MASSIVE OSCILLATORS AS COSMIC ENERGY SOURCES

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Abstract. The usual picture in which a massive object undergoes a gravitational collapse to become a black hole and ultimately end up in space-time singularity, is modified with the introduction of a negative energy force of repulsion effective only at a short range. It is shown that the object executes oscullations between states of high and low densities.

From the view point of high energy astrophysics, such a massive oscillator combines some of the attractive features of black holes and white holes. It is suggested that the energy production and spectral features of quasars, BL-Lacs and the active galactic nuclei might be accounted for by postulating the existence of massive oscillators.

1. Introduction

A few years ago we had advocated white holes as energy machines which power cosmic sources of radiation (Narlikar and Apparao, 1975, hereafter referred to as Paper I). A typical white hole was considered as a mini-explosion from a local space-time singularity, its internal geometry resembling that of a Friedmann universe. It was shown that radiation reaching a remote observer from the surface of the white hole is strongly blue shifted in the early stages of expansion, and this result was used to argue that a white hole can act as a source of high energy particles and quanta. Several instances from high energy astrophysics were discussed in Paper I to illustrate the role of white holes as energy sources.

In Paper I we had mentioned in passing that the singular origin of white holes can be averted if we introduce negative energy and pressures into the equations of general relativity. Such a device can lead to a bounce of a massive object undergoing gravitational collapse. After the bounce, the object will re-expand and its dynamics then resembles that of a white hole with a singular origin. By having successive phases of expansion and collapse, a massive oscillator can generate energy over an extended period instead of in a single burst. In this paper we discuss the properties of such bouncing objects.

One of the difficulties with the classical white hole has been that its origin as a 'delayed mini-bang' or a 'lagging core' is somewhat *ad hoc*. A common criticism is that we do not know what triggers off a white hole explosion. Such a criticism, although surprisingly rarely voiced against the big bang origin of the universe, does bring out the incompleteness of the white-hole scenario. The present picture is free from this criticism, since here the object begins its existence with gravitational collapse in much the same way that leads to the formation of a black hole in the classical black-hole astrophysics.

Another point, usually slurred over in discussions of black-hole astrophysics is that, in the strict mathematical sense, the collapsing object takes infinite time to approach

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the Schwarzschild radius and thus becomes a black hole in the reference frame of a remote observer. Hence, none of the black holes discussed by the astrophysicists are the black holes of relativists. In the present picture, the outer surface of the massive oscillator always stays outside its Schwarzschild radius and, thus, the event horizon is never formed.

There is another respect in which the bouncing, oscillating solution differs from that of the classical white hole of singular origin. Eardley (1974) had given certain geometrical arguments to conclude that the classical white hole has a very short life time. The space-time geometry of our nonsingular oscillator is markedly different from that of the classical white hole. Eardley's arguments leading to a short lifetime do not apply to the former.

We begin with a geometrical description of the bounce of a collapsing object.

2. The Bounce of a Collapsing Object

We will follow the notation used in Paper I and consider the oscillator as a spherical object of uniform density and zero pressure. To produce oscillations, we now introduce a bounce producing term which becomes important when the object is highly compact but whose effect rapidly dwindles as the expansion proceeds. Conventional physics had no such term and we have to introduce a new one. The introduction of such a term would violate the singularity theorems in general relativity (cf. Hawking and Ellis, 1973) and would open to us one of two possible choices. One choice is to place absolute faith in the completeness of physics as we know it now and in the correctness of general relativity. This choice would then rule out the bouncing of collapsing objects and imply that a space-time singularity is an essential feature of the universe. The second alternative is to argue that space-time singularities are undesirable features and their appearance in our calculations indicate incompleteness of relativistic physics. This choice opens up the possibility of new inputs. By writing this paper we are committing ourselves to the latter choice.

New inputs can always be criticized as being *ad hoc*. We accept this criticism although it can be partially countered by two points. First, we will use a formalism which has already been used successfully in producing nonsingular cosmological models. Second, the general quantitative behaviour of the bouncing white hole will not depend critically on what bouncing formalism is used.

The new term we wish to add is the negative energy scalar field C of zero rest mass and zero charge. Its properties have been discussed elsewhere in detail (Hoyle and Narlikar, 1963, 1964) and we will not repeat them here. Its effect on the original white hole solution is as follows. The Einstein field equations with the C-field are

$$R_{ik} - \tfrac{1}{2}g_{ik}R = -\frac{8\pi G}{c^4}[T_{ik} - f\{C_iC_k - \tfrac{1}{2}g_{ik}C^1C_1\}], \tag{1}$$

where T_{ik} is the energy tensor for dust and $C_i \equiv \partial C/\partial x^i$.

As in Paper I we take the interior solution of these equations as the Robertson–

Walker space-time with the line element

$$ds^2 = c^2\,dt^2 - S^2(t)\left[\frac{dr^2}{1-\alpha r^2} + r^2(d\theta^2 + \sin^2\theta\,d\phi^2)\right], \tag{2}$$

where r, θ, ϕ are the comoving spherical polar coordinates of a typical dust particle in the object, with the interior given by $r \leqslant r_b$. $S(t)$ is the scale factor which increases with time during the expansion phase of the oscillator and decreases with time during the contracting stage. For a bouncing solution S must oscillate between finite limits. The parameter α is related to the dust density ρ_0 at $S = 1$ by the relation

$$\alpha = \frac{8\pi G\rho_0}{3c^2}. \tag{2}$$

The C-field inside the white hole is a function of t only, being given by

$$\dot{C} \equiv \frac{dC}{dt} = \frac{A}{S^3}, \quad A = \text{constant}. \tag{4}$$

It is convenient to define another parameter β which relates A to the coupling constant f of the C-field to matter and gravity, by

$$\beta = \frac{4\pi G}{3c^2} fA^2. \tag{5}$$

The physical interpretation of (4) is as follows. Just as a frozen-in magnetic field increases in strength as the object in which it is embedded contracts, so the *C-field* strength increases in the above case as S decreases. Thus, when a massive object begins to contract it picks up and squeezes the ambient cosmic C-field. Although to start with the C-field has negligible strength and no astrophysical effect, it gains in strength according to (4) and ultimately affects the dynamics of collapse.

Instead of Equation (3) of Paper I, $S(t)$ now satisfies the differential equation

$$\dot{S}^2 = \alpha c^2\left(\frac{1-S}{S}\right) - \frac{\beta c^2}{S^4}. \tag{6}$$

The C-field term has all the desirable properties stated earlier. It causes a bounce at a small value of $S > 0$, while its effect diminishes for

$$S \gg \left(\frac{\beta}{\alpha}\right)^{1/3}. \tag{7}$$

Thus, as the maximum value of S is still very close to 1, we require that

$$\beta \ll \alpha. \tag{8}$$

There is another inequality imposed on β by the requirement that the outer boundary of the object *always* remains outside its Schwarzschild radius. This requirement cannot be satisfied for the canonical white hole of Paper I. The main difference between the two cases comes from the fact that in the present case, the gravitational mass of the object changes with time, decreasing as the strength of the C-field increases.

To calculate the gravitational mass we transform (2) to Schwarzschild-type coordinates given by

$$R = rS(t), \qquad T = F(\xi),$$
$$\xi = \int_{\bar{t}}^{t} \frac{\mathrm{d}\tau}{S(\tau)\dot{S}(\tau)} - \int_{r}^{r_b} \frac{x\,\mathrm{d}x}{1 - kx^2}, \tag{9}$$

where $\bar{t}$ is an arbitrary constant and F an arbitrary function of ξ. The resulting line element is given by

$$\mathrm{d}s^2 = c^2 e^{\nu}\,\mathrm{d}T^2 - e^{\lambda}\,\mathrm{d}R^2 - R^2(\mathrm{d}\theta^2 + \sin^2\theta\,\mathrm{d}\phi^2). \tag{10}$$

The form of e^{ν} depends on F; but $e^{-\lambda}$ is given by

$$e^{-\lambda} = 1 - \frac{2GM(R)}{c^2 R} \equiv 1 - \frac{r^2}{c^2}(\dot{S}^2 + \alpha c^2). \tag{11}$$

The gravitational mass of the object is obtained by evaluating $M(R)$ on the boundary $R_b = r_b S(t)$ of the object. From (6) we get

$$M(R_b) = \frac{c^2}{2G}\left(\alpha - \frac{\beta}{S^3}\right) r_b^3. \tag{12}$$

The mass is reduced by the negative energy of the C-field.

The Schwarzschild radius of the object therefore decreases as it contracts. To ensure that R_b always lies outside the Schwarzschild radius we demand that the maximum value of $2GM(R_b)/R_b c^2$ does not exceed 1. A simple calculation then gives

$$\beta \gg \frac{27}{256}\alpha^4 r_b^6. \tag{13}$$

Such a lower limit explains why the canonical white hole of Paper I (with $\beta = 0$) does not always remain outside its Schwarzschild radius. The above limit also implies that the universe has a background energy density of the C-field. Such a background is indeed present in the C-field cosmologies of Hoyle and Narlikar (1966a).

The above condition is important because it ensures that the object is always visible to an external observer. As the object expands and contracts, the external observer (at rest in the Schwarzschild frame) sees it pass through alternate phases of blueshift and redshift. In the following section we will estimate the magnitudes of these spectral shifts and their duration as seen by such an observer.

The conditions (8) and (13) together require that

$$\sin^2\eta_s \equiv \alpha r_b^2 \ll \tfrac{1}{3} \times 4^{4/3} \approx 2.1. \tag{14}$$

The parameter η_s was introduced in Paper I to indicate how collapsed the object is and it was related to the duration of blueshifts and redshifts of the canonical white hole. We will use this parameter in our subsequent work.

3. Spectral Shifts

The line element (10) can be matched to one outside the object (for $r \geqslant r_b$). We can in principle continue λ and ν as functions of R, T in this region of space-time, as was done for the canonical white hole of Paper I. The field Equations (1) will then have to be solved for $T_{ik} = 0$ (empty space) and $C = C(R, T)$ with $\partial C/\partial R$, $\partial C/\partial T$ tending to zero at $R \to \infty$. On the boundary $r = r_b$, C must be continuous at all t (or T). Also, at $r = r_b$

$$e^{-\lambda} = 1 - \frac{r_b^2}{c^2}(\alpha c^2 + \dot{S}^2), \tag{15}$$

$$\left(\frac{\partial T}{\partial t}\right)^2 e^{\nu} = (1 - \alpha r_b^2)e^{\lambda}. \tag{16}$$

To compute the spectral shifts we need to know ν and λ outside the object. We have not been able to solve the field equations together with the above boundary conditions. While it can be demonstrated that a solution exists, its computation involves the difficult (but probably not intractable) task of solving non-linear coupled partial differential equations (cf. Hoyle and Narlikar, 1964). We have therefore followed an approximation procedure. This involves using the broad features of the empty Schwarzschild solution in the expectation that as we move away from the object the effect of the C-field terms will diminish rapidly. Thus, we will assume that

$$\nu + \lambda \approx 0, \qquad \frac{\partial \lambda}{\partial T} \approx 0 \tag{17}$$

outside the object.

Let light leave radially outwards from a surface point $R = R_b$, $T = T_1(t = t_1)$ to reach an observer at $R = R_2$ at $T = T_2$. Then a simple calculation along the lines given in Paper I gives the redshift of this radiation as z, such that

$$1 + z \cong (\sqrt{1 - \alpha r_b^2} + r_b\dot{S}/c)^{-1}_{t=t_1}; \tag{18}$$

the right-hand side being evaluated at the point of emission. Note that during expansion, large values of $\dot{S}$ can make $z < 0$, giving blueshifts. However, provided (13) is satisfied we can avoid $z \to \infty$ for $\dot{S} < 0$, i.e., the actual attainment of the black-hole state. The collapsing object will bounce at a finite but large z and later the external observer will see z change signs.

In (18) S is a function of t_1; and since

$$1 + z = \frac{dT_2}{dt_1}, \tag{19}$$

we can determine z as a function of T_2.

Figure 1 illustrates (z, T_2) curves for a range of parameters. It is convenient to express

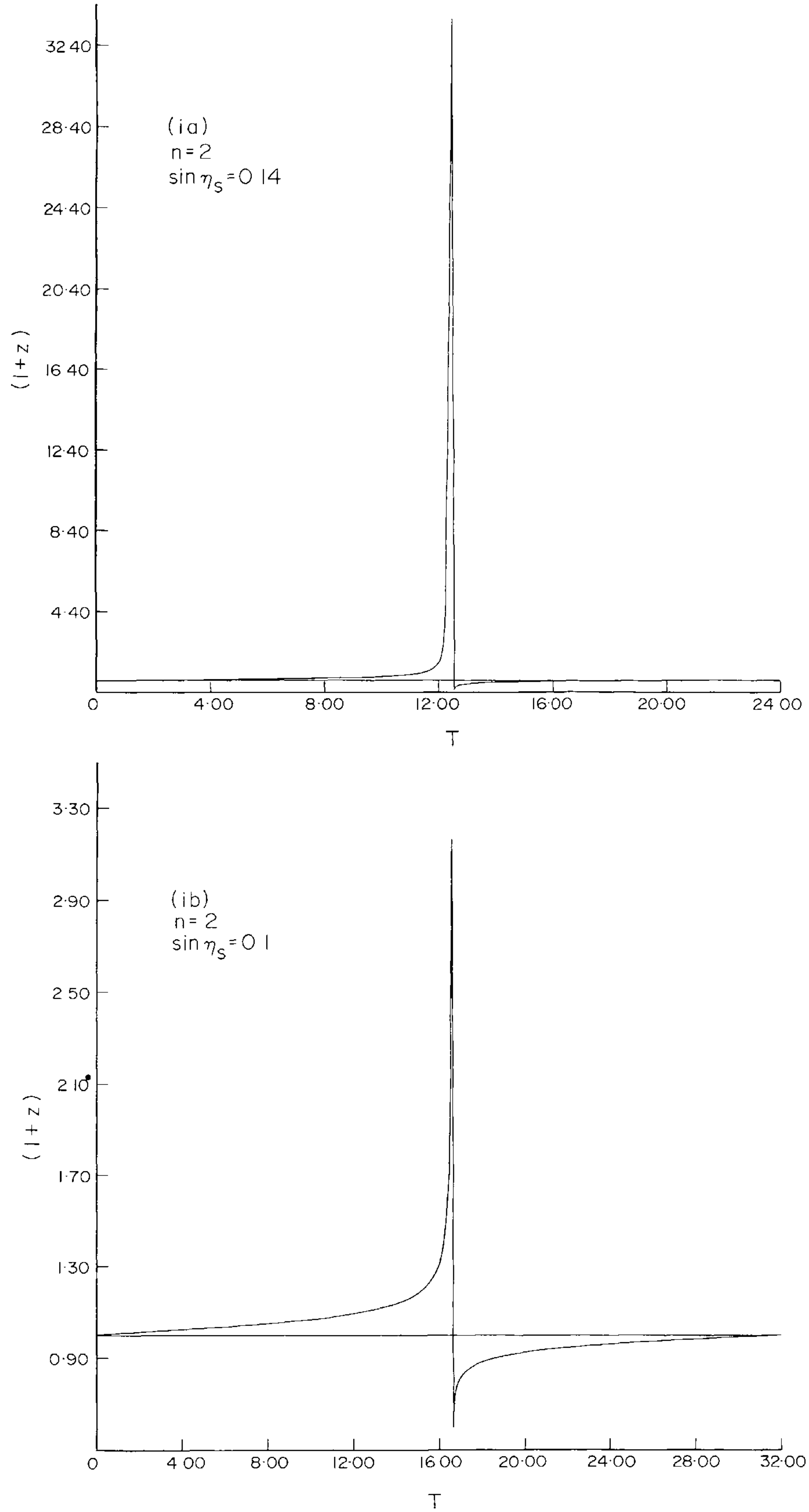


Fig. 1. The variation of line shift with time as seen by an external observer during one cycle of oscillation. (i a, b), $n = 2$ as sin η_s changes from 0.14 to 0.1; (ii a, b, c, d), $n = 1$ as sin η_s changes for 0.46 to 0.1.

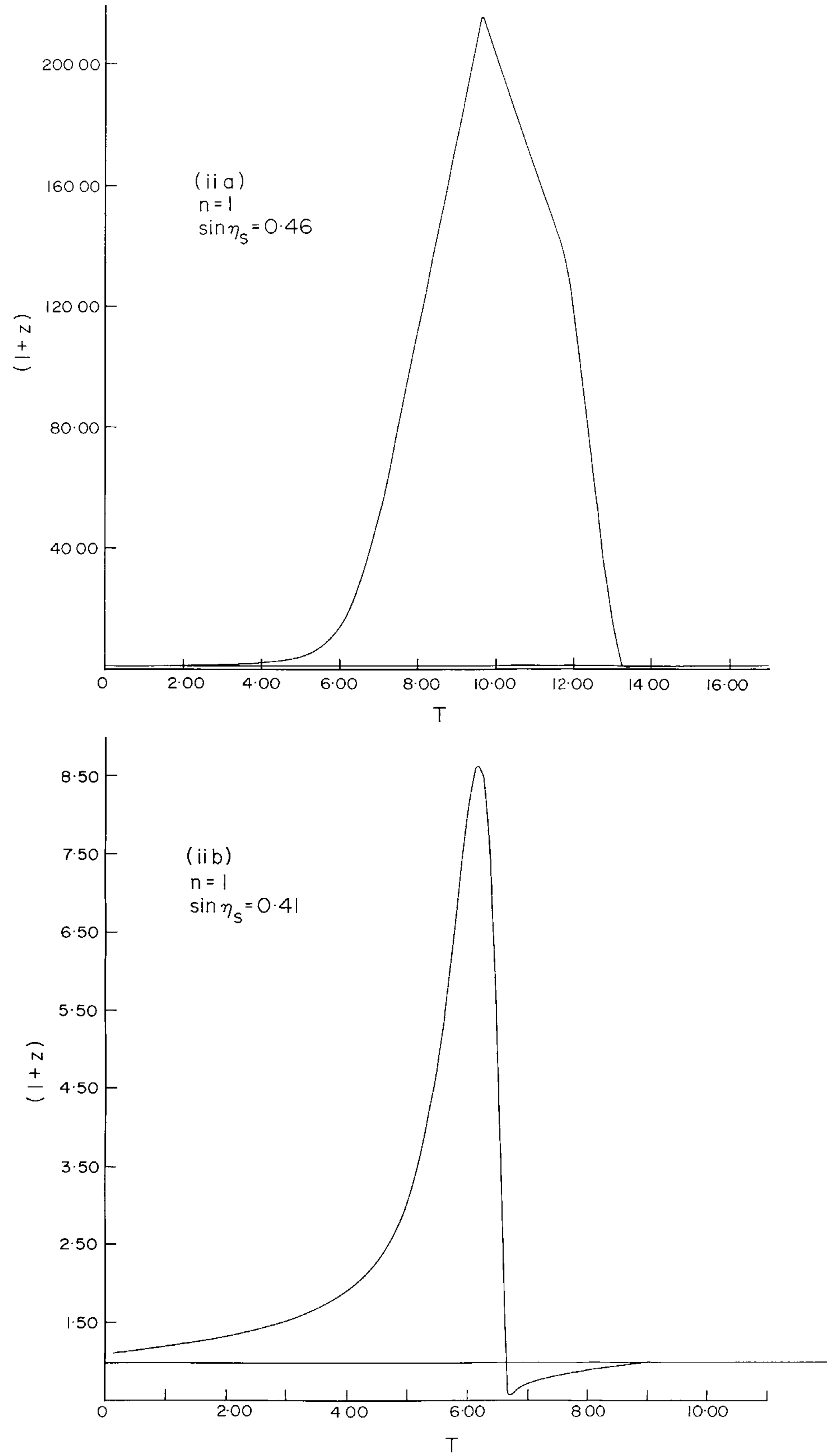
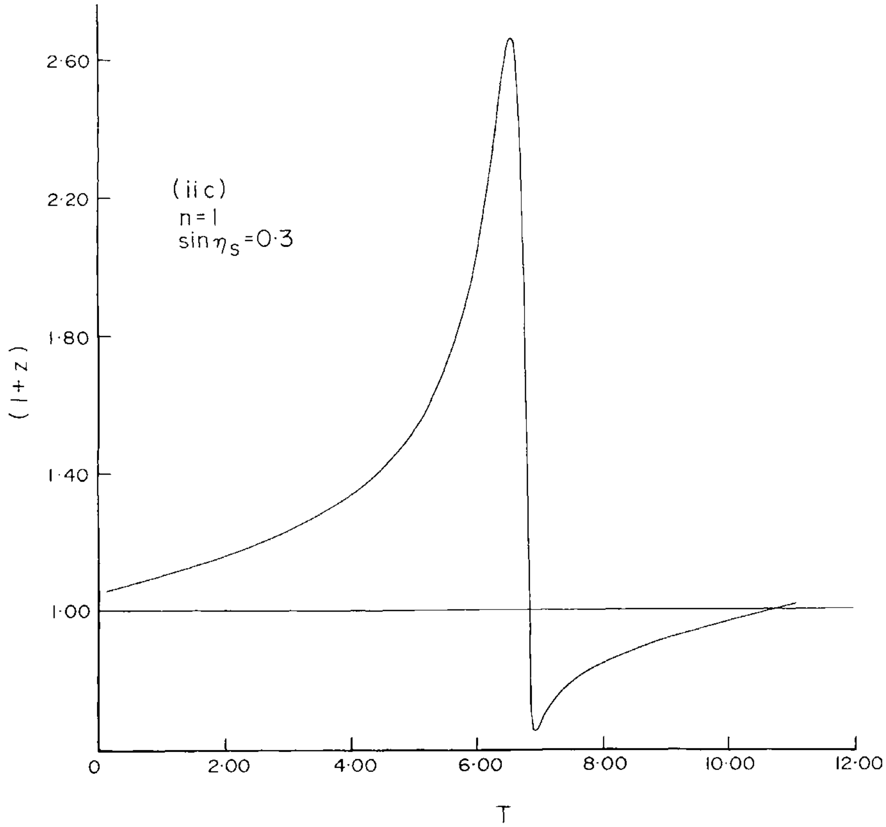


Fig. 1 (*continued*).



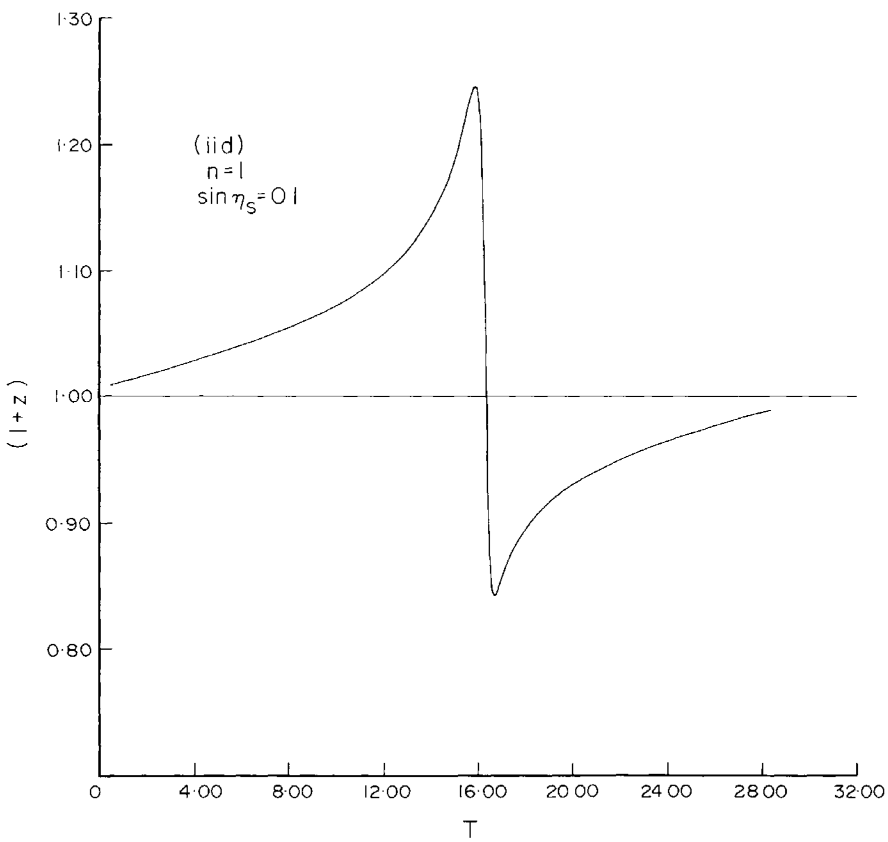


Fig. 1 (*continued*).

(18) and (19) in the form

$$dz = (1 + z)\,\frac{\sin^2 \eta_s}{2S^2}\left\{1 - \frac{4 \times 10^{-3n}}{S^3}\right\} dT_2, \tag{20}$$

where $\sin \eta_s$ is given by (14) and

$$\beta = 10^{-3n}\alpha. \tag{21}$$

Thus large values of n imply that the strength of the C-field is relatively weak. [However, (13) sets an upper limit on n.] A weak C-field causes the bounce to occur at a smaller value of S and the blueshift can be higher and of short duration. The numerical results of integrating (20) are tabulated in Table I and they illustrate the general pattern of behaviour.

TABLE I

Case No.	n	$\sin \eta_s$	Maximum redshift $(1+z)_{\max}$	Maximum blueshift $(1+z)_{\min}$	$\frac{T_R}{T}$	$\frac{T_B}{T}$
1	1	0.46	216.24	0.56	0.82	0.18
2	1	0.41	6.66	0.59	0.76	0.24
3	1	0.36	4.22	0.61	0.72	0.28
4	1	0.30	2.67	0.65	0.68	0.32
5	1	0.22	1.84	0.71	0.62	0.38
6	1	0.10	1.25	0.84	0.55	0.45
7	2	0.14	33.97	0.51	0.59	0.41
8	2	0.10	3.18	0.60	0.56	0.44

In Table I, T_R/T denotes the fraction of the total oscillation period T for which the object is seen redshifted and T_B/T denotes the corresponding fraction for blueshift. The maximum redshifts and blueshifts are also shown. It is obvious from the numerical results of Table I and from Figure 1 that there is a considerable diversity in the overall appearance of these objects, depending on what parameters are chosen. In the following sections we will discuss astrophysical applications of these models.

4. Astrophysical Applications

4.1 Quasars

Ever since their discovery in 1963, the quasars have posed an energy problem. Whether they are at cosmological distances or are local, we have to explain the existence of vast reservoirs of energy in small regions: in the former case, emission of galactic order must come from a region of linear size $\lesssim 10^{-1}$ pc (see Burbidge and Burbidge, 1967) while in the latter case we have to postulate star-like objects containing $\sim 10^6\,M_\odot$.

In the present model the energy source lies in the ram pressure of the surface of the object as it expands outwards. The surface moves with considerably high velocity and if it is surrounded by gas, part of its kinetic energy will be communicated to the gas. We expect that the gas will heat up as it is swept outwards and will eventually radiate thermally. An estimate of the energy available for radiation in this way can be made as follows.

Let ρ be the density of the ambient gas near the expanding surface. The velocity of expansion of the surface may be estimated as

$$v = \frac{\mathrm{d}R_b}{\mathrm{d}t} = r_b\dot{S}(t), \tag{22}$$

and we expect the gas to acquire this velocity. Then as the object expands between the epochs (t_1, t_2), it dumps an energy equal to $\mathscr{E}(t_1, t_2)$ into the surrounding gas, where approximately

$$\mathscr{E}(t_1, t_2) = \int_{t_1}^{t_2} \rho v^2(t) 4\pi R_b^2(t)\, v(t)\, \mathrm{d}t, \tag{23}$$

where we have used the non-relativistic form of energy transferred for ease of integration which should then give a lower limit.

Let $S(t_1) = S_1$, $S(t_2) = S_2$. Then changing the variable to S and using (6) we get

$$\begin{aligned}\mathscr{E}(t_1, t_2) &= 4\pi r_b^5 \int_{S_1}^{S_2} S^2\left[\alpha c^2\left(\frac{1-S}{S}\right) - \frac{\beta c^2}{S^4}\right]\mathrm{d}S \\ &= 4\pi\rho r_b^5 \alpha c^2 \int_{S_1}^{S_2} [S - S^2 - 10^{-n}S^{-2}]\,\mathrm{d}S.\end{aligned} \tag{24}$$

If we set S_1 = the minimum (i.e., bounce) radius and S_2 = the maximum radius (≈ 1) we get

$$\mathscr{E} \approx \frac{2\pi\rho}{3} \times r_b^5 \alpha c^2. \tag{25}$$

(It can be verified that the contribution from the lower end of the integral for $\mathscr{E}$ is negligible.)

We now relate (25) to the mass of the bouncing white hole. For $S \approx 1$, we denote $M(R_b)$ by M and write with the aid of (12)

$$\mathscr{E} \approx \frac{2\pi\rho}{3} GMr_b^2. \tag{26}$$

It is convenient to replace r_b by the time of expansion from S_1 to S_2. Since the C-field is effective for only a short duration, we may approximate the required value by t_0, the total time of expansion for the canonical white hole of Paper I, to be given by

$$t_0 = \frac{\pi}{2c\sqrt{\alpha}}. \tag{27}$$

Using (12) and (27) to eliminate r_b, we get

$$\mathscr{E} \cong \frac{8}{3\pi^{1/3}} \rho (GM)^{5/3} t_0^{4/3}. \tag{28}$$

To re-express $\mathscr{E}$ in astrophysically-relevant units we take $\rho \sim 10^{-17}$ g cm^{-3} as a characteristic gas density. This is the density in the line emitting region of quasars (see Burbidge and Burbidge, 1967). This value of ρ corresponds to a particle density $n \sim 10^{+7}$ cm^{-3}. From spectroscopic considerations of quasar emission lines (which should arise from the hot gas), this value of n is not too high. Although, there is no *prima facie* case for choosing any particular value for t_0, we will take $t_0 \sim 10^{12}$ s for a reason to be discussed shortly. Finally, for M we take a characteristic value of $\sim 10^4 M_\odot$ as lying intermediate between a purely stellar-type quasar mass considered by Terrell (1964) in his galactic ejection hypothesis and a supermassive object needed for the cosmological quasar. We will discuss this choice of M further. We then get

$$\mathscr{E} \approx 2 \times 10^{49} \left(\frac{M}{10^4 M_\odot}\right)^{5/3} \left(\frac{\rho}{10^{-17}\,\text{g cm}^{-3}}\right) \left(\frac{t_0}{10^{12}\,\text{s}}\right)^{4/3} \text{erg}. \tag{29}$$

The average luminosity corresponding to (29) is

$$L \approx 2 \times 10^{37} \left(\frac{M}{10^4 M_\odot}\right)^{5/3} \left(\frac{\rho}{10^{17}\,\text{gcm}^{-3}}\right) \left(\frac{t_0}{10^{12}\text{s}}\right)^{1/3} f' \,\text{erg s}^{-1}. \tag{30}$$

The additional factor f' in (30) is less than unity and represents the fraction of available energy which is actually radiated by the gas.

This expression for L is strongly dependent on M and weakly dependent on t_0. The value for L for $M = 10^4 M_\odot$, $\rho = 10^{-17}$ g cm^{-3} and $t_0 = 10^{12}$ s is in the region of luminosities of pulsars and galactic X-ray sources, suggesting therefore that if quasars are to be discribed by such objects they cannot be very far away from our Galaxy. For example, for an object of this luminosity to produce the observed flux from 3C-273, its distance from us must be of the order of ~ 20 kpc. Thus, we are required to assume that these quasars were ejected from our own Galaxy. However, unlike Terrell's (1964) model, this model does not demand relativistic speeds of ejection since the redshifts of these quasars arise from gravity and not from the Doppler effect.

If we increase M to $\sim 10^8 M_\odot$ and make it comparable to the mass found in galactic nuclei L is increased to $\sim 10^{44}$ erg s^{-1}. The distance of 20 kpc is correspondingly scaled up to ~ 20 Mpc. This distance is of the order of that considered by Hoyle and Burbidge (1966) in their version of the Doppler theory of quasars. Again, we emphasize that in the present theory, the redshifts are not of Doppler origin. It is possible for a massive galactic nucleus to eject a $10^8 M_\odot$-mass object without disrupting itself when we remember that the nucleus of a galaxy like M-87 is believed to contain a mass of $5 \times 10^9 M_\odot$ (Young *et al.*, 1978).

The presence of quasars of high redshifts near galaxies of low redshifts has been emphasized by the observations of Arp (1977, 1978) and also by Burbidge (1979). Such observations are naturally explained in the picture described here.

Our choice of $t_0 \sim 10^{12}$ s has been somewhat arbitrary. A larger value of t_0 of, say, 10^6–10^7 yr cannot be ruled out and is often associated with life-times of quasars and radio galaxies in other theories. A larger value of t_0 will make the luminosities higher, although the t_0-dependence of L in (30) is a weak one. We have chosen the lower value since it makes it just possible to observe one consequence of this model, viz. the change of a quasar redshift with time. For $t_0 \sim 10^{12}$ s, the average change in the redshift of a quasar is $\sim$ 1 part in 3×10^4 yr^{-1}. Improved spectroscopic techniques will make it possible to detect the change in the redshift of a quasar over a span of say 30 yr. It is, of course, possible, as seen from the (z, T_2) curves of Figure 1 that we may be lucky to spot a quasar when its redshift is changing at a more dramatic rate than implied by the average value of 10^{-12} s^{-1}. By contrast, the change in the redshift of a quasar under the cosmological hypothesis is at the much slower rate of $\sim 3 \times 10^{-18}$ s^{-1}.

4.2. Active galactic nuclei

In Paper I we had suggested that white holes in active galactic nuclei, like those in the Seyfert galaxies, act as sources of radiation, primarily in the X-rays. A massive oscillator, being a continuous source of energy (rather than a one-shot phenomenon) is more suitable as the energy machine in such nuclei. The similarity in the properties of quasars and the Seyfert nuclei (as well as the N-galaxies), including the property of strong X-ray emission reinforces our belief that the prime energy source in the two types of objects is the same.

The likelihood of a massive object in M-87 suggests that an oscillator with M in the range $5 \times 10^8 - 5 \times 10^9 M_\odot$ could well generate the high energy activity of many galaxies. For $t_0 \sim 10^7$ yr and $M \sim 10^9 M_\odot$, the total energy reservoir in the oscillator is $\sim 10^{61}$ erg and the luminosity is $\sim 10^{46}$ erg s^{-1}.

Oort (1977) has suggested that our own Galaxy contains expanding features in the nucleus which require an energy of $\sim 10^{55}$ erg at intervals of $\sim 10^6$ yr. A massive oscillator of $M \sim 5 \times 10^6 M_\odot$ could produce this nuclei energy at intervals of $t_0 \sim 10^6$ yr. A mass of this order is well within the upper limit of $10^8 M_\odot$ set by Oort from considerations of the overall dynamics of the Galaxy.

4.3. BL Lac-objects

The lineless objects now classified as BL-Lac objects have posed problems of interpretation which might be solved in the present picture. Why do such objects, with the exception of a handful, show no spectral lines? It is possible to argue that a large redshift ($z \gtrsim 20$, say) or a strong blueshift will make it difficult to identify any lines at all. Further, the continuum radiation from the gas heated by the ram-pressure of the expanding surface of the oscillator will tend to swamp any line emission from these objects. It is also worth investigating whether the lines seen in a few BL-Lacs could have been blueshifted.

In this connection it is worth mentioning that BL-Lacs appear to be present in the centres of massive elliptical galaxies (Miller *et al.*, 1978). A few years ago it was suggested (Hoyle and Narlikar, 1966b) that elliptical galaxies do not form as a result of contraction of pregalactic gas, but that they arise as clouds whose expansion is controlled by a central massive object at the nucleus. The observed lack of rotation in the ellipticals goes to support this picture. A massive oscillator at the nucleus with mass $\sim 10^7 M_\odot$ can effectively control the shape of an elliptical galaxy of mass 10^{11}–$10^{12} M_\odot$.

4.4. Stellar Oscillators

Since the onset of a black hole is believed to occur for masses exceeding $\sim 3 M_\odot$, it is natural to think of massive oscillators in the stellar mass range. For example, could the dark invisible object in the binary system of Cygnus X-1 be a massive oscillator instead of a black hole? In this case the high and low intensity episodes of the X-ray source (Oda, 1977) can perhaps be attributed to oscillations of the massive object. The gravity of the object would help forming an accretion disc although its maintenance and overall equilibrium might depend on the nature of oscillations.

5. Conclusion

We have considered the various uses that a massive oscillator may be put to in high-energy astrophysics. The introduction of a negative energy field is necessary to reverse the gravitational collapse of a massive object. We have shown that the C-field achieves this end and with suitable choices of parameters the collapsing object need not cross the event horizon.

The massive oscillator may be looked upon as alternating between a state when it is like (but not quite) a black hole and a state when it is like (but again, not quite) a white hole. During its white hole like expansion phase it can dump energy into surrounding gas which subsequently radiates. A mass range of $10 M_\odot$–$10^9 M_\odot$ includes energy sources for galactic X-ray sources, through quasars and BL-Lacs to active galactic nuclei. In the case of quasars a change of redshift with time is predicted.

It is interesting to note that while the massive oscillator converts the energy from the C-field reservoir into dynamical energy and dumps it into the ambient medium, it also can act as a centre of external accretion in the usual fashion. All the phenomena associated with accretion discs may therefore occur near the oscillator at least during its contracting phase.

It might be possible to extend the mass range further beyond $10^9 M_\odot$. Hoyle (1980) has suggested that non-singular white holes might be as large as super clusters of galaxies. In such cases it should be possible to apply the above model to study the 'primordial' production of light nuclei and to understand the observed photon/baryon ratio. We hope to consider this problem in a future paper.

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