Resonant Axion Radiation Conversion in Solar Spicules

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It has recently been observed that solar spicules covering almost of all solar surface have strong magnetic field $B \sim 10^5 G$. They are supposed to be plasma jets emitted from chromosphere and they arrive up to $\sim 10^4 km$. Their electron number density is such that $n_e = 10^{10} cm^{-3} \sim 10^{12} cm^{-3}$. Corresponding plasma frequency $m_p = e^2 n_e/m_e$ (electron mass $m_e$) is nearly equal to axion mass $m_a = 10^{-5} eV \sim 10^{-3} eV$. Thus, resonant radiation conversion of axion with the mass can arise in the spicules. We show that radiations converted from axion dark matter possess flux density $\sim 10^{-6} Jy (m_a/10^{-5}eV)(B/3 \times 10^5 G)^2$. The radiations show line spectrum with frequency $\simeq 24 GHz (m_a/10^{-4}eV)$. Our estimation has fewer ambiguities in physical parameters than similar estimation in neutron stars because physical parameters like electron number density have been more unambiguously observed in the sun. But, much strong solar thermal radiations would preclude sensitive observations of such radiations from the axions.

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Axion is the Goldstone boson of Peccei Quinn symmetry[1], which naturally solves strong CP problem. Such an axion is called as QCD axion. The axion is a promising candidate of dark matter in the Universe. Especially, the allowed mass range$[2, 3]$ of the QCD axion is restricted such as $m_a = 10^{-6} eV \sim 10^{-3} eV$.

Many projects$[4, 14]$ for the detection have been proposed and are undergoing at present. Most of them intend to detect the axion dark matter present in the Universe. On the other hand, the axions are abundantly produced$[15]$ in the center of the sun when their masses are order of 1keV. Such axions are converted to X rays under magnetic field. Helioscope of axion are experiments for the detection of such X rays.

The axion is converted to photon under external magnetic field. There are proposals for the detection of such radiations from astrophysical objects like neutron stars$[16, 19]$ and white dwarfs$[20]$. They have strong magnetic fields $\sim 10^{12} G$ and $\sim 10^6 G$, respectively so that the axion radiation conversion is strongly enhanced. The proposals are for the detection of the axion dark matter, which are dilute and ubiquitously present in the Universe. On the other hand, some of the axions may condense$[21, 22]$ in early universe and form axion stars$[23, 25]$. They are much dense localized objects of axions. Thus, they can produce strong radiation when they collide these magnetic stars$[20, 30]$ or K and M types main sequence stars$[31]$ with strong magnetic field $\sim 10^3 G$.

Fast radio bursts (FRBs) are phenomena of radio emission of huge energies with finite bandwidth. Their generation mechanisms are still unknown. A generation mechanism has been proposed as axion star collision$[20]$ with neutron star. Because the axion star is a dense object of axions, radiations from the axion star under strong magnetic field of neutron star are strong enough to be consistent with observed flux of FRBs. Although the generation mechanism is interesting, the model is not convincing. This is because axions themselves are still hypothetical objects and the detail of the physical properties of their magnetosphere has not been observed.

Although the magnetic fields of neutron stars are strong and relevant to axion radiation conversion, their distances from the earth are too far for the detail of their magnetosphere to be clearly observed. Thus, there are ambiguities in the estimation of axion radiation conversion. In this letter, we examine radiation conversion of axion dark matter in the sun. The solar atmosphere is better observed than that of neutron star. In particular, because the recent observation$[32]$ shows the presence of strong magnetic field $\sim 10^5 G$ in solar spicules$[33, 34]$, we can expect that the axion radiation conversion effectively proceeds in the spicules. Indeed, we can show that radiation flux $\sim \mu Jy$ by resonant axion radiation conversion in the spicules is produced. Our target is the axion whose mass is in the range $m_a = 10^{-5} eV \sim 10^{-3} eV$, as explained soon below.

The solar spicules$[33]$ are supposed to be plasma jets emitted from chromosphere to colona. Their width is of the order of $10^3 km$ and they reach at the lower part of colona. They are considered to be objects carrying sufficiently large amount of energies to the colona to keep high temperature. The jets rise up and fall back (or disappear) with their life times $1 \sim 10$ minute. Although their life times are of the order of minutes, their number ($>10^6$) in the solar surface are so large that they are ubiquitous objects covering almost of all solar surface at any time. Thus, we may suppose that they are tubes with radius $\sim 10^3 km$ and lengths $\sim$ several $10^6 km$. In our estimation of axion radiation conversion, we assume that there are many such tubes of spicules covering almost of all solar surface.
It should be noticed that the spicules carry electron plasma whose number density is roughly \(10^{10} \text{cm}^{-3} \sim 10^{12} \text{cm}^{-3}\). The density is higher as the part of the spicule is lower. Thus, corresponding plasma frequency \(m_p = \frac{\sqrt{e}}{m_e}(15 \text{GHz for } n_e = 7 \times 10^{13} \text{cm}^{-3})\) is approximately equal to the axion mass \(m_a = 10^{-5} \text{eV}\). Therefore, we expect that the axion radiation conversion resonantly arises [19] in the spicules when the axion mass is in the range \(m_a = 10^{-9} \text{eV} \sim 10^{-4} \text{eV}\). We notice that the resonance conversion arises when \(m_p \simeq m_a\). The resonant conversion of axion with larger mass arises in the lower part of the spicule. Because the electron number density becomes larger as we approach photosphere closer, the axion with \(m_a = 10^{-4} \text{eV}\) corresponding to \(n_e = 7 \times 10^{12} \text{cm}^{-3}\) can be resonantly converted to radiation even if it is below the spicules. Such a region just below the spicules has stronger magnetic field than that in the spicules themselves, because magnetic flux tube with strong magnetic field \(\sim 10^4 \text{G}\) is present below the spicules.

The point in our discussion is that the spicules, in particular, type II spicules are ubiquitous objects with relatively strong magnetic field of the order of \(10^5 \text{G}\) and that their electron number density \(n_e\) is relevant for resonant axion radiation conversion [19] when \(m_a = 10^{-5} \text{eV} \sim 10^{-4} \text{eV}\). The mass range would be most promising among the allowed QCD axion mass \(m_a = 10^{-9} \text{eV} \sim 10^{-3} \text{eV}\).

First, we show how the axion \(a(t, \vec{x})\) is converted to radiation under magnetized electron plasma. The axion satisfies the field equation,

\[
(\partial^2 - \vec{\partial}^2 + m_a^2)a(t, \vec{x}) = -g_{a\gamma\gamma}\vec{E} \cdot \vec{B}
\]

with \(g_{a\gamma\gamma} = g_{\gamma}\alpha/f_a\pi\), where \(\alpha \simeq 1/137\) denotes fine structure constant and \(f_a\) does axion decay constant satisfying the relation \(m_a f_a \simeq 6 \times 10^{-6} \text{eV} \times 10^{12} \text{GeV}\) in the case of QCD axion under present consideration. The parameter \(g_{\gamma}\) depends on the axion model, i.e., \(g_{\gamma} \simeq 0.37\) for DFSZ model [35, 36] and \(g_{\gamma} \simeq -0.96\) for KSVZ model [37, 38].

We note that the coupling between electromagnetic fields and axion is extremely small. For instance, the equation implies that \(m_a a \sim m_a^{-1} g_{a\gamma\gamma}\vec{E} \cdot \vec{B} \sim 10^{-33} \text{GeV}^2\) for \(E \sim B = 10^3 \text{G}\) and \(m_a = 10^{-5} \text{eV}\). On the other hand, the quantity \(m_a a_d\) in the axion dark matter density \(m_a^2 a_d^2 \sim 0.1 \text{GeV/cm}^3\) is much bigger; \(m_a a_d \sim 10^{-21} \text{GeV}^2\). Therefore, the axion coupling with electromagnetic fields can be treated perturbatively when we consider the axion dark matter in solar magnetic field.

The field equation is derived from the axion photon coupling,

\[
L_{a\gamma\gamma} = g_{a\gamma\gamma}a(t, \vec{x})\vec{E} \cdot \vec{B},
\]

with electric \(\vec{E}\) and magnetic \(\vec{B}\) fields.

They satisfy the modified Maxwell equations,

\[
\vec{\partial} \cdot (\vec{E} + g_{a\gamma\gamma}a(t, \vec{x})\vec{B}) = 0, \quad \vec{\partial} \times (\vec{B} - g_{a\gamma\gamma}a(t, \vec{x})\vec{E}) - \partial_t (\vec{E} + g_{a\gamma\gamma}a(t, \vec{x})\vec{B}) = \vec{J},
\]

\[
\vec{\partial} \cdot \vec{B} = 0, \quad \vec{\partial} \times \vec{E} + \partial_t \vec{B} = 0.
\]

where the electric current \(\vec{J}\) of the magnetized electron plasma is given in terms of electron velocity \(\vec{v}\) such as \(\vec{J} = e n_e \vec{v}\) with electron number density \(n_e\).

We solve the above equations in addition to the equation of motion of electron,

\[
m_e \frac{d\vec{v}}{dt} = e\vec{E}
\]

with electron mass \(m_e\).

The axion dark matter propagates into solar atmosphere from outside of the sun. In particular it passes a solar spicule. Then, it feels the magnetic field of the spicule and its field configuration is modified. Originally, it has an energy \(\omega\) and momentum \(\vec{k}_a\), \(a(t, \vec{x}) = a_0(\vec{x}) \exp(\omega t - i\vec{k}_a \cdot \vec{x})\) with a constant \(a_0\); \(\omega = \sqrt{m_a^2 + k_z^2}\). We write modified field configuration such that \(a(t, \vec{x}) = a_0(\vec{x}) \exp(\omega t - i\vec{k}_a \cdot \vec{x})\) where the dependence on \(\vec{x}\) of \(a_0(\vec{x})\) describes the modification. The electromagnetic fields \(\delta\vec{E}\) and \(\delta\vec{B}\) are produced when the axion passes the electron plasma with external magnetic field \(\vec{B}_{ext}\). They are described in the following equations,

\[
\vec{\partial} \cdot \delta\vec{E} = 0, \quad \vec{\partial} \times \delta\vec{B} - \partial_t (\delta\vec{E} + g_{a\gamma\gamma}a(t, \vec{x})\vec{B}_{ext}) = \vec{J},
\]

\[
\vec{\partial} \cdot \delta\vec{B} = 0, \quad \vec{\partial} \times \delta\vec{E} + \partial_t \delta\vec{B} = 0.
\]
where we have taken into account the fact that the electric charge is immediately screened in the electron plasma, which leads to the equation \( \vec{\nabla} \cdot \delta \vec{E} = 0 \). The axion is governed by the following equation,

\[
(\partial_t^2 - \partial^2 + m_a^2) a(t, \vec{x}) = - g_{a\gamma\gamma} \delta \vec{E} \cdot \vec{B}_{\text{ext}}
\]  

(6)

These coupled equations(4), (5) and (6) describe axion radiation conversion in the electron plasma. We find from these equations that the electric field of the radiation converted from the axion has a component parallel to the magnetic field \( \vec{B}_{\text{ext}} \). In other words, \( \delta \vec{E} \propto \vec{B}_{\text{ext}} \).

It is easy to derive the following equations by using the ansatz \( \delta \vec{E} = \vec{E}_0(\vec{x}) \exp(i\omega t - i\vec{k}_a \cdot \vec{x}) \),

\[
- i \partial_r E(r) + \frac{1}{2k_a} \left( (m_a^2 - m_p^2(r))E(r) + \Delta B a_0(r) \right) = 0
\]

\[
- i \partial_r a_0(r) - \frac{1}{2\omega^2 k_a} \Delta B E(r) = 0 \quad \text{with} \quad E = \frac{\vec{E}_0 \cdot \vec{B}_{\text{ext}}}{|\vec{B}_{\text{ext}}|}
\]

(7)

with \( |\vec{k}_a| = k_a, \Delta_B \equiv \omega^2 g_{a\gamma\gamma} B_{\text{ext}} \) and \( \partial_r \equiv \vec{k}_a \cdot \vec{\nabla}/k_a \) where we assumed that \( \vec{E}_0(r) \) and \( a_0(r) \) do not depend the transverse coordinate \( \vec{x}_t \) defined such as \( \vec{k}_a \cdot \vec{x}_t = 0 \). We note that \( r = \vec{k}_a \cdot \vec{x}/k_a \). Additionally, we used the condition \( k_a \partial_r E(r) \gg \partial_r^2 E(r) \) and \( k_a \partial_r a_0(r) \gg \partial_r^2 a_0(r) \), because the de Broglie wave length \( 1/k_a \sim 10^2\text{cm}(10^{-4}\text{eV}/m_a) \) of the axion dark matter is much small compared with the typical scale \( 10^3\text{km} \) of the spicules.

Now we specify the form of the spicules(33, 34). We suppose that it is a tube with radius \( R = 10^2\text{km} \) and length \( H = 5 \times 10^3\text{km} \) extending to the direction \( z \). (It turns out below that total radiation flux from the sun does not depend on the radius \( R \).) We assume that the external magnetic field \( \vec{B}_{\text{ext}} \) is parallel to the tube and that \( \vec{B}_{\text{ext}} = 3 \times 10^2\text{G} \). Namely, the plasma jet forming the tube makes the magnetic field point to the direction of the jet. We suppose that \( \vec{B}_{\text{ext}} = (0, 0, B_0) \) points to \( z \) direction and has no dependence on the coordinate \( z \). (The axion radiation conversion only arises at the vicinity of a region with \( m_a \approx m_p \) so that the relevant magnetic field is the one present at the region.) Furthermore, we suppose that the distribution of the electron number density \( n_e \) is such that \( n_e(z) = n_e \exp(-z/H) \). The coordinate \( z \) is taken such that the plasma frequency \( m_p(z = 0) = \sqrt{e^2n_e(z = 0)/m_e} = \sqrt{e^2n_e/m_e} \) is equal to the axion mass \( m_a \). That is, the resonant conversion arises in the vicinity at \( z = 0 \). We rewrite the distribution \( n_e = n_e \exp(-z/H) \) such that \( n_e(r) = n_e \exp(-rk_z/(Hk_a)) = n_e \exp(-r/H') \) with \( k_a = (k_x, k_y, k_z), H' = Hk_a/k_z \) and \( z = rk_z/k_a \) because \( z = r \cos \phi \) and \( k_z = k_a \cos \phi \).

Under these conditions, we solve the equations (6) and (7) by neglecting the term of the order of \( \Delta^2_B \),

\[
E(r) = \frac{-ia_0(r = 0)}{2k_a} \int_0^r \Delta_B dr' \exp \left( i \frac{2k_a}{r} \int_0^{r'} (m_a^2 - m_p^2(r''))dr'' \right)
\]

(8)

\[
a_0(r) = a_0(r = 0) + \frac{i}{2\omega^2 k_a} \int_0^r \Delta_B E(r')dr'
\]

(9)

with \( E(r = 0) = 0 \) because there are no radiations before axion comes in. The axion passes through the tube from \( r = 0 \) to \( r = \infty \). The formula in eq(5) shows that the axion radiation conversion resonantly arises in the vicinity at \( r = 0 \). We note that the axion velocity \( k_a/m_a \sim 10^{-3} \) is very small. Thus, the integration over \( r'' \) is controlled only around \( r'' = 0 \) because \( m_a = m_p(r'' = 0) \). Thus, we have \( m_a^2 - m_p^2 = m_a^2(1 - \exp(-r/H')) \approx m_a^2/r/H' \). So, we find

\[
E(r) = \frac{-ia_0(r = 0)\Delta_B}{2k_a} \int_0^{r} dr' \exp \left( \frac{i m_a^2 r'^2}{4k_a H'} \right) = \frac{-ia_0(r = 0)\Delta_B}{2k_a} \int_0^{r} \sqrt{\frac{4k_a H'}{m_a^2}} \int_0^{\frac{m_a^2 r'^2}{4k_a H'}} dx \exp(ix^2).
\]

(10)

We can see that \(|E(r)| \) rapidly increases from \(|E(r = 0)| = 0 \) in the vicinity at \( r = 0 \) and soon becomes constant \(|E(r)| \approx |a_0(r = 0)| \Delta_B/2k_a \sqrt{4k_a H'/m_a^2} \times \sqrt{\pi/4} \). This is because \( r \sqrt{m_a^2/4k_a H'} \approx r(m_a/10^{-4}\text{eV})/10^3\text{cm} \) with typically \( k_z \approx k_a \) (\( k_a/m_a \sim 10^{-3} \)). For instance, for \( r = 10^3\text{cm} \), \(|\int_0^{m_a^2/4k_a H'} dx \exp(ix^2)| = |\int_0^{10} dx \exp(ix^2)| \approx |\int_0^{\infty} dx \exp(ix^2)| = \sqrt{\pi/4} \). Namely the axion radiation conversion only arises in the vicinity at \( r = 0 \).

In order to calculate the flux of the radiation, we need to know corresponding magnetic field \( \delta \vec{B} \),
\[ \partial_t \delta \vec{B} = i \omega \delta \vec{B} = -\vec{E} \times \delta \vec{E} = (-\partial_0 \delta E_z, \partial_0 \delta E_z, 0) \approx (ik_y E_z, -ik_z E_z, 0) \] (11)

with \( E_z = E(r) \exp(i\omega t - ik_a \cdot \vec{x}) \), where we used the relation \( E \gg \partial_r E/k_a \). Therefore, the radiation flux \( F \) is given by

\[ F = \frac{1}{2} \int d\vec{S} \cdot (\delta \vec{E} \times \delta \vec{B}) = \frac{1}{2\omega} \int (dS_z k_z + dS_y k_y) |E(r)|^2 \]

with \( |E(r)| = \frac{|a_0(r = 0)| \Delta_B}{2k_a} \sqrt{\frac{4k_a H}{m_a^4}} \int_0^r \sqrt{\frac{a_0^2}{a_0^2 + 2}} \ dx \exp(i\chi^2) \). (13)

Here we may put \( k_y = 0 \) without loss of generality.

The surface integration is performed over the side surface of the tube with radius \( R = \sqrt{x^2 + y^2} \). The surface is the one such that the radiations with the momentum \( k_a \) produced at the points \( z = 0 \) pass through. We denote the coordinate of the points at \( z = 0 \) as \( \vec{A} = (\vec{0}, 0) \) with \( |\vec{0}| \leq R \) and the coordinates of the points on the side surface of the tube as \( \vec{x} = (\vec{r}, z) \) with \( |\vec{r}| = R \). Then, the side surface at \( z > 0 \) which the radiations emitted at \( \vec{A} \) pass through is defined such that the vector \( \vec{x} - \vec{A} \) is proportional to the momentum \( k_a \). That is, \( \vec{x} - \vec{A} = (\vec{r}, z) - (\vec{0}, 0) = (\vec{r} - \vec{0}, z) = k_a (\vec{r} - \vec{0}, z)/k_a \). Thus, it leads to

\[ \vec{r} - \vec{0} = \frac{k_0}{k_a} \sqrt{(\vec{r} - \vec{0})^2 + z^2}, \quad z = \frac{k_z}{k_a} \sqrt{(\vec{r} - \vec{0})^2 + z^2} \] (14)

with \( k_a = (k_0, k_z) = (k_0, 0, k_z) \). From these equations we can derive the relation,

\[ R^2 = (\rho_{0,t} + \frac{k_0 z}{k_z})^2 + \rho_{0,t}^2 \] (15)

with \( \rho_{0,t} \equiv |\vec{0}_{0,|t} \) and \( \rho_{0,t} \equiv |\vec{0}_{0,|t} \), where we decomposed \( \vec{0}_0 \) into transverse component \( \vec{0}_{0,t} \) ( \( \vec{0}_{0,t} \cdot \vec{k}_a = 0 \) ) and longitudinal one \( \vec{0}_{0,l} \) ( \( \vec{0}_{0,l} \propto \vec{k}_0 \), i.e., \( \vec{0}_0 = \vec{0}_{0,t} + \vec{0}_{0,l} \)). Then, \( \vec{r} = \vec{0}_0 + \vec{k}_0 z/k_z = \vec{0}_{0,t} + k_0 z/k_z + \rho_{0,t}/k_0 \). The equation (15) shows allowed values which \( \rho_{0,t} \) and \( \rho_{0,t} \) can take for \( z \) and \( k_0 \) given; \( R - k_0 z/k_z \geq \rho_{0,t} \geq 0 \) and \( \sqrt{R^2 - (k_0 z/k_z)^2} \geq \rho_{0,t} \geq 0 \) ( \( k_0 \equiv |\vec{k}_0| \)). We denote \( \vec{p} = R(\cos\theta, \sin\theta) \) where the angle \( \theta \) is defined such as \( \vec{r} \cdot \vec{k}_0 = Rk_0 \cos\theta \). So, we have the surface element \( dS_x = R \cos\theta d\theta dz = d(R \sin\theta) dz = \rho_{0,t} dz \) in the surface integral. Therefore, the surface integration in eq (12) is performed such that

\[ F = \frac{1}{2\omega} \int dS_x k_0 |E(r)|^2 = \frac{1}{2\omega} \int d(R \sin\theta) dz k_0 |E(r)|^2 = \frac{1}{2\omega} \int k_0 \rho_{0,t} dz k_0 |E(r)|^2 \]

\[ \simeq \frac{1}{2\omega} \int \sqrt{R^2 - (k_0 z/k_z)^2} k_{0,t} dz \frac{d}\partial_0, k_0 |E(r = \infty)|^2 = \frac{1}{2\omega} \int k_z \pi R^2 E(r = \infty)^2 = \frac{k_z \pi}{8m_a R^2} \left( -i a_0(r = 0) \Delta B \right)^2 \left( \frac{4k_z^2 H}{k_z m_a^2} \right)^2 \frac{\pi}{4} \frac{a_0^2 R^2 H |a_0|^2 (g_{\gamma\gamma} B_{ext})^2}{32} \] (17)

with \( \Delta_B \simeq m_\gamma^2 g_{\gamma\gamma} B_{ext}, a_0 \equiv a_0(r = 0) \) and \( |E(r)| \simeq |E(r = \infty)| \) because typically \( r \sim R = 10^7 \text{cm} \).

The radiation flux \( F \) is the one of radiations emitted from a spicule. There are many spicules which cover almost all solar surface. Thus, their number \( N \) is approximately given such that \( N = 4\pi R_d^2 / \pi R^2 = 4R_d^2 / R^2 \). Total flux \( F_{tot} \) from the solar spicules is \( F N \sim 4F(R_d/R)^2 \). Therefore, it leads to

\[ F_{tot} = \frac{4\pi^2 m_a R_d^2 H |a_0|^2 (g_{\gamma\gamma} B_{ext})^2}{32} = \frac{\pi^2 R_d^4 H |a_0|^2 (g_{\gamma\gamma} B_{ext})^2}{4m_a} \] (18)

where we have expressed \( |a_0|^2 \) by the energy density \( \rho_a \) of the dark matter axion; \( \rho_a = m_a^2 |a_0|^2 / 2 \sim 0.3 \text{GeV/cm}^4 \). The total flux \( F_{tot} \) does not depend on the radius \( R \) of each spicule because we have assumed that the spicules cover almost all solar surface. Numerically,
\[
F_{\text{tot}} \sim 2.3 \times 10^{-4} W \left( \frac{B_{\text{ext}}}{3 \times 10^8 \text{G}} \right)^2 \frac{H}{5 \times 10^4 \text{km}} \frac{m_a}{10^{-4} \text{eV}}
\]

with \( R_\odot \approx 7 \times 10^{5} \text{km} \).

Taking account of the distance \( D \approx 1.5 \times 10^{8} \text{km} \) from the earth to the sun, we obtain the observed flux density,

\[
S_\nu = \frac{F_{\text{tot}}}{4\pi D^2 \delta \nu} \approx 2.3 \times 10^{-6} \text{Jy} \left( \frac{B_{\text{ext}}}{3 \times 10^8 \text{G}} \right)^2 \frac{H}{5 \times 10^4 \text{km}} \frac{m_a}{10^{-4} \text{eV}}
\]

with \( \delta \nu \sim 10^{-6} m_a/2\pi \approx 24 \text{KHz} \). The width \( \delta \nu \) of the radiation frequency is given by the energy width \( \delta \omega \) of the axion dark matter \( \omega = \sqrt{m_a^2 + k_a^2} \approx m_a + k_a^2/2m_a \equiv m_a + \delta \omega \) with \( k_a \sim 10^{-3} m_a \).

In our estimation we have assumed the height of spicules 5 \( \times \) 10^3 km and the strength of the magnetic field 3 \( \times \) 10^5 G. These are based on the observations and are not unrealistic. The ambiguous point is the distribution of the spicules. Namely, we do not know the occupation fraction of the spicules to the solar surface. Especially, among spicules are type II spicules having small widths and abundantly distributed over the surface. Their filling factor \( f \) has not been observed, although it has been speculated; \( N \) denotes the number of the spicules with radius \( R \). (Actually, the type II spicules with width less than 100 km have not yet been observed clearly.) We have assumed that they cover the whole of the solar surface. That is, the filling factor \( f \approx 1 \). But, in reality, it might be much smaller, e.g. \( f \approx 10^{-2} \). Then, the flux density is reduced by a factor \( 10^2 \). The satellites Solar-C launched in near future would make clearer the distribution of the spicules with smaller scales.

Solar magnetic field is much smaller than that of neutron star, but the distance from the sun to the earth is much shorter than that from neutron star. Hence, the observations in detail of the solar surface are possible, which can clarify physical properties of spicules. Such observations are difficult in neutron star. The solar observation makes possible the more precise estimation of the axion radiation conversion in the sun than that in neutron star.

Although we have estimated the radiation flux \( \sim 10^{-6} \text{Jy} \) from the axion dark matter, the radiation is extremely weaker than those of solar radiations with frequencies \( \sim 10^6 \text{GHz} \). They are of the order of \( 10^6 \text{Jy} \). It would be difficult to observe such a radiation from the axion.

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