Polarimetric effects associated with the detection of Goldstone bosons in stars and galaxies

Yu. N. Gnedin and S. V. Krasnikov

Central Astronomical Observatory, Russian Academy of Sciences, Pulkovo
(Submitted 10 June 1992)
Zh. Eksp. Teor. Fiz. **102**, 1729–1738 (December 1992)

Polarimetric methods are proposed for searching for light Goldstone bosons (axions and arions) in the emission from astrophysical objects. The basic idea is to make use of the conversion of photons into Goldstone bosons in the magnetic fields of compact stars and also in the interstellar and intergalactic media. This conversion is very sensitive to the polarization state of the photon, and it may lead to significant polarization of the radiation. Polarimetric observations may yield strong constraints on the Goldstone-boson–photon coupling constant. The best candidates for observation are core-emission pulsars, x-ray binaries with low-mass components, and magnetic white dwarfs. For the interstellar medium the effect may be comparable to the interstellar polarization. The radiation fluxes from quasars and from the nuclei of active galaxies may oscillate as a function of the cosmological redshift Z because of the conversion of photons into arions.

1. INTRODUCTION

A confluence of some seemingly disparate research directions—elementary particle physics, cosmology, and stellar physics—has become the center of attention in astronomical research. A crucial aspect of this confluence is the problem of "dark matter." The mass of radiating matter in the form of stars, galaxies, and the interstellar and intergalactic gas does not match the estimate of the mass in the universe derived from the gravitational interaction of matter. In addition, measurements of the behavior of the orbital velocities of stars as a function of the distance to the centers of galaxies indicate that galaxies seem to contain some "dark" or "hidden" mass. The same conclusion is reached by astrophysicists who have measured the orbital velocities of galaxies which are part of clusters.

What is the nature of this dark matter? If we assume that it consists of baryons, then constraints on its magnitude follow from primordial nucleosynthesis.[1]

Other possible candidates have been advanced for this dark matter: neutron stars, black holes, and also some faintly radiating celestial objects like one of the major planets of the solar system (Jupiter), which would fill a spherical halo around many galaxies (including our own). Such hypotheses, however, run into some serious difficulties from the astronomical standpoint. We will not go into those difficulties here.

In this situation, the idea that the dark matter consists of weakly interacting fundamental elementary particles has become extremely popular. Recent advances in elementary particle physics have significantly expanded the spectrum of such candidates.

The axion has recently been one of the most popular particles. This is a pseudoscalar Higgs boson of exceedingly small mass, which interacts weakly with other matter (a "dark axion"). Elementary particle theory itself tells us next to nothing about possible values of its mass (m_a) or about the coupling constants of an axion with fermions and photons (g_a,g). It is known, however, that they are coupled. For example, we have[1,2] $m_a = 10^{-19}\xi g$, where ξ is a factor on the order of unity which depends on the particular model.

Another particle, very similar to an axion, is the arion.[3] Phenomenologically, its only distinction from the axion is that its mass is strictly zero.

The strongest constraints on m_a, g_{af}, and g have emerged from astrophysical considerations. Specifically, after axions (or arions) are produced in the interior of stars (as a result of processes like the Compton effect, bremsstrahlung etc.) they escape freely from the star, carrying off some energy and thereby influencing stellar evolution. Analysis of these considerations yields[1]

$$g < 10^{-10}\ \mathrm{GeV}^{-1}. \tag{1}$$

An important aspect of an axion (or arion) is the nature of its interaction with phonons (Fig. 1).

The existence of such a vertex has the consequence that an axion (or arion) in a external magnetic field **B** is capable of converting into a photon (polarized parallel to **B**), and vice versa. The energy of the photon which is produced is equal to the total energy of the axion. The inverse process is also possible: a photon can convert into an axion (or arion) in a external magnetic field; this photon–axion (or –arion) conversion is an oscillatory process.

In this paper we discuss, and carry out some related calculations on, some effects associated with the conversion of photons into axions and arions in astrophysical objects with magnetic fields (neutron stars, magnetic white dwarfs, and the interstellar medium). The photon ↔ axion conversion acts as an additional absorption process, one which is

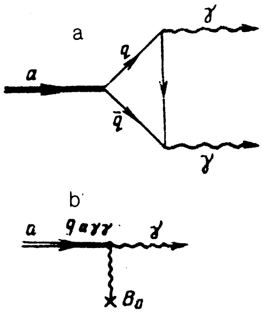

FIG. 1. Interactions of photons with axions (or arions). a—Two-photon decay of an axion (or arion); b—the axion (arion)↔photon conversion.

very sensitive to the polarization state of the photon, since only a single polarization state is subject to the conversion.[1,4] It is thus possible to observe some additional photometric and (especially) polarimetric effects resulting from these processes in radiation from stars which is propagating across magnetic field lines. The process with which we are primarily concerned here is the conversion of a photon into an arion.

2. THEORY OF THE PHOTON ↔ AXION CONVERSION IN A MAGNETIC FIELD

We consider plane axion waves and plane electromagnetic waves which are propagating through a medium with a refractive index n ($\varepsilon \equiv 1 - n \ll 1$), along the z axis, which is perpendicular to the uniform external magnetic field $\mathbf{B}$. We assume that these waves are of the form

$$\begin{pmatrix} A_\parallel \\ \phi \end{pmatrix} \exp\{i(\omega t - kx)\}. \tag{2}$$

Here $A_\parallel$ is the amplitude of the wave potential, which is polarized parallel to $\mathbf{B}$ (a wave polarized perpendicular to $\mathbf{B}$ does not interact with axions). The equations of motion are

$$\begin{pmatrix} k^2-\omega^2 n^2 & iBg\omega \\ -iBg\omega & k^2-\omega^2+m_a^2 \end{pmatrix}\begin{pmatrix} A_\parallel \\ \phi \end{pmatrix}=0. \tag{3}$$

We are using the Lorentz–Heaviside system of units, with $\hbar = c = 1$, and a fine-structure constant $\alpha = e^2/4\pi = 1/137$. The condition for the compatibility of these equations is

$$k^2=(k^\pm)^2=\omega^2(1-m_a^2/\omega^2-\bar{\varepsilon}\pm[\bar{\varepsilon}^2+g^2B^2/\omega^2]^{1/2}), \tag{4}$$

where

$$\bar{\varepsilon}=\varepsilon-m_a^2/2\omega^2,$$

and the solutions of these equations are

$$u^\pm=\begin{pmatrix} \dfrac{(k^\pm)^2-\omega^2+m_a^2}{igB\omega} \\ 1 \end{pmatrix}. \tag{5}$$

If we have a "pure photon" at the origin of coordinates, i.e., if we have $A_\parallel(0) = 1$, $\phi(0) = 0$, then two mixed axion-electromagnetic waves, each with its own amplitude and its own velocity, will subsequently propagate. The resultant potential at an arbitrary point z is given by

$$u(z)=\frac{u^+\exp(-ik^+z)-u^-\exp(-ik^-z)}{(k^+-k^-)/iBg}.$$

The axion component at this point is

$$\phi(z)=-\frac{\exp\left\{-i\dfrac{k^++k^-}{2}z\right\}\sin\left\{\dfrac{k^+-k^-}{2}z\right\}}{(1+x^2)^{1/2}}, \tag{6}$$

where

$$x=\bar{\varepsilon}\omega/Bg.$$

Correspondingly, the electromagnetic component decreases:

$$|A_\parallel|=1-|\phi|^2=1-\frac{1}{1+x^2}\sin^2\left[\frac{1}{2}(x^2+1)^{1/2}Bgz\right]$$
$$=1-A_{osc}^2\sin^2\Lambda. \tag{7}$$

The argument of the sine is written under the assumption $m_a \ll \omega$. Expressions like (7) were derived in Refs. 3 and 4. We can expect effects due to the axion ↔ photon conversion in the case $A_{osc} \sim 1$, i.e., in the case $|x| < 1$. Assuming that the medium is a plasma with an electron density N_e and a neutral-gas (hydrogen) density N_H, we find

$$\varepsilon=\frac{2\pi\alpha N_e}{m_e\omega^2}-4\pi\beta N_H-\frac{14\alpha^2}{45m_e^4}B^2-\frac{m_a^2}{2\omega^2}. \tag{8}$$

The first term in (8) describes the contribution of the plasma polarizability; the second describes the contribution of the gas; the third describes the contribution of the so called polarizability of vacuum in a magnetic field;[5,6] and the last term is the contribution of the axion field. For arions we would have $m_a = 0$.

In the system of units we are using here, the unit for the magnetic field B is $(\text{eV})^2$, and the unit of length is $(\text{eV})^{-1}$. One gauss corresponds to 6.9×10^{-2} $(\text{eV})^2$, and 1 cm to 5×10^4 $(\text{eV})^{-1}$.

The various terms in (8) have different signs, so there is the possibility that they would cancel out completely and that expression (8) would be equal to zero. This important case is discussed in more detail below.

We thus have

$$x=\frac{a}{\omega}-b\omega,\quad a=\frac{10^{-7}N_e}{Bg}-\frac{m_a^2}{2Bg},$$
$$b=2{,}4\cdot10^{-28}\frac{B}{g}+5\cdot10^{-10}\frac{N_H}{Bg}. \tag{9}$$

The requirement $|x| < 1$ is equivalent to the condition

$$\omega\in(\omega^*-\delta,\ \omega^*+\delta), \tag{10}$$

where

$$\omega^*=(\omega_0^2+\delta^2)^{1/2},\quad \delta=1/2b,\quad \omega_0=(a/b)^{1/2}. \tag{11}$$

Here ω_0 is that frequency at which expression (8) becomes $\bar{\varepsilon}(\omega_0) = 0$.

We can estimate the width of this resonant interval in various particular cases.

1. We consider the case in which the plasma polarizability is offset by the contribution of the axion ($m_a \neq 0$) field. We assume

$N_e = (m_a^2/2\cdot10^{-7}) + \Delta N_e$ and $N_H = 0$.

Then we have

$$\frac{\delta}{\omega^*}<\frac{\delta}{\omega_0}=\frac{1}{2(ab)^{1/2}}\leqslant4{,}5\cdot10^{13}\frac{g}{m_a}\left(\frac{N_e}{\Delta N_e}\right)^{1/2}. \tag{12}$$

If we follow the standard model of the axion[1] and estimate the ratio g/m_a as $g/m_a \sim 10^{-19}$, we find

$$\frac{\delta}{\omega^*}\left(\frac{\Delta N_e}{N_e}\right)^{1/2}\leqslant5\cdot10^{-6}. \tag{13}$$

2. We consider the case of a resonance in which the polarizability in (8) is completely cancelled by the polarizability of the magnetized vacuum. In an optically thick plasma with a strong magnetic field, a "vacuum" absorption line can form in this case (Refs. 6–8, for example). This situation is characteristic of accreting neutron stars with strong magnetic fields $B = 10^{12}$–10^{13} G and of accreting magnetic white dwarfs (objects of the AM Her type).[6,9]

The amplitude (7) for photon–axion oscillations can be rewritten in this case as

$$A_{osc}^2=\frac{\gamma^2}{(\omega-\omega_r)^2+\gamma^2}, \tag{14}$$

where

$$\omega_r^2=\frac{45}{28}\frac{m_e^4}{\alpha^2B^2}(\omega_p^2-m_a^2), \qquad \gamma=1,6\frac{m_e^4g}{\alpha^2B} \tag{15}$$

Using (14) and (15), we can easily estimate the width of the resonance for arions under the conditions prevailing at an accreting neutron star with a strong magnetic field:

$$\frac{\delta}{\omega_0}=\frac{3\cdot10^{-6}}{\omega_0}\left(\frac{g}{10^{-13}\,\mathrm{GeV}^{-1}}\right)\frac{10^{12}\,\mathrm{G}}{B}. \tag{16}$$

With $g=10^{-10}\,\mathrm{GeV}^{-1}$ and $N_e=10^{19}\,\mathrm{cm}^{-3}$, the width of the resonance is $\delta/\omega_0=3\times10^{-5}$. This width would be completely within the capabilities of future instruments (we have in mind the Spektr–Roentgen-Gamma project). A variation of the density with distance in the accreting region of the neutron star would of course smear out the resonance.

3. As the radiation propagates through the interstellar medium, a resonance may also occur because the plasma polarizability, determined by N_e and by the neutral gas N_H, appears in the expression for the dielectric constant (8), with different signs.
The resonant frequency is

$$\omega_0=10(N_e/N_H)^{1/2}\,\mathrm{eV}. \tag{17}$$

In the case of the axions described by the standard model,[1] it would hardly be possible to observe the photon $\leftrightarrow$ axion conversion under real astrophysical conditions. The requirement $|x|<1$ imposes definite conditions on the nature of the density variation. For standard axions, photons of a certain frequency convert into axions only on the part of the path on which N_e differs from the equilibrium density by an amount ΔN_e:

$$\frac{\Delta N_e}{N_e}=3\cdot10^{13}\frac{g}{m_a}. \tag{18}$$

For standard axions we would have $g/m_a\sim10^{-19}$ and $(\Delta N_e/N_e)\lesssim3\times10^{-6}$. We thus restrict the discussion below to the case of arions. For them we have $\delta/\omega_0=10^{17}(g/\sqrt{N})$, i.e., $\delta\gg\omega_0$ under the conditions

$$g\gtrsim10^{-13}\,\mathrm{GeV}^{-1},\quad N\ll10^{-10}\,\mathrm{eV}^3. \tag{19}$$

We thus have $|x|<1$ if

$$\omega\in\left(\frac{\omega_0^2}{2\delta},2\delta\right)=\left(a,\frac{1}{b}\right)=\left(\frac{10^{-7}N}{Bg},\frac{2\cdot10^{27}g}{B}\right). \tag{20a}$$

A second condition imposes a limitation on the phase of the sine in (7):

$$\Lambda\approx\frac{Bgz}{2}\geqslant C, \tag{20b}$$

where C is a measure of the sensitivity of the proposed experiment (with $C=0.1$, for example, the light intensity $|A_\parallel|^2$ decreases by $\sim1\%$ after passage through the magnetic field).

As a rule, the magnetic fields of stars are dipole fields. Let us assume that the magnetic field $B(z)$ varies significantly over distances $\sim z_B$. For stellar fields we would typically have

$$\omega z_B\gg1,\quad z_B\leqslant1/gB.$$

Expression (7) does not hold in this case. To find an estimate of the form (20), we assume that the arion $\leftrightarrow$ photon conversion occurs only on the interval $z_0<z<z$, where $|x(z)|<1$, and on this interval we ignore x in comparison with 1. Then for $z>z_1$ we have

$$|A_\parallel|^2=1-\sin^2\frac{1}{2}g\int_{z_0}^{z_1}B(y)dy. \tag{21}$$

Let us consider the simple case in which the radiation is propagating radially in the equatorial plane of a magnetic dipole and N is independent of z. Requiring that the phase of the sine in (11) be greater than C and that ω satisfy (20a) in which we now have $B=B(z)$, we find the conditions for observing oscillations:

$$\Lambda/C=gB_0R_0/4C>1, \tag{22a}$$

$$\omega\in\left(\min\left\{\frac{1}{b(R_0)},\frac{a(R_0)}{(1-C/\Lambda)^{3/2}}\right\},\left(\frac{\Lambda}{C}\right)^{3/4}\frac{1}{b(R_0)}\right), \tag{22b}$$

where R_0 is the radius of the star, and $B_0=B(R_0)$.

The photon $\leftrightarrow$ axion (arion) conversion process in a magnetic field may cause birefringence effects. Specifically, ellipticity may arise when linearly polarized light propagates, and there may be oscillations in the plane of the position angle of the polarization vector (the analog of the Cotton–Mouton effect).[4] We know that the emission from plasma in a magnetic field is characterized by dichroism and birefringence. As light passes through vacuum in a strong magnetic field birefringence effects arise again, in this case because of QED polarization of virtual electron-positron pairs.[5,6,10] The change in the parallel polarization mode as a result of the photon $\leftrightarrow$ axion conversion also leads to birefringence effects, which follow directly from (3)–(5).

Setting $m_a=0$ in (4) and setting $x^2\ll1$, we find, under the condition $\omega\gg Bg$,

$$\frac{k^++k^-}{2}\approx\omega\left(1-\frac{\varepsilon}{2}+\frac{g^2B^2}{4\omega^2}\right). \tag{23}$$

The additional phase difference between the waves $A_\parallel$ and $A_\perp$ at the point z due to the contribution of the arion field is

$$\Delta(\Delta\varphi)=(g^2B^2/4\omega+{}^1/_2\varepsilon\omega)z. \tag{24}$$

After light polarized at an angle α has passed through the magnetic-field region (we are ignoring the ellipticity), it is polarized at an angle α', where

$$\frac{\mathrm{tg}^2\,\alpha'}{\mathrm{tg}^2\,\alpha}=1-A_{osc}^2\sin^2\Lambda. \tag{25a}$$

In general, the light acquires a circular polarization with a degree

$$p_v=\sin2\alpha'\sin(\Delta(\Delta\varphi)). \tag{25b}$$

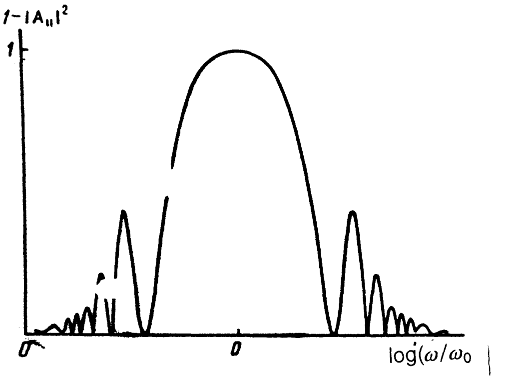

FIG. 2. Probability for the photon↔arion conversion as a function of the frequency.

3. EXPECTED OBSERVATIONAL EFFECTS

In principle, we can expect the following effects.

1) After passage through the magnetic-field region, the intensity of the radiation from the star decreases (except when this radiation is polarized strictly perpendicular to the direction of the magnetic field **B**). The photon ↔ arion conversion process thus acts as an additional cause of photon absorption under astrophysical conditions.

2) As the result of the conversion process, initially unpolarized light from a star acquires a partial linear polarization P_l. At small values of Λ the degree of polarization satisfies $P_l \propto \Lambda^2$. The photon ↔ arion conversion thus acts as an additional mechanism for the appearance of polarized light under astrophysical conditions.

3) The conversion leads to rotation of the polarization plane of light propagating at an angle from the magnetic field. If the polarization initial makes an angle of 45° with **B**, it is rotated through an angle $\theta \sim \Lambda^2/2$ after it emerges from the magnetic-field region (at small values of Λ).

4) As light from stars and other astrophysical sources propagates across the magnetic field, it acquires ellipticity as a result of the photon–arion conversion. The extent of this ellipticity is given by (24) and (25).

5) In principle, light could also be completely depolarized. If the phase difference between the amplitudes $A_\parallel$ and $A_\perp$ is sufficiently large, these waves lose their coherence. Light which is (for example) initially polarized at an angle of 45° from **B** thus emerges completely unpolarized from the magnetic-field region. The contribution of arions to this process is determined by the value of $\Delta(\Delta\varphi)$.

There are numerous mechanisms which would lead to similar effects. The interpretation of observations can be simplified by the very specific dependence of these quantities (P_l,θ,...) on the frequency near the boundaries of interval (20a) at large values of Λ. We get an idea of the situation from Fig. 2, which is a plot of $1-|A_\parallel|^2$ versus $\log(\omega/\omega_0)$ for the case $\Lambda=5$, $\delta/\omega_0=5$.

4. SOME ASTROPHYSICAL APPLICATIONS

The probability for a photon-arion conversion is determined by the product of the magnetic field B and the path length l which the electromagnetic wave traverses in this field. The most suitable candidates for observing effects due to this transition are therefore compact magnetic stars (white dwarfs and neutron stars). Let us look at some examples.

a) Rapidly rotating neutron stars (pulsars). For these objects, conditions (20) become

$$\frac{\Lambda}{C}=\frac{g}{10^{-12}\,\text{GeV}^{-1}}\,\frac{B_0}{10^{12}\,\text{G}}\,\frac{R_0}{10^6\,\text{cm}}\gtrsim 1,$$
$$\omega\lesssim 3\cdot 10^{-5}\left(\frac{\Lambda}{C}\right)^{1/2}\frac{g}{10^{-12}\,\text{GeV}^{-1}}\,\frac{10^{12}\,\text{G}}{B_0}. \qquad (26)$$

Under conditions (26) we would expect a significant additional polarization due to the photon–arion-transition mechanism.

Rankin[11] (see also Ref. 12) studied the profiles and polarization characteristics of the light from a significant number of pulsars, reaching the conclusion that the pulsating radio emission might be generated in two regions, which differ in their depth below the surface of the neutron star. For "core-emission" pulsars, the radiation is generated in the immediate vicinity of the surface. At frequencies ~400 MHz, this radiation predominates, and it is for this radiation that we would expect the most significant manifestation of effects of the photon–arion conversion.

Setting $B_0=10^{13}$ G and $R_0=10^6$ cm, we find that $\omega=4\times10^8\cdot 2\pi$ (s^{-1}) falls in the interval (22b), and in this case we can expect to observe effects due to arion-photon oscillations ($C=1$) over the range to $g=10^{-13}$ GeV^{-1}. Weak effects ($C=0.1$), on the other hand, can be expected even at $g=10^{-14}$ GeV^{-1}.

We also note that under these conditions we have

$$\Delta(\Delta\varphi)=5\cdot10^4\omega+\frac{5\cdot10^{-5}}{\omega}\left(\frac{g}{10^{-10}\,\text{GeV}^{-1}}\right)^2$$

and interactions with arions can have a strong effect on the ellipticity of the radiation.

b) Magnetic white dwarfs. Conditions (20) for an isolated white dwarf take the form

$$\Lambda=0{,}3\,\frac{g}{10^{-11}\,\text{GeV}^{-1}}\,\frac{B_0}{3\cdot10^7\,\text{G}}\,\frac{R_{WD}}{10^9\,\text{cm}}$$
$$\omega\lesssim\omega_{max}=10^{-1}\left(\frac{\Lambda}{C}\right)^{1/2}\frac{g}{10^{-11}\,\text{GeV}^{-1}}\,\frac{3\cdot10^7\,\text{G}}{B_0}\,\text{eV}. \qquad (27)$$

For g we thus find the estimate

$$g\lesssim 3\cdot10^{-11}\left(\frac{B_0}{3\cdot10^7\,\text{G}}\right)\left(\frac{R_{WD}}{10^9\,\text{cm}}\right)\text{GeV}^{-1} \qquad (28)$$

for $C\sim1$.

The magnetic white dwarfs GrW 70°8247 and PG 1031 + 234 are especially interesting because of their exceedingly strong magnetic fields ($B=150$–320 and 200–1000 MG, respectively[13]). Assuming $B=10^9$ G, we find

$$\Lambda=0{,}1 \quad\text{for}\quad g=10^{-13}\frac{10^9\,\text{cm}}{R_{WD}}\,\text{GeV}^{-1};\qquad \omega\lesssim3\cdot10^{-2}\,\text{eV},$$
$$\Lambda=1 \quad\text{for}\quad g=10^{-12}\frac{10^9\,\text{cm}}{R_{WD}}\,\text{GeV}^{-1};\qquad \omega\lesssim0{,}3\,\text{eV}. \qquad (29)$$

It is particularly worth noting that in the latter case we have $\Delta(\Delta\varphi)\sim1.5$ (with $\omega\sim0.3$ eV). In other words, an interaction with arions is capable of converting linearized polarized radiation into circularly polarized radiation, and vice versa.

c) Photon-arion conversion in the interstellar medium. The weakness of the magnetic field here may be offset by the huge distances. Assuming $B\sim3\times10^{-6}$ G, we find $\Lambda\sim1.5(g/10^{-10}\,\mathrm{GeV}^{-1})(z/1\,\mathrm{kpc})$. The polarization properties of the emission from stars should thus vary periodically with distance. The period of these variations should be smaller than the dimensions of the galactic halo (~40 kpc) at $g\geqslant10^{-11}$ GeV^{-1}. Because of the neutral hydrogen [see (14)], however, this effect will be manifested in only narrow frequency range near

$$\omega_0\approx0{,}3\left(\frac{N_e}{10^{-3}\,\mathrm{cm}^{-3}}\Big/\frac{N_H}{1\,\mathrm{cm}^{-3}}\right)^{1/2}. \qquad (30)$$

Interestingly, ω_0 depends on the parameters of the medium but not on g. For disk stars, with $\omega_0\approx2$ eV, the width of the interval is ($A_{\mathrm{osc}}\sim1$)

$$\frac{\delta}{\omega_0}\sim6\cdot10^{-3}\frac{B}{3\cdot10^{-6}\,\mathrm{G}}\,\frac{g}{10^{-10}\,\mathrm{GeV}^{-1}}\Big/\left(\frac{N_e}{10^{-3}\,\mathrm{cm}^{-3}}\,\frac{N_H}{1\,\mathrm{cm}^{-3}}\right)^{1/2}. \qquad (31)$$

d) Photon–arion conversions in the intergalactic medium. The magnetic field and the electron density of the intergalactic medium can be estimated as $B\sim10^{-9}$ G and $N_e\sim10^{-7}$ cm^{-3}. We then have

$$\Lambda\approx0{,}5\,\frac{g}{10^{-10}\,\mathrm{GeV}^{-1}}\,\frac{z}{1\,\mathrm{Mpc}}\,\frac{B}{10^{-9}\,\mathrm{G}},$$
$$\omega>14\,\frac{10^{-10}\,\mathrm{GeV}^{-1}}{g}\,\frac{N_e}{10^{-7}\,\mathrm{cm}^{-3}} \qquad (32)$$
$$\Delta(\Delta\varphi)\big|_{\omega=\omega_{min}}=50\,\frac{10^{-10}\,\mathrm{GeV}^{-1}}{g}\,\frac{z}{1\,\mathrm{Mpc}}.$$

The polarization (and thus the photometric brightness) of galaxies will oscillate as a function of the cosmological redshift Z_c. The polarization will reach a maximum at a cosmological redshift

$$Z_c\sim4\pi H/gB,$$

where H is the Hubble constant.

[1]G. Raffelt, Phys. Rep. **198**, 1 (1990).
[2]M. Turner, Phys. Rep. **197**, 67 (1990).
[3]A. A. Anselm, Phys. Rev. D **37**, 2001 (1988).
[4]G. Raffelt and L. Stodolsky, Phys. Rev. D **37**, 1237 (1988).
[5]S. L. Adler, Ann. Phys. **67**, 599 (1971).
[6]G. G. Pavlov and Yu. N. Gnedin, Sov. Sci. Rev. E **3**, 197 (1984).
[7]G. G. Pavlov and Yu. A. Shibanov, Zh. Eksp. Teor. Fiz. **76**, 1457 (1979) [Sov. Phys. JETP **49**, 741 (1979)].
[8]G. W. Ventura, W. Nagel, and P. Mezaros, Astrophys. J. Lett. **233**, L125 (1979).
[9]G. G. Pavlov, I. G. Mitrofanov, and Yu. A. Shibanov, Astrophys. Space Sci. **73**, 63 (1980).
[10]P. Meszaros and G. Ventura, Phys. Rev. Lett. **41**, 1544 (1978).
[11]G. M. Rankin, Astrophys. J. **274**, 333, 359 (1983).
[12]G. H. Taylor and D. R. Stinebring, Ann. Rev. Astron. Astrophys. **24**, 285 (1986).
[13]G. D. Schmidt, IAU Col. **95**, 377 (1987).

Translated by D. Parsons

 0038-5646/92/120937-12$05.00