IC/87/15

INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

ASTROPHYSICAL AXIONIC LASERS

I.I. Tkachev

INTERNATIONAL ATOMIC ENERGY AGENCY

UNITED NATIONS EDUCATIONAL, SCIENTIFIC AND CULTURAL ORGANIZATION

1987 MIRAMARE-TRIESTE

IC/87/15

International Atomic Energy Agency
and
United Nations Educational Scientific and Cultural Organization
INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

ASTROPHYSICAL AXIONIC LASERS *

I.I. Tkachev
International Centre for Theoretical Physics, Trieste, Italy
and
Institute for Nuclear Research of the USSR Academy of Sciences,
Moscow, 117312, USSR. **

ABSTRACT

The existence of a new type of astrophysical objects is suggested. These are clusters of weakly interacting scalar particles. The central part of such an object can be observed as a cosmic maser due to stimulated decays of the particles. The incident flow of particles from the peripheral area can supply the source with power. Cosmological strings naturally can form such a structure. In the axion case the energy release has the value typical for quasars and so a new model of quasars is proposed.

MIRAMARE – TRIESTE
January 1987

* To be submitted for publication.

** Permanent address.

1. INTRODUCTION

Weakly interacting scalar fields are a natural part of some modern field theories. These are, for example, axion [1], Polonyi fields [2], majoron [3], and so on. Usually, the interaction of these fields with matter is so weak that its detection in the laboratory seems hopeless (see, however, Refs. [4]). But, in any case, cosmology and astrophysics impose severe restrictions on the allowed range of interaction parameters. In this work, we shall show that weakly interacting (pseudo-)scalar fields can form astrophysical objects with peculiar properties. For definiteness we will consider only the axion field.

The axion interaction can be specified by fixing only one unknown parameter, f_a . Allowed values of this parameter are restricted if $f_a > 4\ 10^7$ GeV [5], then stars lose energy too rapidly [6] due to axion emission, but if $f_a < 10^{12}$ GeV then the axion contribution to the energy density of the Universe is too large [7]. At 10^{10} GeV $< f_a < 10^{12}$ GeV, there are enough axions to build up galactic halos and these is the axionic field which has very suitable properties to form the large-scale structure in the Universe [8]. In particular, axions are in a state with a very high phase space density; the axionic field is excited in a coherent state $\vec{p} \approx 0$ at a temperature T ~ 1 GeV due to non-perturbative QCD effects. Later on, axions practically decouple from matter.

The width of a single-particle axionic state is $\Gamma_a = \Gamma_\pi (f_\pi / f_a)^5$, where

$\Gamma_{\pi} = 7$ eV and $f_{\pi} = 93$ MeV are corresponding parameters of π^0 meson, so, the axion lifetime exceeds the age of the Universe. However, stimulated decays $a \to \gamma\gamma$ are possible at large phase space density. This process was considered in papers [7] as unessential in the hot Universe. Indeed, before the recombination moment the exponential growth of photons with energy $\nu_a \equiv \frac{m_a}{2}$ cuts in the cosmological plasma and after recombination the amplitude of the axionic field is so small that, due to the cosmological red shift ,photons leave the resonance zone too rapidly. However, in the epoch of galaxy formation the axion density increases in central parts of corresponding gravitational wells once again. As was suggested in [9], this could lead to the formation of cosmic masers with huge astrophysical activity. Here we consider this possibility.

II. LINEAR EVOLUTION OF AXIONIC PERTURBATIONS. A BRIEF REVIEW.

Considering the early stages of the evolution of perturbations in axions one may restrict ourselves to the "isothermal" ones only (i.e. to the perturbations which persists during some time only in axionic matter). Indeed, at scales $M < 10^{11} M_{\odot}$ only these ones are essential, because adiabatic perturbations have been suppressed [10], but the part of adiabatic perturbations that was reproduced in the axionic field in the QCD phase transition epoch survives, being now isothermal perturbations. Moreover, in the inflating universe scenario, isothermal perturbations can be generated directly [11].

The cosmological phase-space density of axions is very high [8]:

$$n_a(\mathrm{in}) \gg n_{a0} \equiv \frac{N_a(T \sim 1\,\mathrm{GeV})}{m_a^3} \approx \frac{f_a^5\, m_p^2}{M_{pl}\, m_\pi^6} \sim 10^{47} \qquad (1)$$

and concentrated in the vicinity $\vec{p} \approx 0$, i.e. axions are very cold. This leads to the following.

i) Axions can "stream out" perturbations only on scales which are far too small to be of astrophysical interest, namely, $M_{FS} \ll 10^{-3} M_{\odot}$ [8] (note that neutrinos "stream out" perturbations on mass scales $< 10^{15} M_{\odot}/(m_\nu/30\ \mathrm{eV})^2$ [12]).

ii) Being collisionless, axions do collapse in overdense regions. However, axionic perturbations do not grow till non-relativistic matter gives the dominant contribution to the average energy density [13] (i.e. at $f_a \sim 10^{12}$ GeV they do not collapse up to the temperature $T \sim T_a \equiv 10$ eV). That is, at the moment $T \sim T_a$ the magnitude of perturbations coincides with that reached at the moment when the corresponding scale enters the horizon (this is valid for all scales $M < 10^{11} M_{\odot}$). The corresponding spectrum determines the moment when the perturbation goes over to the non-linear stage of evolution. For example, in the case of Zeldovich flat spectrum, perturbations go over to the non-linear stage simultaneously. It is a flat spectrum that one obtains in inflationary scenarios. However, in papers [8] there have been concluded that axionic perturbations should rather have a Poisson spectrum with a characteristic mass scale of order $10^{7} M_{\odot}$.If stimulated decays of axions are unessential, then perturbations on this scale would collapse to black holes. However, at $M < 10^{10} M_{\odot}(10^{12}\ \mathrm{GeV}/f_a)^2$ stimulated decays in collapsing axionic "star" come into play before the system has locked itself in a black hole [9], this may prevent a hole formation.

A perturbation which enters the horizon with $\delta\varepsilon/\varepsilon < O(1)$ does not in general form a black hole and for clouds of axions we may use the standard picture of collisionless evolution. The whole picture is complicated and may be quite diverse. So, for definiteness we choose a single fashionable scenario of initial perturbations. Namely, we shall consider perturbations caused by the cosmological strings.

III. STIMULATED DECAYS OF AXIONS

Axionic matter is a medium inversely populated with respect to the decay $a \to \gamma\gamma$. Therefore, the radiation with frequency ν_a is enhanced when passing through this matter. The question arises as to what is stimulated emission rate?

The relevant quantity for such a problem (as well as for the problem of the collisionless collapse) is the density of particles in the phase-space $n^A(x^\mu, p_i)$, where A are intrinsic quantum numbers, x^μ are some coordinates in space-time, p_i are conjugated momenta. Note, that $g^{\mu\nu}p_\mu p_\nu = m^2$ and p_o is the integral of particle motion in the stationary metric. Using the geodesic equation

$$\frac{dp_\mu}{dt} = \frac{1}{2} g_{\nu\lambda,\mu} \frac{dx^\nu}{dt} p^\lambda \tag{2}$$

and the relation $p^\mu = p^o\, dx^\mu/dt$ we find

$$p^o \frac{dn}{dt} = \frac{\partial n}{\partial x^\mu} p^\mu + \frac{1}{2} \frac{\partial n}{\partial p_\alpha} g_{\mu\nu,\alpha} p^\mu p^\nu . \tag{3}$$

This is valid for photons as well as for axons. In accordance with Liouville-theorem theorem $dn/dt = 0$.

Dealing with interacting particles we introduce the collision integral in the right-hand side of the above equation. We shall consider only weak gravitational field

$$ds^2 = (1+2\varphi)\,dt^2 - (1-2\varphi)(dx^i)^2, \qquad \varphi \ll 1 , \tag{4}$$

so one can neglect gravity contribution into the collision integral. Then the kinetic equation describing the evolution of the photon distribution function takes the form

$$\frac{dn_\gamma(k,x)}{dt} = \int \frac{d^3p}{(2\pi)^3} \frac{d^3q}{(2\pi)^3} M(p,k,q)\Big[n_a(p,x)\big(1+n_\gamma(k,x)+n_\gamma(q,x)\big) - n_\gamma(p,x)n_\gamma(q,x)\Big] \equiv I_\gamma(k,x), \tag{5}$$

where $M(p,k,q) = (2\pi)^4 \delta^4(p-k-q)\, |M_{fi}|^2/8p^o k^o q^o$,

and M_{fi} is the Lorentz-invariant matrix element. A similar collision integral enters the axionic kinetic equation.

One can carry out the integrations explicitly using δ-functions and owing to the fact that $|M_{fi}|^2 \propto$ const here. Let decaying axion has the momentum p_i. Then energy k_o of one of the photons depends on the angle between $\vec{p}$ and $\vec{k}$ in the following way

$$m_a^2 = 2k_o\big[(1-2\varphi)p_o - |\vec{p}|\cos\theta\big]. \tag{6}$$

So, at fixed $|\vec{p}|$ axion decay products contain photons with energy k_o in the region $k_- \leq k_o \leq k_+$, where

$$k_\pm \equiv \frac{p_o \pm \sqrt{p_o^2 - m_*^2}}{2}, \tag{7}$$

$$m_* \equiv m_a (1 + \varphi) .$$

For further use let us consider collision integral in the space region where particle distribution is isotropic in the momentum space, i.e. $n = n(x^\mu, p_o)$. All axions with

$$p_o > \frac{4k_o^2 + m_*^2}{4k_o} \equiv p_{min}(k_o) \tag{8}$$

contribute to the change of the number of photons with energy k_o, so, we obtain for I_γ

$$I_\gamma^\sigma(x, k_o) = \frac{\Gamma_a m_a}{2k_o^2} \int\limits_{p_{min}(k_o)}^{\infty} dp_o \left\{ n_a(p_o)\left[1 + n_\gamma(k_o) + n_\gamma(p_o - k_o)\right] - n_\gamma(k_o) n_\gamma(p_o - k_o) \right\}. \tag{9}$$

We omit the index of photon chirality since this equation contain n_γ with the same chirality everywhere.

Now we are at the point where we must make some assumptions about the axionic distribution function. In weak gravitational field this function satisfies the equation

$$\frac{\partial n_a}{\partial x^\mu} p^\mu + m_a^2 \, \partial_i \varphi \, \frac{\partial n_a}{\partial p_i} = I_a(p, x). \tag{10}$$

Gravitational potential should be determined from the Poisson equation

$$\partial_i^2 \varphi = 4\pi \varkappa \mathcal{E}, \tag{11a}$$

where

$$\mathcal{E}(x) = m_a \int n_a(x, p) \frac{d^3p}{(2\pi)^3} . \tag{11b}$$

There is a lot of papers concerning star dynamics and the system (10), (11) (with $I_a = 0$, of course) have been carefully investigated, see e.g. [14]. This problem is not yet solved. In stationary state the density in phase space is an arbitrary function of motion integrals $n = n(p_o, L^2, L_z)$ and that is why even stationary states can not be investigated completely. Moreover, unlike thermodynamic, final stationary state in such selfgravitating systems depends upon initial non-equilibrium one (we know that there exist very different spherical galaxies and spirals). However, one can think that phase space density in the case of our interest is still non-stationary. The most rapid relaxational process is the "violent" relaxation (which occurs due to time dependent mean field φ) with the relaxation time of order of the cross-over time of the system, and, if the axionic astrophysical laser does work, it is most likely that the period of its brightest shining is also of order of the cross-over time.

For simplicity we consider spherically-symmetric configuration . For our purposes axionic configuration with practically radial motion of particles at sufficiently large distance from the center is the most appropriate. However, due to the velocity dispersion the density n_a is isotropic within some region near the center. We discuss below the scenario of formation of such a configuration assuming that cosmological strings are seedings of gravitational condensation.

So, for the axion energy density we admit the following profile at $r< R$ with some $R >> r_c$

$$\mathcal{E}(r) = \frac{M_{Pl}^2 v_o^2}{4\pi r_c^2 (1 + r^2/r_c^2)} , \tag{12}$$

where v_o coincides with the circular velocity at the orbit $r_c < r < R$. Note, that observational data indicate that the same energy density profile with $v_o \sim 10^{-3}$ should be characteristic for dark matter in extended galactic halos [15]. There are observational data for the central part of the Milky Way in the range 0.1 pc $< r < 1$ kpc which also give the r^{-2} dependence for $\mathcal{E}(r)$ with $v_o \sim 10^{-3}$ [16]. However, there is a lot of distribution functions $n(x,p)$ which give the law (12). For example, such an energy density will be produced by "isothermal" distribution $n \propto \exp(p_o / \theta)$.

For simplicity, we shall place all the axions on a single energy level. Inside the isotropic core region we have

$$n = n_o \, \delta (E - \varphi(R)) , \tag{13}$$

where $$E \equiv \frac{v^2}{2} + \varphi(R) .$$

Comparing now the expressions (11b) and (9) for energy density and collision integral we obtain in the core region (assuming $n_\gamma(k_o) = n_\gamma(p_o - k_o) << n_a(p_o)$)

$$\frac{dn_\gamma}{dt} = \frac{\Gamma_a \pi^2 \mathcal{E}}{m_a^4 v_e} (1 + n_\gamma) , \tag{14}$$

where $v_e(r)$ is the escape velocity, $v_e \equiv \sqrt{-2\varphi}$, which characterize the depth of the gravitational well. In fact, this relation up to a numerical factor is valid for a wide class of distribution functions n_a.

Let us introduce the amplification coefficient

$$\mathcal{D}(k_o , r) = \int_0^r I_\gamma (k_o , r', n_\gamma = 0) \, dr' . \tag{15}$$

Applying (7) we see that the width of resonance zone, $\Delta k = k_+ - k_-$, is equal to $\Delta k(r) = \sqrt{p_o^2 - m_*^2(r)}$. So, the deeper is the gravitational well, the wider is the resonance zone at fixed axion energy p_o. One may neglect the growth of photon number density outside the core region and a maximal amplification is achieved for photons inside the band $\Delta\nu \equiv \Delta k(r_c) = mv_e$ with

$$\mathcal{D} = \frac{\Gamma_a \pi^2 \mathcal{E}(0) r_c}{m_a^4 \, v_e(0)} . \tag{16}$$

Note, that one may rewrite (16) in the form familiar in ordinary lasers theory

$$\mathcal{D} \sim \frac{\Gamma_a N}{\nu_a^2 \Delta\nu} r , \tag{17}$$

where $N = N_2 - N_1$, and N_1, N_2 are populations of the ground and exited levels (in our case $N = N_2 = \mathcal{E}/m_a$). However, owing to the fact that an axion decays into two photons (with "opposite" momenta), a bounded axionic medium effectively behaves as inversely populated atomic matter placed in a resonator. Indeed, let one inject N photons with a resonance energy into axionic matter from the left. The increment of the photon number density with the same momenta is $\Delta N = N \exp(D)$. But it is also the number of emerged photons with the opposite momenta. Therefore, at $D > 1$, the number of photons moving from the right in the next cycle exceeds that of injected photons, and so on. Therefore, in the region where the condition $D > 1$ is satisfied, n_γ increases exponentially with time (as in an infinite medium) till the condition $D = 1$ is reached.

The formulae (16) gives $D \to \infty$ for the axions at rest ($v \to 0$). This mean that the kinetic assumptions, which was made, breaks here. To obtain the correct result one may investigate the creation of photons by the classical axionic field $a = A \cdot \sin(m_a t)$. Then one have [7,9]

$$\frac{dn_\gamma}{dt} = \alpha m_a \frac{A}{f_a} \sqrt{n_\gamma (1+n_\gamma)}\ ,$$

$$\Delta\nu = 2\alpha m_a \frac{A}{f_a} \qquad (18)$$

and

with being QED fine structure constant. Note, that the expression for the inverse amplification length may be rewritten here in the form (17) as well, with N being now $N = m_a A^2$.

However, in the case of our interest the formulae (16) do work, and we are searching now for a configuration which would yield $D_{core} > 1$. Then, we can expect $L_\gamma \sim L_a(in)$, where $L_a(in)$ is the value of the axionic energy inflow in a halo and L_γ is the total photon luminosity of the object. With energy densi-

ty profile (12) and with the assumption of pure radial axionic flow we get $L_a \approx M_{Pl}^2 v_0^3$. For a typical value $v_0 \sim 10^{-3}$ we obtain $L_\gamma \sim 10^{17} L_\odot$ which is even grater than the maximum of the observed luminosity of quasars ($L_{\gamma\,max} \sim 10^{14}\ L_\odot$).

It is the core region which is responsible for the stimulated emission, but the axionic inflow from the halo region supplies the core with power.

Let us establish the conditions whose fulfillment ensures $D > 1$. The Eq. (16) can give the deceptive feeling that the grater is r_c, the larger is D. However, in the case of selfgravitating system, r_c can not be arbitrarily large, because r_c, $\mathcal{E}$, and v_e are related to each other. Using (12) we obtain

$$\mathcal{D} \sim \frac{\Gamma_a v_e M_{Pl}^2}{r_c\, m_a^4} = \frac{\Gamma_\pi M_{Pl}^2 v_e}{r_c\, m_\pi^4}\left(\frac{f_\pi}{f_a}\right). \qquad (19)$$

It is splendid that D depends upon (f_π/f_a) only linearly. The requirement $D > 1$ gives

$$r_c < 10^7 v_e \left(10^{12} GeV/f_a\right)\ cm. \qquad (20)$$

(At $v_e = 10^{-1}(f_a/10^{12}$ GeV) the expression on r.h.s. of (20) coincides with the radius of a neutron star). So, to obtain stimulated emission we must concentrate axions in the core with a very small radius. Was it possible in the Universe ?

Since axion phase-space density is conserved (up to "collisions") we find (using (1) and (12)) that spherically-symmetric axionic perturbation produces a core with

$$r_c \simeq \frac{M_{pl}}{m_p m_a} \sqrt{\frac{M_{pl} v_o}{f_a\, n_a(in)}} \simeq 10^{10} \sqrt{v_o 10^{12}\,GeV/f_a} \sqrt{n_{a0}/n_a(in)} \quad cm. \tag{21}$$

This does not contradict to the inequality (20) in view of $n_a(in) \gg n_{a0}$.

IV. STRING-INITIATED FORMATION OF STRUCTURES

Now we are going to discuss concrete abilities of that perturbations which are produced by cosmic strings. The cosmological evolution of strings, which may appear as a result of some phase transitions in the early Universe, has been widely discussed (see [17]). Strings have a mass per unit length $\sim \mu^2$, where μ is the characteristic energy scale of symmetry breaking. Network of strings at a time t produces a number of oscillating loops of size $r_s \sim t$.

Following [18] one may approximate the gravitational effect of the string loop by that of a spherical shell with radius r_s and mass $2\pi\mu r_s$. Let us assume that the field of velocities of the axions describes a pure cosmological expansion at the moment $T = T_a$. Then the evolution of axions can be described in the spatially flat Friedman universe by the equation [19],[18] :

$$\frac{1}{2}\dot{r}^2 - \frac{\varkappa(M + m_s)}{r} = E, \tag{22}$$

where M is the mass of axions within the radius r(t). For those

axions which was inside a loop at a moment $T = T_a$ we have $E = -\frac{\varkappa m_s}{r_s} = -2\pi\mu\varkappa$. This leads to a flat rotational curves with $v_o = \sqrt{4\pi\mu}\, M_{pl}^{-1}$ in the corresponding part of a halo after virialization, i.e. a local energy density contrast will take a form (12), $\mathcal{E}(r) = \mu r^{-2}$. This would correspond to observational data if

$$\varkappa\mu \sim 10^{-7}. \tag{23}$$

Axions which at $T = T_a$ occurred outside the loop have $E_a(r_o) = -\varkappa m_s/r_o$ and produce a power law $\mathcal{E} \propto r^{-9/4}$ [19].

If we assume, that those axions which initially were inside the loop form the core satisfying Eq.(21), then we obtain the object we desire. Indeed, the loop may contain as many axions as were needed to construct extended galactic halos [18] and they would maintain energetics of the core. However, it is natural to assume that in this case r_c is equal not to its possible minimal value (21) but to the loop size [18], because rapidly moving string can intensively perturbate the axion trajectories (this question needs special investigation). Following this way we obtain that a halo which can accommodate $M \sim 10^6 M_\odot$ (this would correspond to the typical quasar lifetime $t \sim 10^6$ year with typical luminosity $L_\gamma \sim 1\ M_\odot$/year) do correspond to $r_c \sim 10^{18}$ cm which is too large (see (20)) to provide stimulated emission.

Nevertheless, one may think that string motion does not disturb axions in the distant regions because the effect of higher multipole moments will be small there [20] and the core which those axions will form can reach its minimal radius (21), or be even smaller due to gravitational attraction to the loop. So, let us consider axions inflow from the region with $\mathcal{E} \propto r^{-9/4}$. The expression (19) for the enhancement coefficient is still valid but we should know what value of the luminosity this region could provide.

Axions which were at a distance r_i at a moment T_a will produce an energy flux $L_i = 4\pi\mathcal{E}\, r_i^2 (dr_i/dt_c)$, where t_c is the recollapse time for those axions

$$t_c = \frac{2\pi æ M_i}{(2E_i)^{3/2}} . \tag{24}$$

This gives $L_i \approx M_{pe}^2 \, E_i^{3/2}$, or

$$L \approx M_{pe}^2 \left(\frac{T_a}{M_{pe}}\right)^{4/3} \frac{m_s}{M_{pe}^{4/3}\, t_c^{1/3}} \sim M_{pe}^2\, 10^{-15} \left(\frac{m_s}{M_\odot}\right)\left(\frac{year}{t_c}\right)^{1/3}, \tag{25}$$

where

$$t_c \gtrsim (æ\mu)^{-9/2} \left(\frac{T_a}{M_{pe}}\right)^4 \frac{m_s^3}{M_{pe}^4} .$$

String loops decay by the gravitational wave emission [17]. Typical lower mass of the loop surviving to the moment $T = T_a = 10$ eV is equal to $m_{s\,min} \sim 10^5 M_\odot$. From the other hand this is the mass of the most probable loop because the size spectrum of loops is estimated to be $n_s \sim (t^2 m_s)^{-1}$ for $m_{s\,min} < m_s < 2\pi\mu t$.

With $t_c \sim 10^6$ year and $m_s \sim 10^5 M_\odot$ the formulae (25) gives for L the value of the maximal observed quasars luminosity $L \sim 10^{14} L_\odot$.

Of course, the estimate (25) is valid not only for the case of structure formation initiated by the strings but also in the case of compact massive object of any other nature which was present at $T \sim T_a$ (the rate of the stimulated axionic decays in the course of their accretion on the black hole should be estimated separately).

Formulae (25) predict too slow t_c dependence, i.e. at $t_c \sim 10^{10}$ year there will be too many bright quasars. However, the model under consideration is oversimplified. The loop moving with the velocity $v \sim |E|$ with respect to the axionic background can not capture axions with $E_i > E$. This will cut luminosity at some t . From the other hand, there are mechanisms [20] which can slow down the loop to the speed at which it can absorb a large halo. To collect a halo with $M \sim 10^6 M_\odot$ for a string with $m_s \sim 10^5 M_\odot$ it is sufficient at $T = T_a$ to capture axions within radius which is only by an order of magnitude grater then the string radius. At the moment T_a the velocity of the loop is slowered by the Hubble expansion down to the value $v \sim v_i (r_s/t_a)^{1/2}$ [20], so to capture those axions it is sufficient for string to have initial velocity of order $v_i \sim 1/3$. These values seem to be reasonable for model to work.

It has been argued [20] that the matter accretion on the loop could produce a black hole in the center of (every) galaxy. Applying the same arguments we may conclude that axions accretion on strings will result to a core with very high stimulated emission rate. It is worth noting that the same ingredients (axions and strings) which have currently been used in explanations of galactic-scale , and large-scales structure in the Universe can lead to the quasar-like energy release.

Objects with smaller luminosity and longer life-time could also emerge in this model. To study this possibility more careful calculations are needed, e.g. one should investigate the growth of axion velocity dispersion in a course of real galaxy evolution.

Thus, we predict objects with energetics which is typical for quasars. Could we identify them with quasars?

V. A MODEL OF QUASARS

A pure axionic object will produce a monochromatic spectrum with $\nu = \nu_a$.

The spectrum of quasars is very complicated with characteristic emission lines and non-thermal continuous part. One may hope that their spectra can be produced if the axionic core with high luminosity is surrounded by ordinary matter [9] (recently, the idea that axion decays could be responsible for the quasar phenomena have been also expressed in [21]). Being in the center of the gravitational well axionic core should be, in general, obscured by the barionic matter. It is the radiation pressure which could give then an observational quasar features.

Here we consider some peculiarities of matter dynamics in the axionic halo of such objects. There are two main forces acting on the particle with mass m placed at a distance R from the center of the axionic object: the gravity, taking in the halo the form $F_g = mv_o^2/R$, and the radiation pressure, which is equal to $F_{rp} = \sigma L/4\pi R^2$. Since these forces have a different R-dependence, the point $R = R_{eq}$ where the particle is in equilibrium can exist: $R_{eq} = \sigma L/4\pi m v_o^2$. For a plasma, we mast take $m = m_p$, since the proton is the heaviest particle; and $\sigma = \sigma_T$, where $\sigma_T = 8\pi e^4/3m_e^2$ is the Thompson cross section for the electron. If we take for L the quasars luminosity we obtain that $R_{eq\,p}$ exceeds the size of a galaxy. We conclude that the plasma is accelerated away from the object.

For a star with mass M and radius r we obtain $R_{eq\,s} \sim 10^6\, r_\odot\, v_o \cdot$

$(M_\odot/M)\ (r/r_\odot)^2$, taking $\sigma = 4\pi r^2$ and, as before, assuming quasars luminosity. That is, the star could be on a circular orbit in the immediate vicinity of the core with isotropic emission. If the orbit satisfies the inequality $R < \sqrt{R_{eq\,p}\ r}\ \ (v_o/v\ )$, where v is the escape velocity from the star then the radiation pressure, acting on a plasma, exceeds the gravitational attraction to the star. The stars can supply the zone of the halo with plasma which, when accelerated by the radiation, could produce the spectra of the "quasars". This is a main mechanism for producing the complicated spectra in the model. Absorbing the radiation the stars themselves expand and cool in accordance with the virial theorem, and these regions of matter at some circumstance could give characteristic contribution to the total spectrum.

So far we have considered spherically symmetrical configuration. However a normal quasar has relativistic jets which are very narrow, cold and stable. One may hope to explain such a feature if the radiation from stimulated decays of axions is concentrated in beams [9]. The core which have a spindle-like form will produce such a radiation in a natural way. In this case axion distribution function should be of the type $n_a = n_a(p_o,\ L^2,\ L_z)$.

One may hope that conditions for the stimulated emission is more comfortable in this case because in non-spherical geometry the relation between $v_e, \mathcal{E}$, and the length of the major axis r is not so tight. Note also, that jets are so cold that their spectra contain emission lines. Under such circumstances one may speculate that some resonance phenomena (not necessary at the axion frequency) could play a (crucial) role stabilizing, squeezing and cooling a jet.

Is the spindle-like axionic core unusual? As concerns the barionic matter we know that spiral galaxies have two components:

spherical halo and thin disk with spirals. The way in which the real collapse of the axionic matter will proceed remains to be studied yet.

VI. CONCLUSION

We have shown that the Bose-fields with very long-lived quanta can form gravitationally bound objects with a very high luminosity due to the process of stimulated decays. At present the axionic field seems to be an excellent (even the best) tool for achieving this. We can conjecture that the acceleration of matter by a coherent maser beam is the origin of quasar phenomena. In view of the fact that this question is yet unexplored, by "quasars" I mean here an active galactic nuclei, BL Lacertae and SS 433 object also . In the dynamical center of our Galaxy there is a radiosource of non-thermal radiation Sgr A West with the size less than 10 a.u. I hope that the specific spectra of these objects, the extreme narrowness and low temperature of relativistic jets, and the evolutionary features can be naturally explained in this model. The discovery of pure cosmic masers or the lines in discrete spectra of quasars which cannot be identified with any molecular or atomic levels, would be of crucial importance for the model, and then the axion mass will be precisely determined. The most likely frequency band for the search is the intermediate one between radio- and infrared astronomy.

On the other hand, identification of these objects may turn out to be the only way to obtain evidence for the existence of very weakly interacting Bose-fields.

ACKNOWLEDGMENTS

I am grateful to V.A.Berezin, V.S.Berezinsky, A.Yu.Ignatiev, V.A.Kuzmin , N.G.Kozimirov, V.A.Rubakov, H.Sato, M.E.Shaposhnikov and M.V.Sazhin for helpful discussions. I wish to thank L.Abbott for bringing to my attention the preprint [21]. I would like to thank Professor Abdus Salam, IAEA and UNESCO for kind hospitality at the International Centre for Theoretical Physics, Trieste, where this work was completed.

REFERENCES

1. R.D.Peccei,H.Quinn, Phys.Rev.Lett. 38 (1977) 1440
   S.Weinberg, Phys.Rev.Lett. 40 (1978) 223
   F.Wilczek, Phys.Rev.Lett. 40 (1978) 279
2. J.Polonyi, Budapest preprint KFKI-1077-83 (1977)
3. Y.Chikashige, R.N.Mohapatra, R.D.Peccei, Phys.Lett 98B (1981) 265
   G.B.Gelmini, A.Roncadelli, Phys.Lett. 99B (1981) 411
4. V.A.Kuzmin, M.E.Shaposhnikov, I.I.Tkachev, Sov.Phys.-JETP Lett. 36 (1982) 49
   P.Sikivie, Phys.Rev.Lett. 51 (1983) 1415
5. M.Fukugita, S.Watamura, M.Yoshimura, Phys.Rev.Lett. 48 (1982) 1522
6. G.V.Domogatsky, Sov.Nucl.Phys. 8 (1968) 759
7. L.Abbott, P.Sikivie, Phys.Lett. 120B (1983) 133
   J.Preskill, M.Wise, F.Wilczek, Phys.Lett. 120B (1983) 127
   M.Dine, W.Fischler, Phys.Lett. 120B (1983) 137
8. J.Ipser, P.Sikivie, Phys.Rev.Lett. 50 (1983) 925
   F.W.Stecker, Q.Shafi, Phys.Rev.Lett. 50 (1983) 928
   M.S.Turner, F.Wilczek, A.Zee, Phys.Lett. 125B (1983) 35
9. I.I.Tkachev, Sov.Astron.Lett. 12 (1986) 726
   I.I.Tkachev, in: Proceedings of the Quarks-86 Seminar, Tbilisi, April 1986; VNU Science Press, in press.

10. J.Silk, Astrophys.J. 151 (1968) 459
11. A.D.Linde, Sov.Phys.-JETP Lett. 40 (1984) 496
12. J.Bond, G.Efstattcon, J.Silk, Phys.Rev.Lett. 45 (1980) 1980
13. P.Meszaros, Astron.Astrophys. 38 (1975) 5
14. A.M.Fridman, V.L.Polyachenko, Physics of Gravitating Systems, Springer-Verlag, 1984, v.1,2
15. J.R.Primak, Dark matter, galaxies and large scales structure in the Universe, SLAC-PUB-3387, 1984
16. J.H.Oort, Ann.Rev.Astron.Astrophys. 15 (1977) 295
17. A.Vilenkin, Phys.Reports 121 (1985) 263
18. H.Sato, Modern Physics Lett. A1 (1986) 9
19. J.E.Gunn, Astrophys.J. 218 (1977) 592
20. C.J.Hogan, Phys.Lett. 143 (1984) 87
J.Silk, A.Vilenkin, Phys.Rev.Lett. 53 (1984) 1700
21. T.W.Kephart, T.J.Weiler, Axion stars, preprint VAND-TH 86-2