New method of galactic axion search

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ABSTRACT

An important and appealing candidate of the galactic dark matter is the axion, which was postulated to solve the CP (Charge-conjugation Parity) violation problem in strong interaction of the standard particle theory. A new experimental method is proposed to determine both the axion mass and its velocity distribution on Earth. The method uses collectively and coherently excited atoms or molecules triggered by strong laser field, resulting in galactic axion absorption along with signal photon emission to be detected.

Keywords Cold dark matter, Axion, Strong CP problem, Physics beyond standard model, Superradiance
Introduction  If the QCD (Quantum Chromo-Dynamics) axion \[^1\], \[^2\], is the solution to strong CP problem \[^3\] in the standard particle theory, the axion may become the most dominant component of our galaxy \[^4\], both in terms of its mass density \(\rho_G \simeq (0.3 \sim 0.45) \text{GeV}/c^2 \text{cm}^{-3}\) and by far the largest number density \(n_a \sim 10^{13} \text{cm}^{-3}\) (its precise value depending on the Peceei-Quinn (PQ) symmetry breaking scale \(f_a\)). The current parameter region explored in major experiments is in the range, \(f_a = 10^8 \sim 10^{13} \text{GeV}\), the corresponding axion mass in the range, \(m_a = 10^2 \sim 10^{-3} \text{meV}\) (inversely proportional to \(f_a\)).

Origin of the galactic axion is traced back to the QCD epoch of cosmic temperature \(100 \text{MeV}/k_B \sim 10^{12} \text{K}\), since the axion mass is generated by the QCD chiral symmetry breaking. Its couplings to ordinary quarks, leptons, and gauge bosons are all suppressed by the PQ symmetry breaking scale \(\propto 1/f_a\), hence axions are born in cold, namely their initial velocities at creation are much less than the velocity of light. Cold axions decouple from radiation and matter throughout the later epoch of cosmological evolution \[^4\]. After the radiation energy density drops below the axion matter density, gravitational self-interaction begins to dominate and axions become virialized (process of violent relaxation \[^5\]), leaving a velocity dispersion of order \(10^{-3} \times \text{the light velocity}\).

The symmetry breaking scale may be raised if it is higher than the reheat temperature due to inflation \[^6\]. This widens the parameter space to be explored up to \(\sim 10^{16} \text{GeV}\).

Current experiments \[^7\], \[^8\], \[^9\], of galactic axion search sensitive to low mass axions use magnetic axion conversion into microwave in cavity \[^10\]. In the present work we propose an entirely new experimental method of axion detection using atoms or molecules, which is sensitive both to the axion mass and to its velocity distribution at the location of earth. If the feasibility of the method is demonstrated, it provides a powerful tool to directly link physics beyond the standard theory and the dark matter cosmology/astrophysics.

We use the natural unit of \(\hbar = c = 1\) unless otherwise stated.

Microscopic process of galactic axion search  We propose a detection method using atomic or molecular process, \(a_G + |i\rangle \rightarrow |f\rangle + \gamma_t + \gamma_s\): a galactic axion \((a_G)\) collides against an atom or a molecule in an excited state \(|i\rangle\), which then de-excites to a lower energy state \(|f\rangle\), emitting a signal photon \(\gamma_s\) under a strong trigger electromagnetic (EM) field \((\gamma_t)\). The process is called Triggered Radiative Absorption of Cosmic Axion (TRACA for short). Relevant diagrams are depicted in Fig(1). The virtual photon of four momentum \(k\) may be off or on the mass shell, depending on its invariant mass squared \(k^2 = k_0^2 - \vec{k}^2 \neq 0\) or 0. The two cases have different experimental merits. In the on-shell case one may view TRACA occurring in two real steps: the triggered decay axion produces a real photon, which is inelastically scattered off excited atom/molecule to generate a signal photon \(\gamma_s\). The on-shell TRACA may occur in the trigger photon frequency \((\omega_t)\) range at \(m_a/2\), hence in the rf (radio frequency) or microwave region, while the trigger frequency of the off-shell TRACA is in the optical to the infrared (IR) range depending on the level spacing \(\epsilon_{if}\).

We also considered TRACA arising from diagrams containing direct electron coupling with the axion. In most axion parameter region we found that rates of axion emission are smaller than those of Fig(1) even for DFSZ model \[^2\] of large axion-electron coupling.

The method can be applied to axion-like particles \[^12\] as well.
In laboratory experiments on Earth, axions collide with a velocity vector \( \vec{\beta}_a = -\vec{\beta}_e \) where \( \vec{\beta}_e \) is the earth velocity vector in the Galaxy. Its flux is, when one fine-tunes the detecting system to the direction \( -\vec{\beta}_e \), of order, \( 10^{29} \text{cm}^{-2}\text{sec}^{-1}(10\mu\text{eV}/m_a) \) by taking the normalized Maxwellian distribution \( F_M(\vec{\beta}_a) \) of dispersion \( \beta_0 \simeq 0.8 \times 10^{-3} \), which may well be an excellent approximation except for a possible BEC component \[11\], separately detectable. The uniform \( -\vec{\beta}_e \) wind is superimposed by diurnal and sidereal modulation, which can be detected, too. One of our objectives of the proposed experiment is to precisely determine these important parameters. The required pointing accuracy should be better than of order \( O(\beta_0) \) radians, which is more quantified later.

To maximize TRACA rates, we fully utilize effects of large axion occupation and coherence \( \rho_{if} \) between \(|i\rangle \) and \(|f\rangle \): the total TRACA rate for entire atoms is proportional to the product of three factors, (1) axion number density \( n_a \), giving the enhancement \( 3 \times 10^{13} m_a/10\mu\text{eV} \) compared with the single axion, (2) field power \( |E_t|^2 = \omega t n_t \) (its photon number density denoted by \( n_t \)) driving the axion decay, (3) macro-coherently prepared atomic/molecular target number squared \( N^2 T \rho_{if}^2 \). The principle of amplification by macro-coherence \[13\] related to the second and the third issue has been experimentally demonstrated in \[14\] for a weak QED (Quantum ElectroDynamics) process of two-photon emission in para-hydrogen vibrational transition (the quantum number change of \( v = 1 \rightarrow 0 \)). The process was amplified by a factor \( \sim 10^{18} \) in rate, with a realization of atomic coherence \( \rho_{if} \sim 0.03 \) \[15\]. The process was termed PSR (Paired Super-Radiance) in an obvious analogy to Dicke’s super-radiance \[16\].

![TRACA diagram](image.png)

**Figure 1:** TRACA diagram of axion detection. Virtual or real photon emitted by the axion \( a_G \) decay triggered by \( \gamma_t \) induces inelastic photon or Raman scattering producing a signal photon \( \gamma_s \). Diagram of exchanged atomic/molecular vertexes (s-channel \(|n\rangle \) here being replaced by u-channel \(|n\rangle \) not depicted here) adds coherently to this contribution. This pair diagram is further doubled by \( \gamma_t \leftrightarrow \gamma_s \) exchange.

**Derivation of rate formula** We shall first treat the off mass shell case and later discuss photon emission on the mass shell. Suppose that the intermediate atomic state \(|n\rangle \) related to electric dipole transition is far away in energy from the initial and the final states \(|i\rangle, |f\rangle \); \( \epsilon_{ni}, \epsilon_{nf} \gg \epsilon_{if} \). This holds for the example of p-H\(_2\). We further assume the frequency relation, \( \omega_t, \omega_s \) (signal frequency) \( \gg m_a \) for the off-shell TRACA.

The axion-photon-photon vertex in Fig(1) has a form \( g_{a\gamma\gamma} a \vec{E} \cdot \vec{B} \) in terms of field operators, since the axion field \( a \) is a pseudo-scalar. Its coupling constant is suppressed by \( g_{a\gamma\gamma} \sim \alpha/(\pi f_a), \alpha = 1/137 \). The dominant atomic coupling in the internal vertex is of the electric dipole form \( \vec{d} \cdot \vec{E}' \), giving rise to an external field coupling of the quantum field theoretic form, \( \vec{B}' \cdot \langle 0|T(\vec{E}'\vec{d} \cdot \vec{E})|0 \rangle \) with propagation. The central part
of the probability amplitude is thus

$$\sqrt{\frac{\rho_G}{2\pi m_a f_a}} \left( \frac{d_{nf} \cdot \vec{E}_s d_{ni} \cdot \vec{B}_t (k_t - q)^2}{\epsilon_{ni}} + (s \leftrightarrow t) \right) \times 2,$$

(1)
to a good approximation. Here $d_{nf}, d_{ni}$ are dipole operators related to the decay rates $\gamma_{nf}, \gamma_{ni}$. The propagator factor is approximately $(k_t - q)^2/(k_t - q)^2 \sim -\omega_t/(2m_a)$.

For simplicity we take the atomic transition of spin parity change, $0^+ \rightarrow 0^+$. Using the probability amplitude given above, the off-shell TRACA rate is calculated as

$$\frac{d\Gamma_1}{d\Omega_s} = \frac{9}{4\pi^2} \rho_G N_f^2 \rho_{1f}^2 \left( \frac{c_{\alpha\gamma\gamma} \alpha}{m_a f_a} \right)^2 \frac{\gamma_{ni} \gamma_{nf} \epsilon_{1f}^2}{\epsilon_{ni}^3 \epsilon_{nf}^3} \omega_n^3 \omega_t n_t \sin^2 \theta_{ts} F_M \left( \frac{\vec{p}_{1f} - \vec{k}_t - \vec{k}_s}{m_a} \right),$$

(2)

with $|\vec{k}_s| = \omega = \epsilon_{1f} - \omega_t$ and $\vec{p}_{1f}$ the phase imprinted at the excitation [19]. Various level spacings, $\epsilon_{ab} = \epsilon_a - \epsilon_b$, are used together with the radiative decay rate $\gamma_{ab}$ for $a \rightarrow b$. We assumed that only one dominant upper level $|n\rangle$ contributes, but extension to include many relevant levels is straightforward. Model dependent factor $c_{\alpha\gamma\gamma}$ [17], [18] is typically -0.97 for KSVZ-model [1], and 0.36 for DFSZ-model [2].

Some striking features of the off-shell TRACA formula are as follows. The angular dependence $\sin^2 \theta_{ts}$ with the relative polarization angle $\theta_{ts}$ between $\gamma_n$ and $\gamma_s$ indicates a strong correlation of two linear polarizations: the rate is maximal at the perpendicular $\theta_{ts} = \pi/2$ and vanishes at the parallel configuration. There is a one-to-one correspondence of angular distribution to the axion velocity distribution $F_M(\vec{\beta}_a)$ to the differential rate $d\Gamma_1/d\Omega_s$.

The on-shell TRACA rate is given by $n_a \Gamma_a T \sigma_{scat} N_f^2 \rho_{1f}^2$, where $\Gamma_a$ is the triggered axion decay rate. The effective interaction time $T$ is the target length $L$ divided by the light velocity, since a short interaction time limit applies (the long interaction time requires $L/c > 1/\Gamma_a$ which is practically impossible for a realistic target). The on-shell TRACA rate for atomic transitions is then calculated as

$$\Gamma_2 = n_a \Gamma_a L \sigma_{scat} N_f^2 \rho_{10}^2 F_M(\vec{\beta}_a), \quad \sigma_{scat} = 36\pi m_a \frac{\gamma_{ni} \gamma_{nf} \epsilon_{1f}^2}{\epsilon_{ni}^3 \epsilon_{nf}^3},$$

(3)

$$\Gamma_a = \frac{1}{(2\pi)^2} \left( \frac{c_{\alpha\gamma\gamma} \alpha}{m_a f_a} \right)^2 m_a^3 n_t \frac{\Gamma_t}{m_a - 2\omega_t + \Gamma_t^2/4} F_M \left( \frac{k_t + k_a}{m_a} \right),$$

(4)

with $|k_t| = m_a/2, |k_a| = \epsilon_{1f} + m_a/2$. Note that the axion mass dependence of the on-shell rate is $\propto n_a \Gamma_a \sigma_{scat} \propto m_a^3$, since we take $m_a n_a = \rho_G, m_a f_a \sim m_\pi f_\pi$ as fixed. The axion mass dependence $\propto m_a^3$ in this formula differs from the off-shell rate $\propto 1/m_a^2$.

**TRACA rates for p-H\textsubscript{2} vibrational transition** A typical experimental scheme we consider uses two excitation lasers and a trigger field, all of which may be irradiated near the same axis direction. The target is of a cylindrical shape whose length and volume are denoted by $L, V$.

Consider para-hydrogen molecule as a high density gas target, its gas density at the liquid nitrogen temperature 77 K being $\sim 5 \times 10^{19}$ cm$^{-3}$. The vibrational transition of electronically ground state, $v = 1 \rightarrow 0$ of level spacing 0.52 eV, may be used. The electronically excited level $(1s)(2p)$ is far away from the ground vibrational levels, separated by $\sim 10$ eV.
To calculate the signal rate for molecules, we need to take into account vibrational modes associated with electronically excited states such as (1$s$)(2$p$). The effect of Franck-Condon factor related to the potential curve of this first electronically excited state was calculated in [20], and this gives a reduction \( \sim 0.019 \) for \( \gamma_{ni} \gamma_{nf} \), which is included in the following.

The two rates, those of the off-shell and the on-shell TRACA rates, are numerically

\[
\frac{d\Gamma_1}{d\Omega} \sim 3.1 \times 10^8 \text{sec}^{-1} \left( \frac{10 \mu\text{eV}}{m_a} \right)^2 x_t (1 - x_t)^3 X, \quad \Gamma_2 \sim 1.5 \times 10^{-4} \text{sec}^{-1} \left( \frac{m_a}{10 \mu\text{eV}} \right) \frac{m_a}{\Gamma_t} \frac{L}{10 \text{cm}} X, \quad (5)
\]

\[
X = c_{a\gamma\gamma}^2 \sin^2 \theta_{ts} \left( \frac{n_T}{5 \times 10^{19} \text{cm}^{-3}} \right)^2 \left( \frac{V}{0.5 \text{cm}^3} \right)^2 \frac{n_t}{10^{18} \text{cm}^{-3}} \left( \frac{\rho_{10}}{0.1} \right)^2 \xi, \quad x_t = \frac{\omega_t}{\epsilon_{if}}, \quad (6)
\]

with \( \xi = (\sqrt{2}\pi\beta_0)^3 F_M(\tilde{\beta}_a), \tilde{\beta}_a = (\vec{p}_{if} - \vec{k}_t - \vec{k}_s)/m_a, |\vec{k}_s| \sim \epsilon_{if} - \omega_t \). The on-shell TRACA rate depends on the rf width \( \Gamma_t \).

We illustrate total rates of the off-shell and the on-shell TRACA contributions in Fig(2). There are higher sensitivities for smaller axion masses less than 100 \( \mu\text{eV} \).

![TRACA counting rate](image)

**Figure 2: TRACA rate for p-H\(_2\).** Assumed parameters are the target number density \( 5 \times 10^{19} \text{cm}^{-3} \), the target volume \( 0.5 \text{cm}^3 \), the target length \( 10 \text{cm} \), the coherence \( \rho_{10}/\Gamma_t = 10^4 \), \( \xi = 1 \), and \( \omega_t = 0.124 \text{eV} \). We used \( c_{a\gamma\gamma} \) values for the KSVZ and DFSZ models, and the perpendicular polarization of the trigger and the signal was taken, using the formula (5).

**Raman excitation and trigger** We consider Raman-type of excitation from the ground \( v = 0 \) to \( v = 1 \) state, with frequencies, \( \omega_1 > \omega_2, \omega_1 - \omega_2 = 0.52\text{eV} \). In the on-shell process the trigger frequency may be in the rf or microwave frequency range, and the momentum conservation dictates the signal photon emission very near the excitation axis. On the other hand, in the off-shell case it can be in a near infra-red range. For example, using CO\(_2\) laser of \( \omega_t = 0.124\text{eV} \) for the trigger, one can relate the angles of trigger irradiation and signal emission as functions of the axion mass: \( \theta_t = 11 \text{mrad} \sqrt{m_a/10 \mu\text{eV}} \) and \( \theta_s = -3.5 \text{mrad} \sqrt{m_a/10 \mu\text{eV}} \) away from both the trigger and the excitation lasers.

The off-axis configuration and its kinematic constraint is illustrated in Fig(3). This much of the off-axis configuration may be crucial in rejecting various QED backgrounds (see below on more of this).

**Backgrounds and experimental scheme** By far the largest background may arise from (macro-
Figure 3: Kinematic configuration of initial Raman excitation in the left, the final decay processes in the right. Both energy and momentum conservations are satisfied for PSR only when the trigger ($\omega'_t$) is injected collinearly with the excitation lasers, while they are satisfied for TRACA when the trigger ($\omega_t$) is injected with angle due to the axion absorption. These different configurations help separate two processes.

coherently amplified) PSR process, whose rate is $\sim 2 \times 10^{23} \text{sec}^{-1} \frac{m_a}{2 \pi \Gamma} \exp\left[-\frac{m_a^2}{2 \Gamma^2}\right]$ assuming a gaussian tail of Doppler and collisional broadening and taking the same set of target and trigger parameters as indicated in Fig(2). To the best of our knowledge, the tail of line-width far away from a normal dispersion is not well known. We assume a gaussian function certainly valid for the Doppler broadening. The relevant relaxation parameter for p-H$_2$ is $\Gamma = 200 \text{MHz} \times 2\pi$. For an axion mass range of a few times $10 \mu\text{eV}$ (its precise value to be determined with actual experimental design), the gaussian tail of relaxation is not a serious problem as the background.

Solid p-H$_2$ target is promising from a number of reasons; (1) a larger molecule density like $2.6 \times 10^{22} \text{cm}^{-3}$, (2) expected small relaxation caused by phonon emission. Relevant phonon emission involves the axion-proton coupling predominantly of the spin operator. The spatial extension of axion wave function is of order cm for $10 \mu\text{eV}$ mass, hence the axion field feels a collective body of para-hydrogen molecule. Since the molecule has a total nuclear spin 0, this coupling, hence the phonon relaxation, is highly suppressed.

Perhaps the most important message of this work is a short observation time for axion detection. Our estimate for solid para hydrogen shows that observation of a few months is sufficient to search the mass range of 5-15 $\mu\text{eV}$ with sensitivity reaching the DFSZ model.

In summary, the axion can be used to probe how the galactic dark matter is distributed and to determine the fundamental symmetry breaking scale beyond the standard particle theory. We proposed a novel experimental method using excited atoms or molecules which can be implemented in small scale laboratories.

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