hole $L = \eta\dot{M}c^2$			
b.h. accreting from a star cluster, sub-Eddington regime (brightening)	v^2, dissipated via tidal disruption and cooling	AL^{1+p} $p=0.25$	$\frac{R_{ipc}}{N_{i8}} \frac{N_{i8}^{1/2}}{m_{i8}^{1/2}}$
b.h. accreting from a dense star cluster, super-Eddington regime (brightening)	v^2, dissipated via disruptive collisions and cooling	$a(\Lambda-L)^{-k}$ $1\leq k\leq 4.5$	$\frac{L_i}{\Lambda}\left(\frac{R_{ipc}}{N_{i8}}\right)^{7/2} N_{i8}^2$
sub-Eddington regime (brightening)	" "	AL^{1+p} $1.22\leq p<2$	$\left(\frac{R_{ipc}}{N_{i8}}\right)^{7/2} N_{i8}^2$
(dimming)	star dynamics determined by the b.h.	$-AL^{1+p}$ $p=0.5$	$\left(\frac{R_{ipc}}{N_{i8}}\right)^{7/4} \frac{R_{ipc}^{7/4} N_{i8}^{1/2}}{m_8^{1/4}}$
Class 2) Rotating magnetized configurations $L = cR_c^2 B_p^2$			
spinar, spinned up by contraction (brightening)	J, removed via $B_\perp$ electro-dynamically	AL^{1+p} $0.75\leq p<1$	$0.1\left(\frac{GM^2}{J^2}\right)^{1/4} \frac{M_8}{L_{42}^{3/4}}$
" " (dimming)	$\underline{B}$ aligns to $\underline{J}$	$-AL^{1+p}$ $p=0$	$\frac{t_{s+}}{\cos^2\chi_i}$
Class 3) Accretion and rotation combined $L = cR_c^2 B_p^2$			
magnetized accretion disk fed by a star cluster	v^2, dissipated via tidal disruption; and J, electrodynamically	AL^{1+p} $p\simeq 0.1$	$\frac{0.7}{m_8^{1/3}}$

References. Class 1: McMillan, Lightman and Cohen 1981. Class 2: Cavaliere, Morrison and Wood 1971; Ozernoy and Usov 1973. Class 3: Blandford 1976. N.B. The table is taken from Cavaliere, Giallongo, Messina and Vagnetti, in preparation. To class 3, it may be added the energy release by a rotating hole in a magnetic field, discussed by Blandford and Znajek 1977 and by Rees et al. 1982.

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