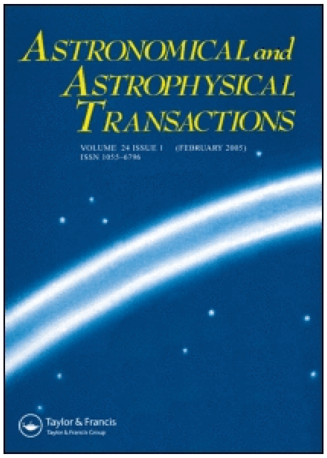

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ASTROPHYSICAL CONSEQUENCES OF THE EXISTENCE OF GOLDSTONE BOSONS

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(2 December 1992)

We suggest routine astronomical methods of searching for light Goldstone bosons (axions) including massless arions. The basic idea is to use the conversion of photons into low-mass bosons in the magnetic fields of compact stars and also of interstellar and intergalactic media. This process depends strongly on the polarization state of a photon and may produce a noticeable amount of linear polarization. Polarimetric observations may give strong constraints on the axion masses and the axion coupling constant. The best candidates for the observations are core-emission pulsars, low-mass X-ray binaries, magnetic white dwarfs including AM Her type objects. For the interstellar medium, this effect can be comparable with the interstellar polarization. X-ray fluxes from quasars and AGNs may oscillate with the cosmological redshift z due to the conversion of X-ray photons into arions.

KEY WORDS Elementary particles—galaxies, general—stars, general.

1. INTRODUCTION

Astrophysical and cosmological considerations have led to numerous constraints on new Particle Physics phenomena. The conservation of the CP symmetry in strong interactions has been a long-standing puzzle of Particle Physics in view of CP-violating effects observed in the K^0-meson system. Axions have been proposed to solve the strong CP problem (t'Hooft 1976; Jackiw and Rebbi 1976; Callan *et al.*, 1976), i.e., to answer the question why the neutron electric dipole moment is so small compared to the natural value predicted by QCD (Baluni 1979; Crewther *et al.*, 1979). QCD predicts a value of roughly $d_n = 5 \times 10^{-16}\,\Theta$ e cm, where Θ is the parameter which determines the strength of the CP invariance violation. The experimentally determined value is $d_n \leq 10^{-25}$ e cm. This leads to the constraint $\Theta \leq 10^{-9} \div 10^{-10}$.

To explain naturally why Θ should be small, Peccei and Quinn (1977) introduced an approximate global symmetry into the QCD Lagrangian making Θ a dynamical variable. A necessary consequence of the spontaneous breaking of the approximate PQ symmetry is the generation of a nearly massless pseudoscalar particle dubbed the axion. The axion is a pseudo-Goldstone boson which is called arion or omion if the PQ symmetry had been exact before spontaneous symmetry breaking occurred (Anselm and Uraltsev 1982; Sikivie 1988). The PQ symmetry also arises naturally in theories of supersymmetry and superstrings (Davis 1986).

Observational and experimental searches for the axions (arions or omions) depend crucially on the symmetry-breaking scale f_a, because both the mass and

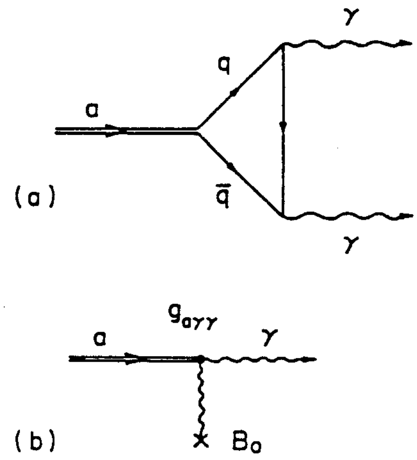

Figure 1 (a) Two-photon coupling of the axion through a triangle anomaly; (b) Axion-photon transition in magnetic field.

the couplings of axion are related to f_a. It is therefore quite important to get bounds on f_a which is *a priori* a free parameter. Since for nonmassless axions we have $f_a \sim 1/m_a$, this is equivalent to obtaining bounds on the axion mass.

There are two generic types of axions. The DFSZ axions couple to both quarks and leptons at the tree level (Dine *et al.*, 1981). The hadronic, or the KSVZ axion has no tree-level couplings to leptons but does couple to them through loop diagrams (Kim 1979; Shifman *et al.*, 1980).

Astrophysical and cosmological arguments have been extremely useful in narrowing the allowed range of the axion mass. Due to many clever arguments there are only three "windows" left open for the axion: $10^{-6}\,\mathrm{eV} \leq m_a \leq 10^{-3}\,\mathrm{eV}$, for the hadronic axion only $3\,\mathrm{eV} \leq m_a \leq 8\,\mathrm{eV}$ and $m_a = 0$ for the arion (omion) (see, e.g., Turner 1990; Reffelt 1991; Ressel 1991). Observations from the recent supernova SN 1987A provide the most stringent upper limit $m_a \leq 4 \times 10^{-4}\,\mathrm{eV}$ (Turner 1990; Raffelt 1991; Mayle *et al.*, 1988; Burrows *et al.*, 1989).

The detection of axions is possible through their coupling to photons. The direct annihilation of axion into two photons (Figure 1a) has a finite although extremely low cross section (Wuensch *et al.*, 1989; Ressel *et al.*, 1991). However, the interaction with magnetic field via the Primakoff process can lead to the photon production which energy is comparable to the axion total energy (Figure 1b). The reverse process, i.e., photon-axion transitions in magnetic field also exists.

We shall discuss in detail only one problem: the indirect effects of photon propagation in magnetic fields of many astrophysical objects including neutron starts, magnetic white dwarfs and magnetic Ap stars, and also interstellar and intergalactic media. The process of photon-axion conversion acts as an additional absorption process strongly depending on the polarization state. Therefore, it can lead to noticeable photometric and polarimetric effects for stellar radiation which propagates across magnetic field lines. Hereafter we restrict ourselves basically to the photon-arion conversion.

2. THEORY OF THE PHOTON-AXION CONVERSION IN MAGNETIC FIELD

The theory of the photon-axion transitions in magnetic field was developed in detail by Raffelt and Stodolsky (1988). Their main equation for the description of the photon-axion system is given by

$$\left[\omega + \begin{pmatrix} \Delta_\perp & 0 & 0 \\ 0 & \Delta_\parallel & \Delta_F \\ 0 & \Delta_F & \Delta_a \end{pmatrix} - i\frac{\partial}{\partial z}\right]\begin{pmatrix} A_\perp \\ A_\parallel \\ a \end{pmatrix} = 0, \tag{1}$$

where $A_\perp$, $A_\parallel$ and a are the amplitudes of the perpendicular and parallel photon states and the axion state, respectively, and

$$\begin{aligned}
\Delta_{\perp,\parallel} &= \Delta^v_{\perp,\parallel} + \Delta^p_{\perp,\parallel},\\
\Delta^v_\perp &= 2\omega\delta\sin^2\phi,\\
\Delta^v_\parallel &= \tfrac{7}{2}\omega\delta\sin^2\phi,\\
\Delta^p_\perp &= -\left(\frac{\omega_p^2}{2\omega}\right) f_\perp(\omega, T_e, B),\\
\Delta^p_\parallel &= -\left(\frac{\omega_p^2}{2\omega}\right) f_\parallel(\omega, T_e, B),\\
\Delta_a &= -m_a^2/2\omega,\\
\Delta_F &= (B/2F)\sin\phi,\ F \text{ is the axion coupling constant}\\
\delta &= \frac{\alpha}{45\pi}\left(\frac{B}{B_{cr}}\right)^2, \qquad B_{cr} = \frac{m_e^2}{e} = 4.414\times 10^{13}\,\mathrm{G}.
\end{aligned} \tag{2}$$

Here the natural Lorentz–Heaviside units are used, i.e. $\hbar = c = 1$ and α is the fine-structure constant. The Δ factors represent the momentum differences of the respective normal modes j compared to the photons of the same energy in a field-free vacuum. In the absence of the axion terms, there are two independent normal modes (Adler 1971; Meszaros and Ventura 1978; Pavlov and Gnedin 1984) characterized by the quantities $\Delta_{\perp,\parallel}$. Subscripts $\perp$, $\parallel$ refer to the photon polarization states with the electric vectors predominantly perpendicular or parallel to the plane containing the magnetic vector and the wave vector k_j, $\Delta_j = k_j - \omega$, and superscripts p and v refer to the plasma and vacuum QED contributions, respectively. ω_p is the electron plasma frequency. The angle ϕ is that between **k** and **B**. Adding the photon-axion interaction term to the Lagrangian introduces the third (axion) mode Δ_a as well as a non-diagonal Δ_F mixing term that couples the axion a and photon $\parallel$ modes. The functions $f_{\parallel,\perp}$ are of order unity and depend on temperature T_e. It is very important that the inverse quantities Δ^{-1} are the natural length scales of the problem and are interpreted as the oscillation lengths. Another important thing is different signs between different Δ_j (for example, vacuum $\Delta^v_\parallel$ and plasma $\Delta^p_\parallel$, $\Delta^v_\parallel$ and Δ_a, etc). This fact allows us to expect resonant effect that should probably enhance the photon-axion conversion process.

The solution of the wave Eq. (1) for a region where the field is approximately homogeneous is accomplished by Fourier transformation, which gives the dispersion relation $k_\perp = n_\perp \omega$ for the first component, while the second and the third components can be diagonalized by a rotation through the mixing angle θ given by (Raffelt and Stodolsky 1988)

$$\tfrac{1}{2}\tan 2\theta = \frac{\Delta_F}{\Delta_\parallel - \Delta_a}. \tag{3}$$

For arions, (3) transforms into a simpler form:

$$\tfrac{1}{2}\tan 2\theta = \frac{\Delta_F}{\Delta_\parallel}, \tag{4}$$

for a weak mixing $\theta \approx \Delta_F/(\Delta_\parallel - \Delta_a)$.

For strong mixing, the transition rate is (Raffelt and Stodolsky 1988):

$$P(\gamma_\parallel \to a) = \sin^2(\Delta_F l), \tag{5}$$

where l is the size of region where magnetic field is approximately homogeneous. The oscillation length is

$$L_{\text{osc}} = \frac{\pi}{\Delta_F} = \frac{2\pi F_a}{B_\perp}, \tag{6}$$

where $B_\perp = B \sin\phi$.

For the weak mixing case:

$$P(\gamma_\parallel \to a) = 4\theta^2 \sin^2(\tfrac{1}{2}\Delta_{\text{osc}} l), \tag{7}$$

where

$$\theta = \frac{2\omega\,\Delta_F}{2\omega\,\Delta_\parallel^v + 2\omega\,\Delta_\parallel^p + m_a^2}, \qquad \Delta_{\text{osc}} = \Delta_\parallel - \Delta_a, \tag{8}$$

and the oscillation length is $L = 2\pi/\Delta_{\text{osc}}$.

Now we consider separately the case of arions ($m_a = 0$) and rewrite Eqs. (5)–(8) in a more convenient form:

$$P(\gamma_\parallel \to a) = A^2 \sin^2\left(\frac{l}{L_{\text{osc}}}\right), \tag{9}$$

where $A = l_p/l_a$, $L_{\text{osc}} = \min\{l_0, l_a\}$ and

$$\begin{aligned} \frac{1}{l_0} &= \left|\frac{1}{l_v} - \frac{1}{l_p}\right|, \\ l_p &= \frac{2\omega}{\omega_p^2}, \qquad l_v = \frac{2}{7}\frac{1}{\omega\delta\sin^2\phi}, \\ l_a &= \frac{2F_a}{B_\perp}. \end{aligned} \tag{10}$$

The amplitude of the photon-arion transition is $A \approx 1$ if $l_a < l_p$ and $A \approx l_0/l_p$ if $l_0 < l_p$. In the more common case ($m_a \neq 0$) one should consider one more oscillation length $l_m = 2\omega/m_a^2$.

For the weak-mixing case, when $l_a \gg l_0$, l_m, Reffelt and Stodolsky (1988) have also considered birefrigence effects to search for axions. Then the amplitude $A \ll 1$ can be approximated as

$$A = \frac{2B_\perp \omega}{F_a m_a^2}. \tag{11}$$

The ellipticity ϵ and rotation angle φ are:

$$\begin{aligned} \epsilon &= A^2\left[1 - \frac{\sin(l/L_{osc})}{l/L_{osc}}\right]\frac{l}{L_{osc}}, \\ \varphi &= \tfrac{1}{2}A^2 \sin^2(l/2L_{osc}). \end{aligned} \tag{12}$$

The case of very small axion masses and (or) essential plasma (or magnetized vacuum) contribution $l/L_{osc} \ll 1$ is especially interesting. The results of (12) can be expanded as

$$\begin{aligned} \epsilon &= \frac{(B_\perp m_a)^2}{48\omega F_a^2} l^3, \\ \varphi &= \frac{B_\perp^2 l^2}{8F_a^2}. \end{aligned} \tag{13}$$

From (13) it is clear that the ellipticity is produced only for nonmassless axions. For arions, one should observe only the rotation of the electric vector of the propagated radiation.

3. THE PHOTON-ARION TRANSITION IN MAGNETIZED COMPACT STARS

The probability of the photon-arion transition is derived as a product of magnetic field strength $B_\perp$ and the size of the magnetic field region. Therefore, magnetized compact stars (neutron stars and magnetized white dwarfs) are probably good candidates for the observation of this effect. One should expect the excess of linear polarization due to the photon-arion conversion and the predominant direction of the electric field lies perpendicularly to the plane of the magnetic dipole axis of the star and the line of sight.

Our first step is an estimation of the spectral range within which the strong photon-arion transition exists. In the high-frequency range, the vacuum polarization effect can be predominant. This fact restricts the high-frequency range by the condition

$$\omega \ll \omega_{max} = \frac{5 \times 10^{27}}{B_\perp F_a}. \tag{14}$$

At low frequencies, the plasma effect should suppress the photon-arion transition. It requires the following condition on frequency to strengthen the photon-arion conversion:

$$\omega \gtrsim \omega_{min} = 5 \times 10^{-8} \frac{F_a N_e}{B_\perp}, \tag{15}$$

where N_e is the electron number density.

The final condition of a strong photon-arion transition is the oscillation length being comparable to the size of the compact star R_s:

$$L_{\rm osc} \approx l_a \leq R_s. \tag{16}$$

Now we use Eqs. (14)–(16) for analysis of the concrete situations.

3.1. *Fastly Rotating Neutron Stars—Radiopulsars*

The conditions (14)–(16) are

$$\begin{gathered}
\omega \leq 3 \times 10^{-5}\left(\frac{10^{12}\,\mathrm{GeV}}{F_a}\right)\left(\frac{10^{12}\,\mathrm{G}}{B_s}\right)\left(\frac{R}{R_s}\right)^3 \mathrm{eV}, \\
\omega \geq 1.8 \times 10^{-11}\left(\frac{F_a}{10^{12}\,\mathrm{GeV}}\right)\left(\frac{10^{12}\,\mathrm{G}}{B_s}\right)\left(\frac{R}{R_s}\right)^3\left(\frac{N_e}{1\,\mathrm{cm}^3}\right) \mathrm{eV}, \\
\left(\frac{B_s}{10^{12}\,\mathrm{G}}\right)\left(\frac{F_a}{10^{12}\,\mathrm{GeV}}\right)\left(\frac{10^6\,\mathrm{cm}}{R}\right) \leq 1,
\end{gathered} \tag{17}$$

where R is the distance from the NS.

If the conditions (17) are fulfilled one should expect an additional excess of linear polarization at the level $p_l \approx 50\%$ that is produced by the photon-arion transition.

Rankin (1983, see also Taylor and Stinebring 1986) has studied the profile morphology and polarization characteristics of a large number of pulsars and has come to the conclusion that pulsar radio emission arises in two distinct beams generated at different distances from the NS. The so called core-emission one which dominates at frequencies around 400 MHz is generated in the close vicinity of the NS. It exhibits circular polarization predominantly, shows no evidence of a periodic subpulse modulation and has a steeper spectrum than the emission from the hollow cone. The hollow cone emission is generated at the light-cylinder distances.

Keeping this fact in mind, one can expect a stronger effect of the photon-arion transition especially for core-cone emitted pulsars. As a result, the following limits on the symmetry-breaking scale F_a and the magnetosphere plasma density N_e can be obtained ($\lambda > 0.5$ cm):

$$\begin{gathered}
F_a \leq 10^{12}\left(\frac{B_s}{10^{12}\,\mathrm{G}}\right) \mathrm{GeV}, \\
N_e \leq 10^{7}\left(\frac{B_s}{10^{12}\,\mathrm{G}}\right)\left(\frac{10^{12}\,\mathrm{GeV}}{F_a}\right) \mathrm{cm}^{-3}.
\end{gathered} \tag{18}$$

For example, if $F_a = 5 \times 10^{10}$ GeV, the restriction on the plasma effect requires $N_e < 2 \times 10^8\,\mathrm{cm}^{-3}$.

For cone-emission pulsars we have a situation when magnetic field falls down with distance, so one should not expect a strong bound on the symmetry-breaking scale:

$$F_a \leq 10^{10}\left(\frac{B_s}{10^{12}\,\mathrm{G}}\right)\left(\frac{10^{-3}\,\mathrm{s}}{P}\right)^2 \mathrm{GeV}, \tag{19}$$

where P is the pulsar rotation period.

The basic features of pulsar polarization are known (Taylor and Stinebring 1986); orthogonal polarization modes are ubiquitous, but polarization behavior remains complicated and is not still totally understood. Therefore, we need a more detailed work to distinguish the photon-arion transitions effect from plasma and magnetosphere effects on the poarlized radiation.

3.2. *Intrinsic Emission of an Isolated NS*

A NS in vacuum should acquire an excess of polarization due to the photon-arion transition. In this case the most favourable conditions of (14)–(16) are:

$$\omega \leq 25\left(\frac{10^{10}\,\mathrm{GeV}}{F_a}\right)\left(\frac{3\times 10^{8}\,\mathrm{G}}{B_s}\right)\mathrm{eV},$$

$$N_e \leq 1.5\times 10^{16}\left(\frac{F_a}{10^{10}\,\mathrm{GeV}}\right)\left(\frac{3\times 10^{8}\,\mathrm{G}}{B_s}\right)\mathrm{cm}^{-3}, \tag{20}$$

$$L_{\mathrm{osc}} = 2\pi\frac{F_a}{B_s} = 5\times 10^{6}\left(\frac{F_a}{10^{10}\,\mathrm{GeV}}\right)\left(\frac{3\times 10^{9}\,\mathrm{G}}{B_s}\right)\mathrm{cm} \approx R_s.$$

If these conditions are fulfilled, one should expect a high degree of linear polarization ($p_l \approx 100\%$) for the intrinsic emission of the NS in the far UV spectral range ($\lambda > 500$ Å).

We believe that the best candidates are the nearest old NS and probably low mass binaries whose magnetic field is estimated to be not much greater than 10^9 G.

3.3. *Accreting NS—X-ray Pulsars*

These objects are not good candidates for the observations of the hard-energy photon-arion transition because of a high contribution of both magnetized vacuum polarization and plasma effects. But there is the only possibility for the observation in the range of the so-called "absorption vacuum features" (Pavlov and Gnedin 1984) where the plasma and vacuum polarizabilities compensate one another. These features are located near the cyclotron frequency ω_B and in the low-energy range:

$$\omega_1 = \omega_B\left(1-\frac{1}{2V}\right),$$
$$\omega_2 = \omega_B/\sqrt{V}, \tag{21}$$

where V is the so-called "vacuum" parameter defined as (Pavlov and Gnedin 1984)

$$V = \frac{1}{60\pi^2 N_e}\frac{\omega_B^4}{m_e} = \frac{3\times 10^{28}\,\mathrm{cm}^{-3}}{N_e}\left(\frac{B}{B_{\mathrm{cr}}}\right)^4. \tag{22}$$

The photon-arion transition strengthens the absorption into the "vacuum feature" due to the conversion of one orthogonal mode ($\|$) and also produces a noticeable degree of polarization in X-rays in these spectral "windows". The estimations show that only very narrow absorption features are probable. It requires the spectral resolution as:

$$\frac{\omega}{\Delta\omega} = 10^4 \, |\sin\phi| \left(\frac{F_a}{10^{12}\,\mathrm{GeV}}\right), \tag{23}$$

for the magnetic field strength $B_s = 10^{10}$ G, the longward wavelength vacuum feature being $\omega_2 = 116$ eV.

If the magnetic field strength is $B = 3 \times 10^{12}$ G, then $\omega_2 = 210$ eV and one requires the resolution:

$$\frac{\omega}{\Delta\omega} = 10^7 \, |\sin\phi| \left(\frac{F_a}{10^{12}\,\mathrm{GeV}}\right). \tag{24}$$

Unfrotunately, there are many effects that can smear this feature because of the inhomogeneity of the plasma density distribution and the electron thermal motion.

4. OPTICAL RADIATION OF MAGNETIC WHITE DWARFS (MWD)

A more advantageous situation one can meet in the case of MWDs. For an isolated MWD in pure vacuum,

$$\omega \leq \omega_{\max} = 2.5\left(\frac{10^{12}\,\mathrm{GeV}}{F_a}\right)\left(\frac{3 \times 10^7\,\mathrm{G}}{B_s}\right)\mathrm{eV}, \tag{25}$$

and one should obtain the following upper limits for F_a:

$$F_a \leq 2 \times 10^{10}\left(\frac{B_s}{3 \times 10^7\,\mathrm{G}}\right)\left(\frac{R_{\mathrm{WD}}}{10^9\,\mathrm{cm}}\right)\mathrm{GeV} \tag{26}$$

at the level $p_l \sim 100\%$ and

$$F_a \leq 2 \times 10^{11}\left(\frac{B_s}{3 \times 10^7\,\mathrm{G}}\right)\left(\frac{R_{\mathrm{WD}}}{10^9\,\mathrm{cm}}\right)\mathrm{GeV} \tag{27}$$

at the level $p_l \sim 1\%$.

If $F_a = 10^{12}$ GeV and $B_s = 3 \times 10^7$ G, then $L_{\mathrm{osc}} \simeq 50 R_{\mathrm{WD}}$ and the expected value of polarization is $p_l \simeq 0.04\%$ up to $\omega = 25$ eV.

The MWDs GrW 70° 8247 and PG 1031 + 234 present a special interest because of their especially strong magnetic field: $B_s = (150 \div 320)$MG and $B_s = (200 \div 1000)$MG, respecively (Schmidt 1988). If one adopts $B_s = 10^9$ G, one should expect

$$\begin{aligned} p_l &\sim 100\%, \qquad F_a = 6.4 \times 10^{11}\left(\frac{R_{\mathrm{WD}}}{10^9\,\mathrm{cm}}\right)\mathrm{GeV}, \\ p_l &\sim 1\%, \qquad F_a = 6.4 \times 10^{12}\left(\frac{R_{\mathrm{WD}}}{10^9\,\mathrm{cm}}\right)\mathrm{GeV}. \end{aligned} \tag{28}$$

These results are valid only for a MWD in pure vacuum and for $\omega \leq 1.2$ eV (J band) and $\omega \leq 0.12$ eV, respectively.

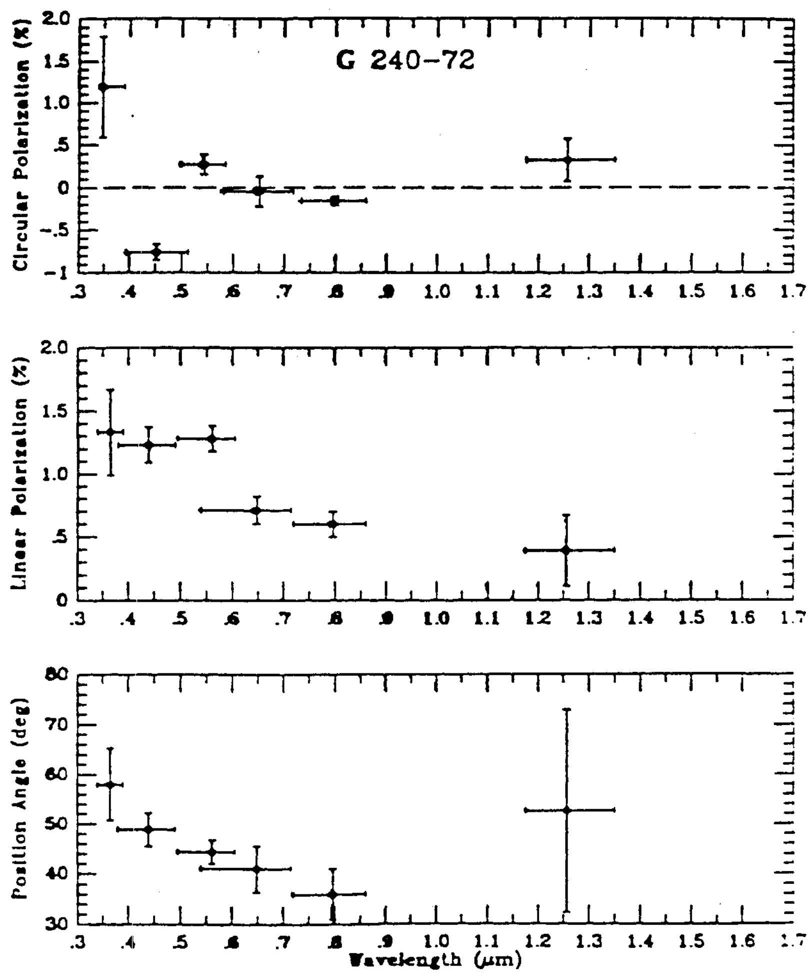

Figure 2 Optical and near-infrared polarization measurements of MWD G240-72 (from West 1989).

For the photon-arion transition ($m_a = 0$), in the wide range of frequencies $\omega \ll \omega_{\max}$ from (25) the degree polarization p_l does not depend on frequency ω. If the density of plasma in the nearest vicinity of the MWD is close to $N_e \simeq 10^{16}(F_a/10^{10}\,\mathrm{GeV})(10^9\,\mathrm{G}/B_s)\,\mathrm{cm}^{-3}$ (AM Her type objects) or we deal with the photon-axion transition ($m_a \neq 0$) one should expect an unusual, for the magnetobremsstrahlung frequency, dependence of the degree of polarization, $p_l \sim \omega^2$. One should note that this dependence qualitatively corresponds to the last observable data (Figure 2).

Finally, we would like to consider an interesting situation with spectrum of MWD VV Pup. Wickramasinghe and Visnavatan (1979) have explained the observable broad absorption features in the spectrum of this MWD due to the cyclotron absorption in high cyclotron harmonics up to 6. Pavlov *et al.* (1980)

have noticed the difficulties for this explanation because of a strong decrease of the absorption coefficients with increasing harmonic number. We would like to note that, for optically thick plasma, cyclotron absorption features are mainly produced by the ordinary mode whose polarization lies in the plane containing the magnetic field and the line of sight. But namely this mode is tested for the additional absorption via the photon-arion (axion) transition.

5. MAGNETIC Ap STARS

Generally speaking, magnetic Ap stars are also suitable candidates for the arion (axion) search. Through the magnetic field of an Ap star is not so high as that of a MWD, the size of a region of magnetic field is noticeably larger than the MWD radius. At this case the probability of the photon-arion conversion is

$$P(\gamma_{\parallel} \to a) \approx \left(\frac{B_s R_s}{2F_a}\right)^2. \tag{29}$$

The net intrinsic polarization can be estimated as

$$p_l = 0.25\left(\frac{B_s}{5\times 10^4\,\mathrm{G}}\right)^2\left(\frac{R_s}{10^{11}\,\mathrm{cm}}\right)^2\left(\frac{3\times 10^{10}\,\mathrm{GeV}}{F_a}\right)^2\%. \tag{30}$$

For the photon-axion transition,

$$p_l \approx \left(\frac{\omega R_s}{m_a}\right)^2\left(\frac{B_s}{F_1}\right)^4 \approx \left(\frac{\omega}{3\,\mathrm{eV}}\right)^2\left(\frac{R_s}{10^{11}\,\mathrm{cm}}\right)^2\left(\frac{10^{-6}\,\mathrm{eV}}{m_a}\right)^2\left(\frac{B_s}{5\times 10^4\,\mathrm{G}}\right)^4\left(\frac{10^{10}\,\mathrm{GeV}}{F_a}\right)^4. \tag{31}$$

6. INTERSTELLAR PHOTON-ARION TRANSITIONS

Now we estimate an additional polarization that can be produced via the photon-arion transition in the interstellar medium. The low value of magnetic field in the ISM is compensated by larger distances for light propagation. The oscillation lengths in the ISM are:

$$\begin{aligned} l_a &= 2(F_a/10^{10}\,\mathrm{GeV})(3\times 10^{-6}\,\mathrm{G}/B_\perp)\,\mathrm{kpc}, \\ l_p &= 2\times 10^{-4}(\omega/3\,\mathrm{eV})(1\,\mathrm{cm}^{-3}/N_e)\,\mathrm{kpc}, \\ l_m &= 1.2\times 10^{6}(\omega/3\,\mathrm{eV})(10^{-5}\,\mathrm{eV}/m_a)^2\,\mathrm{cm} \end{aligned} \tag{32}$$

This implies that one has to consider only the photon-arion transition in the ISM practically for the whole spectrum of electromagnetic waves except for probably very high energy γ-rays ($\gtrsim 10^{12}$ eV) for which the photon-axion transition has also the noticeable probability and l_m becomes comparable with l_a and l_p.

For the Galactic disk, one cannot expect any noticeable effect for visual photons because of $N_e \sim 10^{-2}\,\mathrm{cm}^{-3}$ and $l_p \ll l_a$. But the starlight from the Galactic halo should be tested for the photon-arion transition. For the Galactic halo, the electron density is estimated as $N_e \sim 10^{-3}\,\mathrm{cm}^{-3}$ (hot gas regions, Cowie and Songaila 1986). Then $l_p = 0.2(\omega/3\,\mathrm{eV})$ kpc and one should expect the amount of polarization at the level of $p_l \approx 1\%$, which is quite comparable with the observed interstellar polarization of stars produced in the Galactic disk by interstellar grains. Since the Galactic halo dimension is estimated as ~40 kpc, the expected interstellar polarization due to the photon-arion transition could be oscillatory with the characteristic length determined by Eq. (32).

Another example is the Local Bubble in the Galactic meridional plane, whose mean physical parameters are (Cox and Reynolds 1987) $N_e = 5 \times 10^{-3}\,\mathrm{cm}^{-3}$ and $B_\perp \sim B \approx (5 \pm 3)\mu$ G. As a result one has:

$$\begin{aligned} l_a &= 1.2(F_a/10^{10}\,\mathrm{GeV})(5\times 10^{-6}\,\mathrm{G}/B_\perp)\,\mathrm{kpc}, \\ l_p &= 0.04(\omega/3\,\mathrm{eV})(5\times 10^{-3}\,\mathrm{cm}^{-3}/N_e)\,\mathrm{kpc}. \end{aligned} \tag{33}$$

For the optical photons, the net polarization can be estimated as

$$p_l^{\max} = 0.1(10^{10}\,\mathrm{GeV}/F_a)^2\%. \tag{34}$$

For UV photons with $\lambda = 1000$ Å one has

$$p_l^{\max} = 1(10^{11}\,\mathrm{GeV}/F_a)^2\%.$$

7. INTERGALACTIC PHOTON-ARION TRANSITIONS

The oscillation length in the IGM can be estimated as

$$\begin{aligned} l_a &= 60(F_a/10^{11}\,\mathrm{GeV})(10^{-9}\,\mathrm{G}/B_\perp)\,\mathrm{Mpc}, \\ l_p &= 8(\omega/3\,\mathrm{eV})(0.005/\Omega_b h_{50})\,\mathrm{Mpc}. \end{aligned} \tag{35}$$

Eqs. (35) are valid for the medium between galaxy clusters. We have used the estimations of the IGM magnetic field from Fujimoto (1990).

The expected degree of polarization is at the level of $p_l \sim 0.5\%$ and should increase in the UV range as $p_l \sim \omega^2$. The polarization (and, correspondingly, the photometric brightness) of galaxies (QSOs and AGNs) should oscillate with the cosmological redshift z, the maximum polarization corresponding to the cosmological redshift

$$z \approx \pi H F_a / c B_{\mathrm{IGM}}, \tag{36}$$

where H is the Hubble constant (Figure 3).

One meets an especially interesting situation for the IGM in galaxy clusters. Kim *et al.*, (1991) have detected an excess rotation measure due to intracluster magnetic fields in clusters of galaxies. Using also the X-ray data for selected Abell clusters of galaxies they have estimated not only the magnetic field $B \approx 1\mu$ G but also the core electron number density, $N_e \approx 10^{-3}\,\mathrm{cm}^{-3}$, and the reversal scale of the magnetic field, $r_0 \approx 10$ kpc.

The case when $l_p \leq l_a \sim r_0$ exists only for $\omega \geq 150(F_a/10^{10}\,\mathrm{GeV})$ eV. It means that X-ray QSOs and AGNs projected onto a cluster core should be tested for an additional absorption due to the photon-arion transition comparable with a source located outside the core.

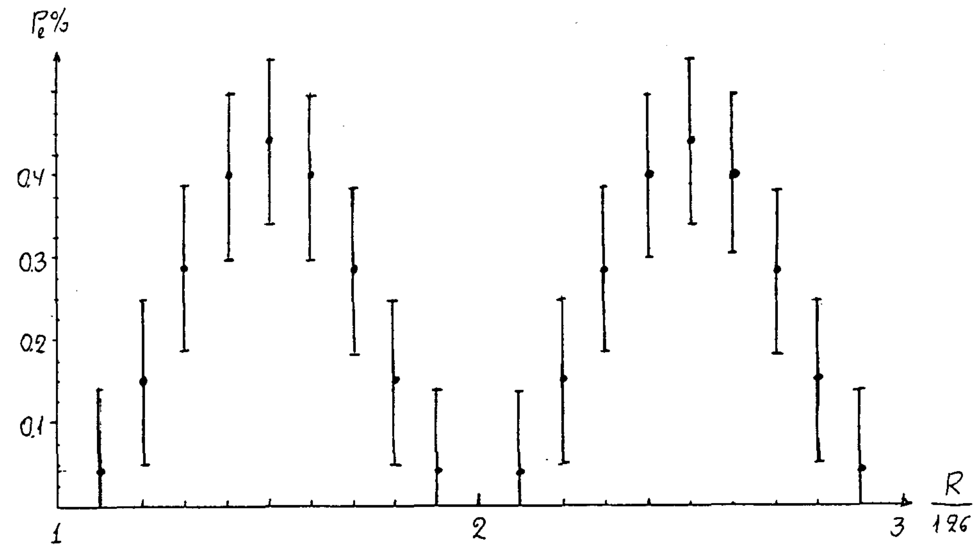

Figure 3 Oscillation of AGN's polarization with the cosmological distance (Mpc); $\Omega_6 h_{50} = 0.003$; $F_a = 10^{11}$ GeV; $B_1 = 10^{-10}$ G.

8. THE INFLUENCE OF NONIONIZED GAS

We restrict ourselves to the consideration of hot plasma regions in the interstellar and intergalactic media. The question is the contribution from nonionized material, for example, HI.

For the disk of the Galaxy, $N_e/N_H \leq 0.03$. In such physical situation one needs to consider the oscillation length l_p as a gas parameter but not plasma, and, as a rule, in many cases $l_a \gg l_p$.

Generally speaking, the transition rate is the largest when axion and photon amplitudes are coherent throughout the detection volume. For the massive axions, this coherence can be achieved though the photons acquire an effective mass m_γ when they propagate through the medium. For arions, the coherence condition can also take place due to the different signs and different wavelength dependences of gas and plasma polarizabilities, respectively.

For example, in the V band ($\lambda = 5500$ Å) free and bound electrons make comparable contributions to the dielectric constant of the ISM and produce a resonance in (3) and (4). This means that one can expect a noticeable excess of linear polarization in this band, the photon-arion transition being responsible for this excess.

An approximate estimation for the resonance width gives for the disk of the Galaxy

$$\frac{\Delta\lambda}{\lambda} \approx 10^{-2}\left(\frac{\omega}{3\,\mathrm{eV}}\right)\left(\frac{0.01\,\mathrm{cm}^{-3}}{N_e}\right)\left(\frac{B}{3\mu\,\mathrm{G}}\right)\left(\frac{10^{10}\,\mathrm{GeV}}{F_a}\right). \tag{37}$$

Of course, this situation requires a more detailed consideration.

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