Probing microscopic origins of axions by the chiral magnetic effect

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Workshop on Physics of Dark Cosmos, Busan, Korea, Oct 22nd, 2022

Outline

- 3 possible classes of microscopic origins of axions
- (Measurable) distinctive patterns of axion couplings depending on the microscopic origins
- Measuring the axion-electron coupling by the chiral magnetic effect

Strong CP problem and QCD axion

$$y_u H Q_L u_R^c + y_d H^* Q_L d_R^c + \frac{g_s^2}{32\pi^2}\theta G\tilde{G}$$

$$\bar{\theta} = \theta + \arg\,\det\,(y_u y_d) < 10^{-10}$$

Non-observation of neutron EDM

[Abel et al ’20]

CPV in the QCD sector

while $\delta_{\rm CKM} = \arg\,\det\left[y_u y_u^\dagger, y_d y_d^\dagger\right] \sim \mathcal{O}(1)$

The QCD vacuum energy is minimized at the CP-conserving point ($\bar{\theta} = 0$).

[Vafa, Witten ’84]

$$V_{\rm QCD} = -\Lambda_{\rm QCD}^4 \cos\bar{\theta}$$

Promote $\bar{\theta}$ to a dynamical field (=QCD axion) : $\frac{g_s^2}{32\pi^2}\left(\theta + \frac{a}{f_a}\right)G\tilde{G}$

[Peccei, Quinn ’77, Weinberg ‘78, Wilczek ’78]

QCD axion lagrangian

$$\mathcal{L} = \frac{1}{2}(\partial_\mu a)^2 + \frac{g_s^2}{32\pi^2} c_G \frac{a}{f_a} G^{\mu\nu}\tilde{G}_{\mu\nu}$$

$$+ \frac{a}{f_a} \sum_{A=W,B,\ldots} \frac{g_A^2}{32\pi^2} c_A F^{A\mu\nu}\tilde{F}^A_{\mu\nu} + \frac{\partial_\mu a}{f_a} \left(\sum_{\psi=q,\ell,\ldots} c_\psi \psi^\dagger \bar{\sigma}^\mu \psi + \sum_{\phi=H,\ldots} c_\phi \phi^\dagger i \overleftrightarrow{D}^\mu \phi \right)$$

$$U(1)_{PQ}: \quad a(x) \to a(x) + \alpha$$

broken by $c_G \neq 0$ non-perturbatively

➡

$$m_a^2 \simeq c_G^2 \frac{m_u m_d}{(m_u + m_d)^2} \frac{m_\pi^2 f_\pi^2}{f_a^2}$$

The axion couplings to the other SM particles

$c_W, c_B, c_q, c_\ell, c_H$ are UV model-dependent.

Axion-Like Particles (ALPs)

- Cousins of the QCD axion, while not being necessarily involved in the strong CP problem (so c_G can be 0)
- Ubiquitous in many BSM scenarios, in particular, string theory

[Arvanitaki, Dimopoulos, Dubovsky, Kaloper, Marsh-Russell, '09]

$$\boxed{\frac{1}{2}(\partial_\mu a)^2 - \frac{1}{2}m_a^2 a^2 + \frac{a}{f_a}\sum_A \frac{g_A^2}{32\pi^2} c_A F^{A\mu\nu}\tilde{F}^A_{\mu\nu} + \frac{\partial_\mu a}{f_a}\left(\sum_\psi c_\psi \psi^\dagger \bar{\sigma}^\mu \psi + \sum_\phi c_\phi \phi^\dagger i \overset{\leftrightarrow}{D}^\mu \phi\right)}$$

i) approximate shift symmetry $U(1)_{PQ}$ $\quad a(x) \;\to\; a(x) + c \;\; (c \in \mathbb{R})$

: ALP can be naturally light.

ii) periodicity $\quad \dfrac{a(x)}{f_a} \equiv \dfrac{a(x)}{f_a} + 2\pi$

: f_a characterizes typical size of ALP couplings up to dimensionless parameters c_A, c_ψ, c_ϕ.

KSVZ model

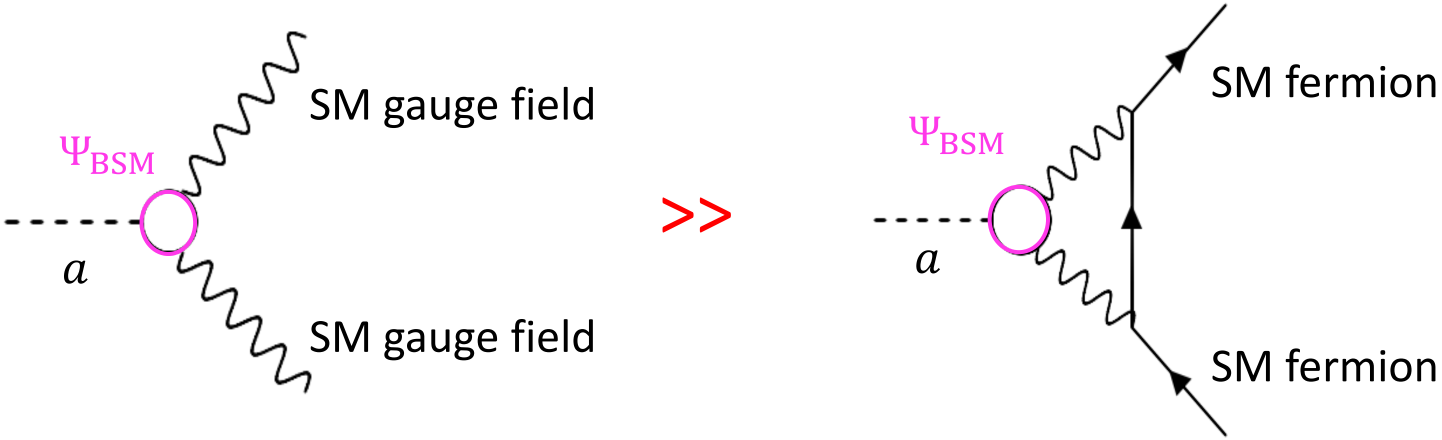
Kim '79, Shifman, Vainshtein, Zakharov '80

The axion *couples to SM fields via a heavy BSM fermion* charged under the SM gauge group.

$$\langle \Phi \rangle = \frac{1}{\sqrt{2}} f_a$$

$$\boxed{y \Phi \Psi_{\rm BSM} \Psi^c_{\rm BSM} + {\rm h.c.}} \qquad \Phi = \frac{1}{\sqrt{2}} (\rho + f_a) e^{ia/f_a}$$

$$m_\Psi = \frac{y}{\sqrt{2}} f_a$$



"KSVZ-like models"

: *no tree-level couplings* to the SM fermions

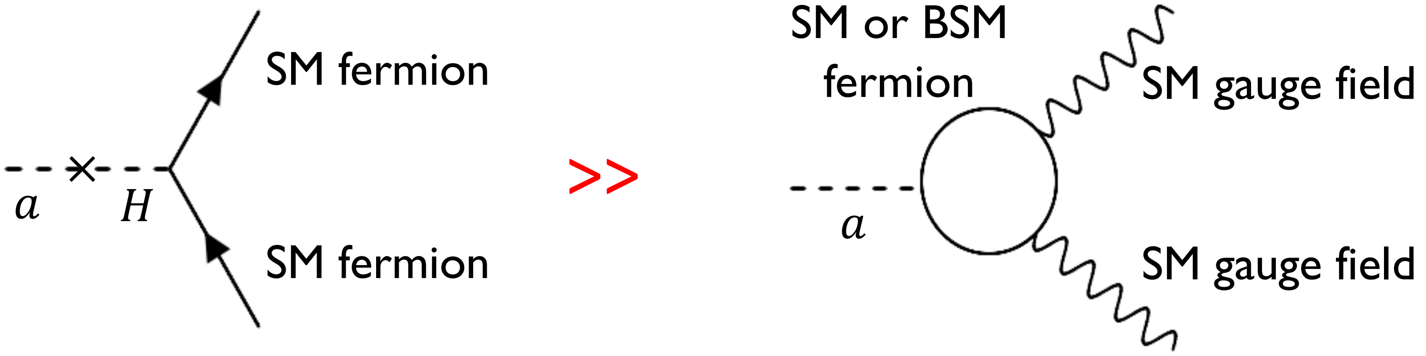
DFSZ model

Dine, Fischler, Srednicki ’81, Zhitnitsky ’80

The axion *couples to the SM sector at tree-level* through the Higgs portal.

$$\boxed{y_u u_R^c Q_L H_u + y_d d_R^c Q_L H_d + y_e e_R^c L H_d + \lambda \Phi^2 H_u H_d + \text{h.c.}}$$

$$\Phi = \frac{1}{\sqrt{2}}(\rho + f_a) e^{ia/f_a}$$



“DFSZ-like models”

: $O(1)$ *tree-level couplings* to the SM fermions

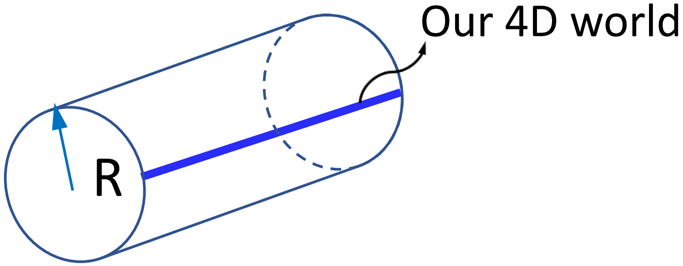
String-theoretic models

$$A_{[m_1 m_2 .. m_p]}(x^\mu, y^m) = a(x^\mu)\Omega_{[m_1 m_2 .. m_p]}(y^m)$$

Ω : harmonic p-form on the compact internal space

4D axions identified as zero modes of higher-dimensional p-form gauge field

- Simplified 5D toy model



Compact extra dimension

$$A_M(x^\mu, y) \qquad M = 0, 1, 2, 3, 5$$

$$\Rightarrow \quad A_5(x^\mu, y) = a(x^\mu)\Omega_5(y)$$

$$F_{MN} = \partial_M A_N - \partial_N A_M$$

M_5: 5D Planck mass (String scale)

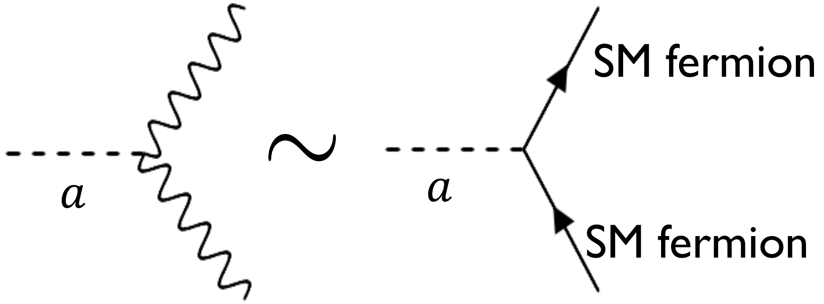
$$S = \int d^5x\sqrt{-g}\left(-\frac{1}{4g_5^2}F^{MN}F_{MN} + \underbrace{c_1\frac{\epsilon^{MNPQR}}{\sqrt{-g}}A_M G^a_{NP}G^a_{QR}}_{\substack{\text{5D Chern-Simon term}\\ \text{(axion-gauge field coupling)}}} + \underbrace{\frac{c_2}{M_5}F_{MN}\bar{\Psi}\gamma^M\gamma^N\Psi}_{\substack{\text{5D gauge-matter coupling}\\ \text{(axion-matter coupling)}}} + \cdots\right)$$

$$c_1 \sim c_2 \sim O(1)$$

Integrating over the extra dimension

$$S_{\text{eff}} = \int d^4x\left(-\frac{1}{4g^2}F^{\mu\nu}F_{\mu\nu} + \frac{1}{2}\kappa M_5^2\partial^\mu a\partial_\mu a + c_1 a G^{a\mu\nu}\widetilde{G}^a_{\mu\nu} + c_2\partial_\mu a\bar{\Psi}\gamma^\mu\gamma^5\Psi + \cdots\right)$$

$$\kappa \sim O(1)$$



String-theoretic axion couplings to matter fields and gauge fields are comparable.

$$S_{\text{eff}} = \int d^4x \left(-\frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} + \frac{1}{2} \kappa M_5^2 \partial^\mu a \partial_\mu a + c_1 a G^{a\mu\nu} \widetilde{G}^a_{\mu\nu} + c_2 \partial_\mu a \bar{\Psi} \gamma^\mu \gamma^5 \Psi + \cdots \right)$$

$$\kappa \sim c_1 \sim c_2 \sim O(1)$$

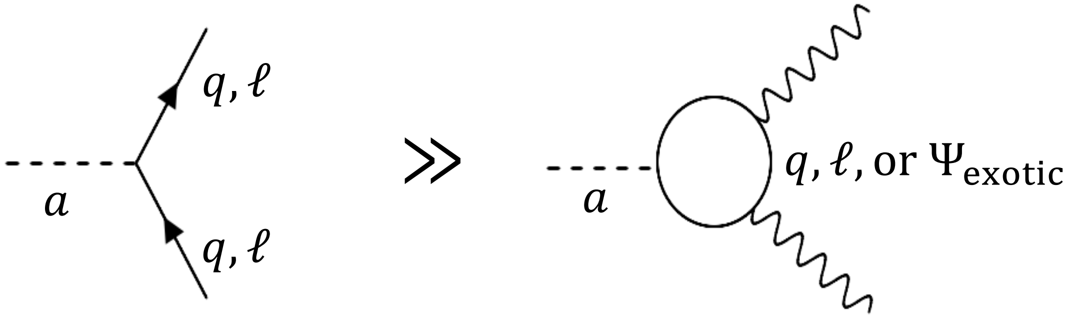
Canonical normalization $\quad a \to \dfrac{a}{32\pi^2 f_a} \qquad f_a = \dfrac{\sqrt{\kappa}}{32\pi^2} M_5$

$$S_{\text{eff}} = \int d^4x \left(-\frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} + \frac{1}{2} \partial^\mu a \partial_\mu a + \frac{c_1}{32\pi^2} \frac{a}{f_a} G^{a\mu\nu} \widetilde{G}^a_{\mu\nu} + \underbrace{\frac{c_2}{32\pi^2} \frac{\partial_\mu a}{f_a}} \bar{\Psi} \gamma^\mu \gamma^5 \Psi + \cdots \right)$$

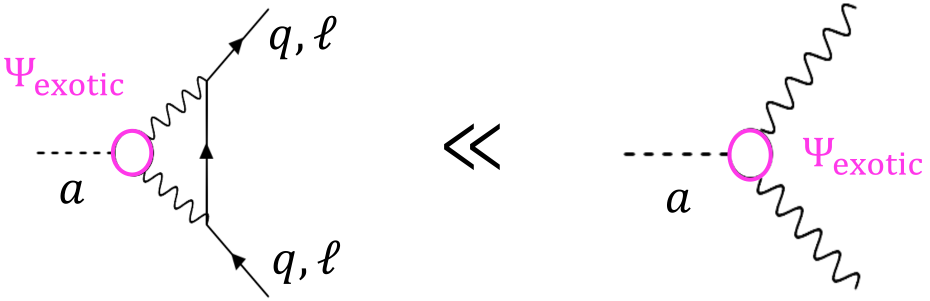
Axon-matter couplings are suppressed by the 1-loop factor.

Summary: characteristic patterns of axion couplings to the SM depending on the microscopic origins

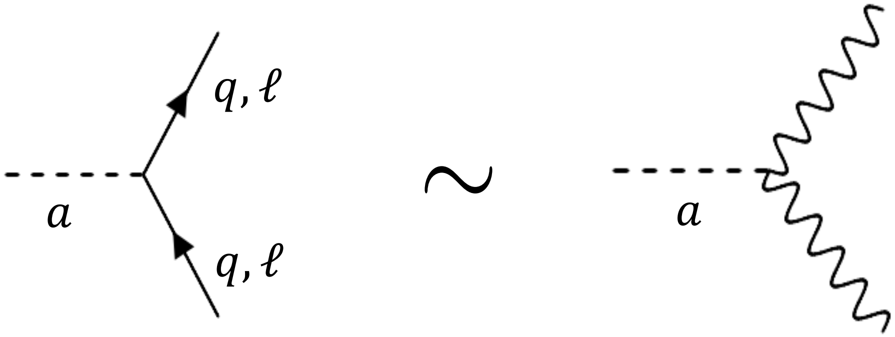
DFSZ-like models



KSVZ-like models

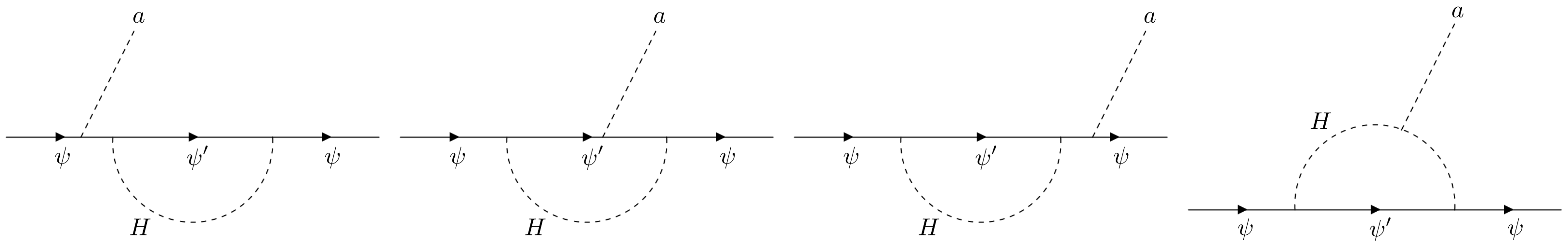


String-theoretic models



Running of axion couplings by Yukawa interactions

K Choi, SHI, CB Park, S Yun '17, Camalich, Pospelov, Vuong, Ziegler, Zupan '20
Heiles, König, Neubert '20, Chala, Guedes, Ramos, Santiago '20



$$\frac{\partial_\mu a}{f_a}\left(\sum_\psi c_\psi \psi^\dagger \bar\sigma^\mu \psi + \sum_{\alpha=1,2} c_{H_\alpha} H_\alpha^\dagger i \overset{\leftrightarrow}{D}{}^\mu H_\alpha\right) + \sum_A \frac{g_A^2}{32\pi^2} c_A \frac{a}{f_a} F^{A\mu\nu}\widetilde{F}^A_{\mu\nu}$$

$y_t u_3^c Q_{L3} H_u$ ⇒

$$\frac{dc_{Q3}}{d\ln\mu} \approx \frac{\xi_y}{16\pi^2} y_t^2 n_t$$

$$\frac{dc_{u_3^c}}{d\ln\mu} \approx \frac{\xi_y}{8\pi^2} y_t^2 n_t$$

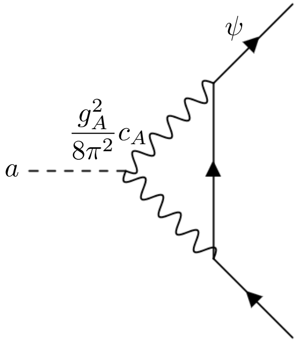
$$\frac{dc_{H_u}}{d\ln\mu} \approx \frac{3}{8\pi^2} y_t^2 n_t$$

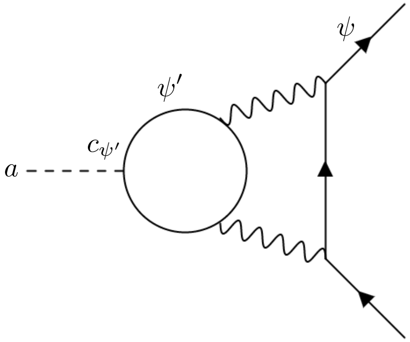
$$n_t \equiv c_{Q_3} + c_{u_3^c} + c_{H_u}$$

$$\xi_y = \begin{cases} 1 & \text{for non-SUSY models} \\ 2 & \text{for SUSY models} \end{cases}$$

Running of axion couplings by gauge interactions

Srednicki ’85, S Chang and K Choi ’93
K Choi, SHI, CS Shin ’20,
Chala, Guedes, Ramos, Santiago ‘20
Bauer, Neubert, Renner, Schnubel, Thamm ‘20





$$\frac{\partial_\mu a}{f_a}\left(\sum_\psi c_\psi \psi^\dagger \bar{\sigma}^\mu \psi + \sum_{\alpha=1,2} c_{H_\alpha} H_\alpha^\dagger i \overset{\leftrightarrow}{D}^\mu H_\alpha\right) + \sum_A \frac{g_A^2}{32\pi^2} c_A \frac{a}{f_a} F^{A\mu\nu}\widetilde{F}^A_{\mu\nu}$$

$$\left.\frac{dc_\psi}{d\ln\mu}\right|_{\rm gauge} = -\xi_g \sum_A \frac{3}{2}\left(\frac{g_A^2}{8\pi^2}\right)^2 \mathbb{C}_A(\psi)\,\tilde{c}_A \qquad \tilde{c}_A \equiv c_A - \sum_{\psi'} c_{\psi'}$$

$$\left.\frac{dc_{H_\alpha}}{d\ln\mu}\right|_{\rm gauge} = -\xi_H \sum_A \frac{3}{2}\left(\frac{g_A^2}{8\pi^2}\right)^2 \mathbb{C}_A(H_\alpha)\,\tilde{c}_A \qquad \mathbb{C}_A(\Phi) : \text{quadratic Casimir}$$

$$\xi_g = \begin{cases} 1 & \text{for non-SUSY models} \\ 2/3 & \text{for SUSY models} \end{cases}, \quad \xi_H = \begin{cases} 0 & \text{for non-SUSY models} \\ 2/3 & \text{for SUSY models} \end{cases}$$

Consequences in low energy observables

Axion couplings to the photon, electron, neutron, and proton below GeV

$$\frac{1}{4}g_{a\gamma}a\vec{E}\cdot\vec{B}+\partial_\mu a\left[\frac{g_{ae}}{2m_e}\bar{e}\gamma^\mu\gamma_5 e+\frac{g_{an}}{2m_n}\bar{n}\gamma^\mu\gamma_5 n+\frac{g_{ap}}{2m_p}\bar{p}\gamma^\mu\gamma_5 p\right]$$

$$\begin{aligned}
g_{a\gamma} &\simeq \frac{\alpha_{\rm em}}{2\pi}\frac{1}{f_a}\left(c_W+c_B-\frac{2}{3}\frac{m_u+4m_d}{m_u+m_d}c_G\right) \simeq \frac{\alpha_{\rm em}}{2\pi}\frac{1}{f_a}\Big(c_W+c_B-1.92c_G\Big),\\
g_{ap} &\simeq \frac{m_p}{f_a}\left(C_u\Delta u+C_d\Delta d-\Big(\frac{m_d}{m_u+m_d}\Delta u+\frac{m_u}{m_u+m_d}\Delta d\Big)c_G\right),\\
&\simeq \frac{m_p}{f_a}\Big(0.90(3)\,C_u(2\,\mathrm{GeV})-0.38(2)\,C_d(2\,\mathrm{GeV})-0.48(3)\,c_G\Big),\\
g_{an} &\simeq \frac{m_n}{f_a}\left(C_d\Delta u+C_u\Delta d-\Big(\frac{m_u}{m_u+m_d}\Delta u+\frac{m_d}{m_u+m_d}\Delta d\Big)c_G\right),\\
&\simeq \frac{m_n}{f_a}\Big(0.90(3)\,C_d(2\,\mathrm{GeV})-0.38(2)\,C_u(2\,\mathrm{GeV})-0.04(3)\,c_G\Big),\\
g_{ae} &\simeq \frac{m_e}{f_a}C_e(m_e),
\end{aligned}$$

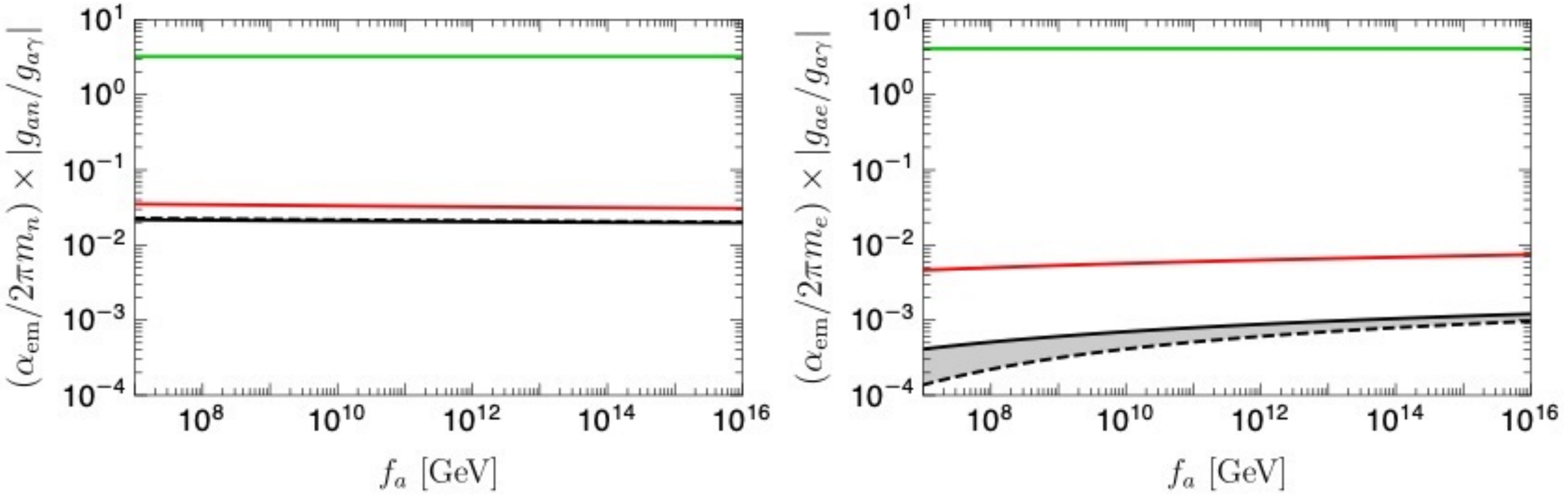
$$\underbrace{\langle p|\bar{u}\gamma^\mu\gamma_5 u|p\rangle}_{s^\mu\Delta u}$$

$$\underbrace{\langle p|\bar{d}\gamma^\mu\gamma_5 d|p\rangle}_{s^\mu\Delta d}$$

Cortona, Hardy, Vega, Villadoro ‘15

Distinguishing the models of an axion by coupling ratios

For QCD axion ($c_G \neq 0$), $g_{ap} \sim \frac{\mathrm{m_p}}{\mathrm{f_a}}$ regardless of the classes of models

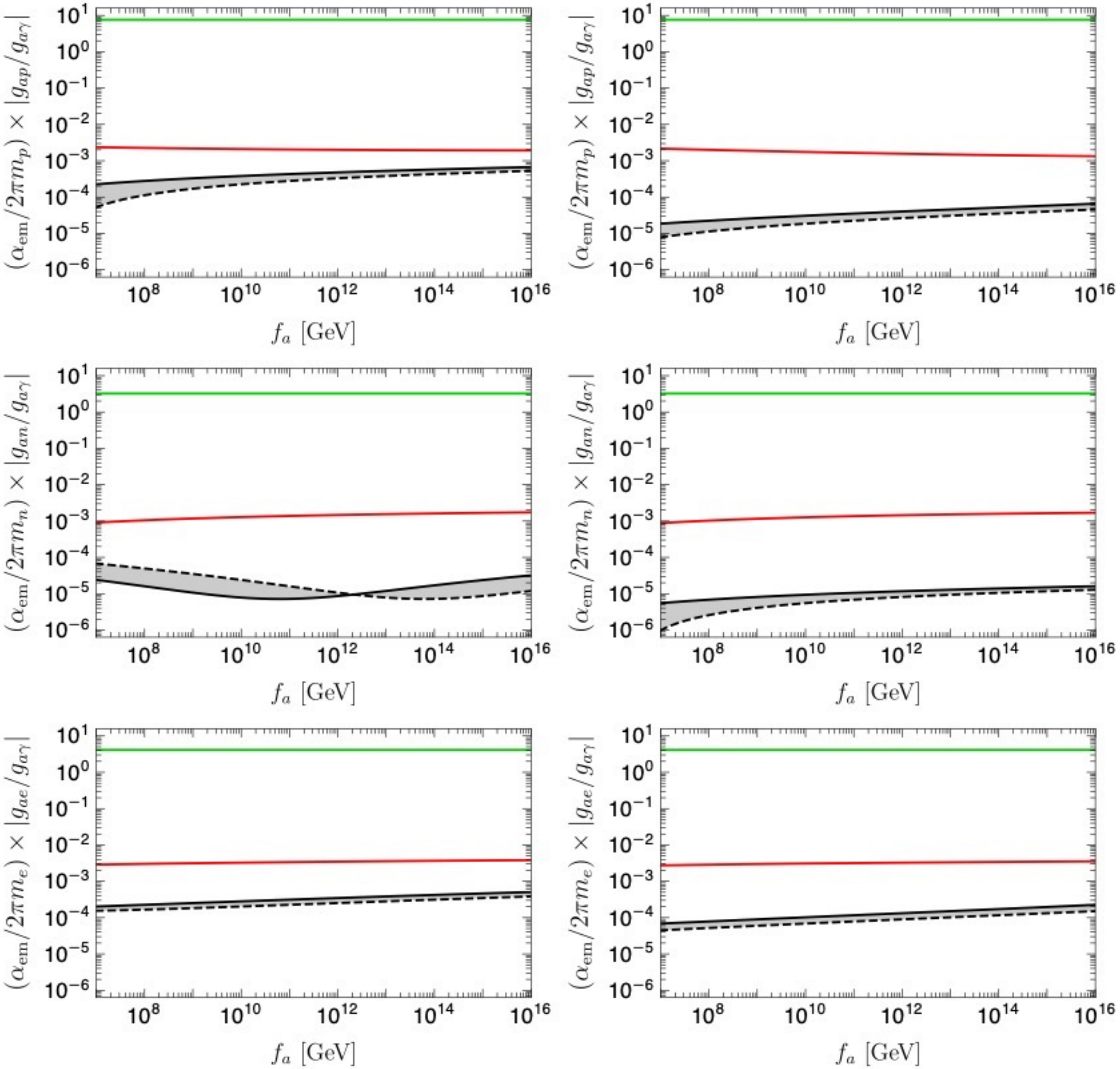


Green : DFSZ-like model

Red : String-theoretic model

Black : KSVZ-like model (dashed : $m_\Psi = 10^{-3} f_a$, solid : $m_\Psi = f_a$)

For ALPs with ($c_G = 0$),



$c_W = 1$ ($c_G = c_B = 0$)

$c_B = 1$ ($c_G = c_W = 0$)

Green : DFSZ-like model

Red : String-theoretic model

Black : KSVZ-like model

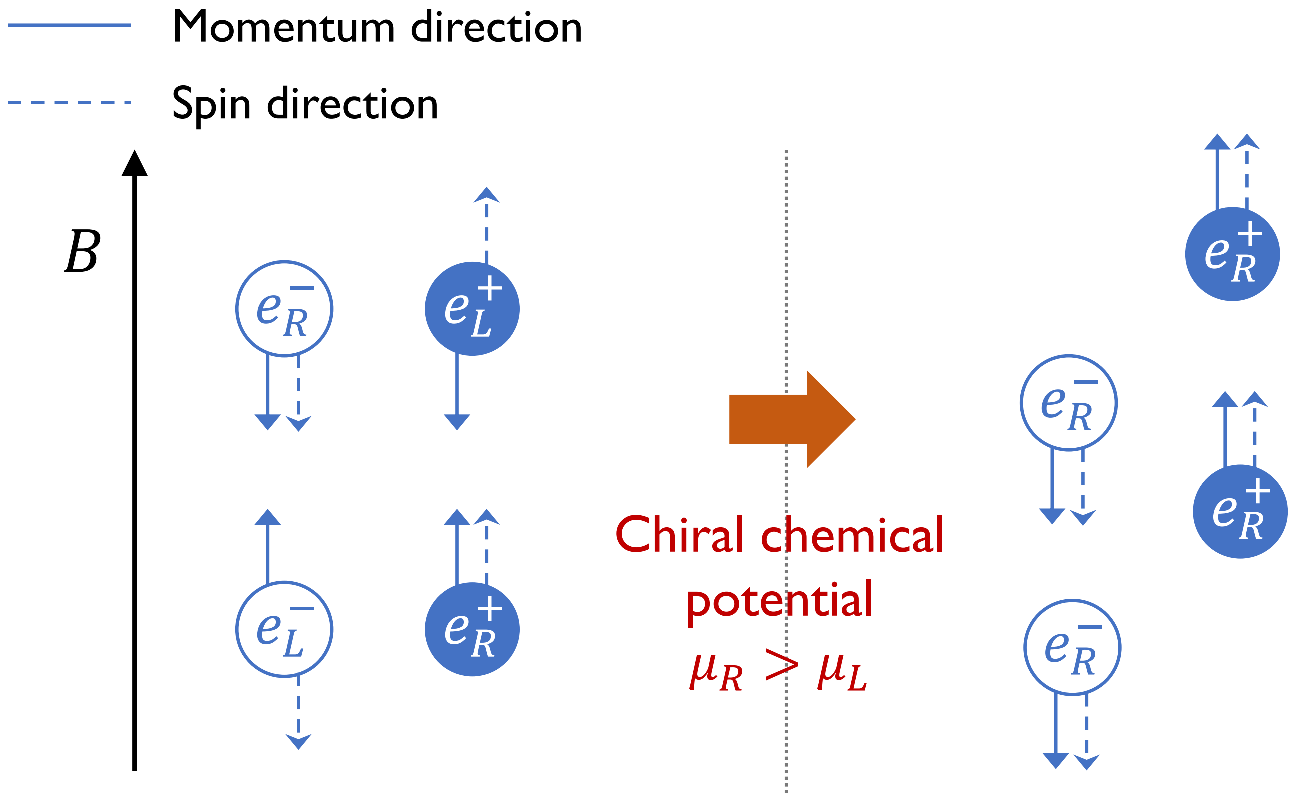
(dashed : $m_\Psi = 10^{-3} f_a$,

solid : $m_\Psi = f_a$)

Take-home message I

- In principle, we have 3 possible classes of UV physics for an axion : KSVZ-like, DFSZ-like, and string-theoretic.
- Those 3 classes of UV physics may be experimentally distinguishable by measuring the ratio of an axion-fermion coupling to the axion-photon coupling.
- For the QCD axion, *the measurement of the axion-electron coupling is crucial* for the distinction.

Chiral Magnetic Effect (CME) in a nutshell



Net current $\vec{J} = 0$

Net current $\vec{J} \propto (\mu_R - \mu_L)\,\vec{B}$

- The magnetic field aligns the spin directions depending on particles and antiparticles.
- The helicity imbalance causes a non-zero electric current along the B-field direction.

- Chiral chemical potential

$$\mathcal{L} \supset \mu_5(n_R - n_L) = \mu_5\left(\psi_R^\dagger\psi_R - \psi_L^\dagger\psi_L\right)$$

$$= \mu_5\bar{\Psi}\gamma^0\gamma^5\Psi \qquad \Psi = \begin{pmatrix}\psi_L \\ \psi_R\end{pmatrix}$$

It makes *helicity imbalance*.

- (Vector) chemical potential

$$\mathcal{L} \supset \mu(n_R + n_L) = \mu\left(\psi_R^\dagger\psi_R + \psi_L^\dagger\psi_L\right)$$

$$= \mu\bar{\Psi}\gamma^0\Psi$$

It makes *charge imbalance* (i.e. particles vs antiparticles).

The *charge imbalance alone cannot induce a current*. However, μ might be still relevant for the magnitude of the current for a given helicity imbalance.

Time-varying axion field as a source for μ_5

Suppose that we have an axion-fermion coupling:

$$c_\Psi \frac{\partial_\mu a}{f_a} \bar{\Psi} \gamma^\mu \gamma^5 \Psi$$

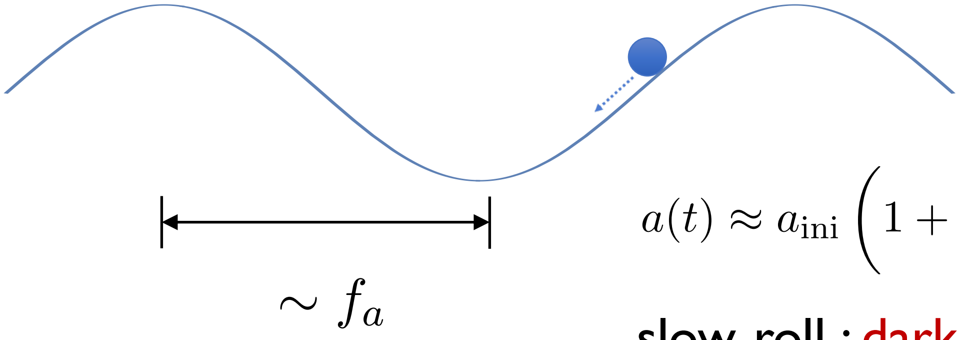
➡

$$c_\Psi \frac{\dot{a}}{f_a} \bar{\Psi} \gamma^0 \gamma^5 \Psi$$

$$\mu_5(t) = c_\Psi \frac{\dot{a}}{f_a}$$

Cosmological evolution of an axion field

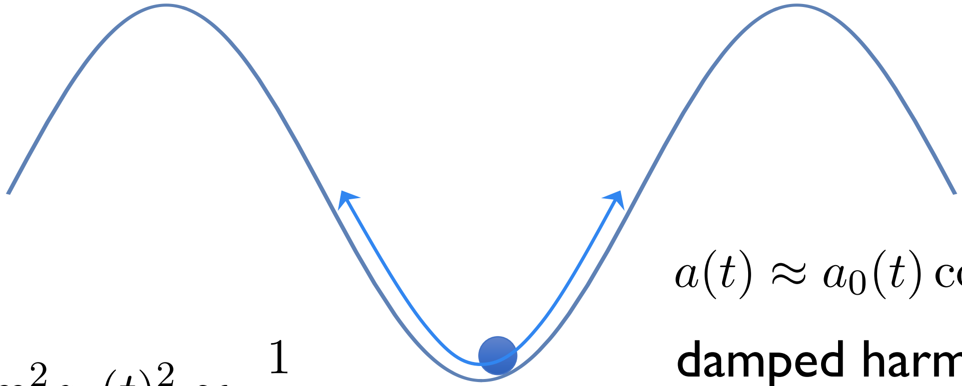
$m_a \ll H(t)$



$$a(t) \approx a_{\text{ini}} \left(1 + \frac{1}{20} \left[\frac{m_a^2}{H_{\text{ini}}^2} - \frac{m_a^2}{H(t)^2} \right] \right)$$

slow-roll : dark energy

$m_a \gg H(t)$

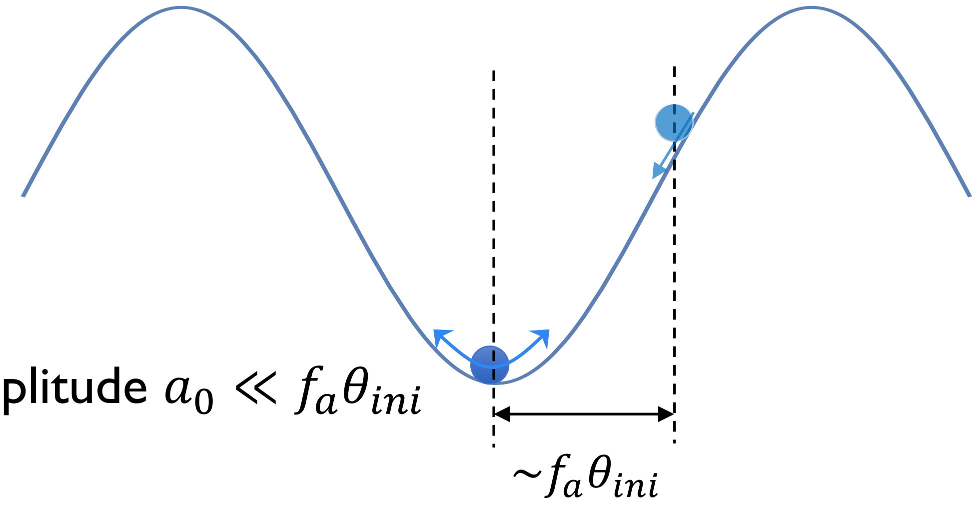


$$a(t) \approx a_0(t) \cos(m_a t)$$

damped harmonic
oscillator : cold dark matter

$$\rho_a \approx \frac{1}{2} m_a^2 a_0(t)^2 \propto \frac{1}{R^3}$$

Misalignment production of axion dark matter and chiral chemical potential



Present oscillation amplitude $a_0 \ll f_a \theta_{ini}$

$$a(t, \vec{x}) \approx a_0 \cos m_a t$$

$$\rho_a = \frac{1}{2} m_a^2 a_0^2$$

➡

$$\boxed{\mu_5(t) \approx -c_\Psi \frac{\sqrt{2\rho_a}}{f_a} \sin m_a t}$$

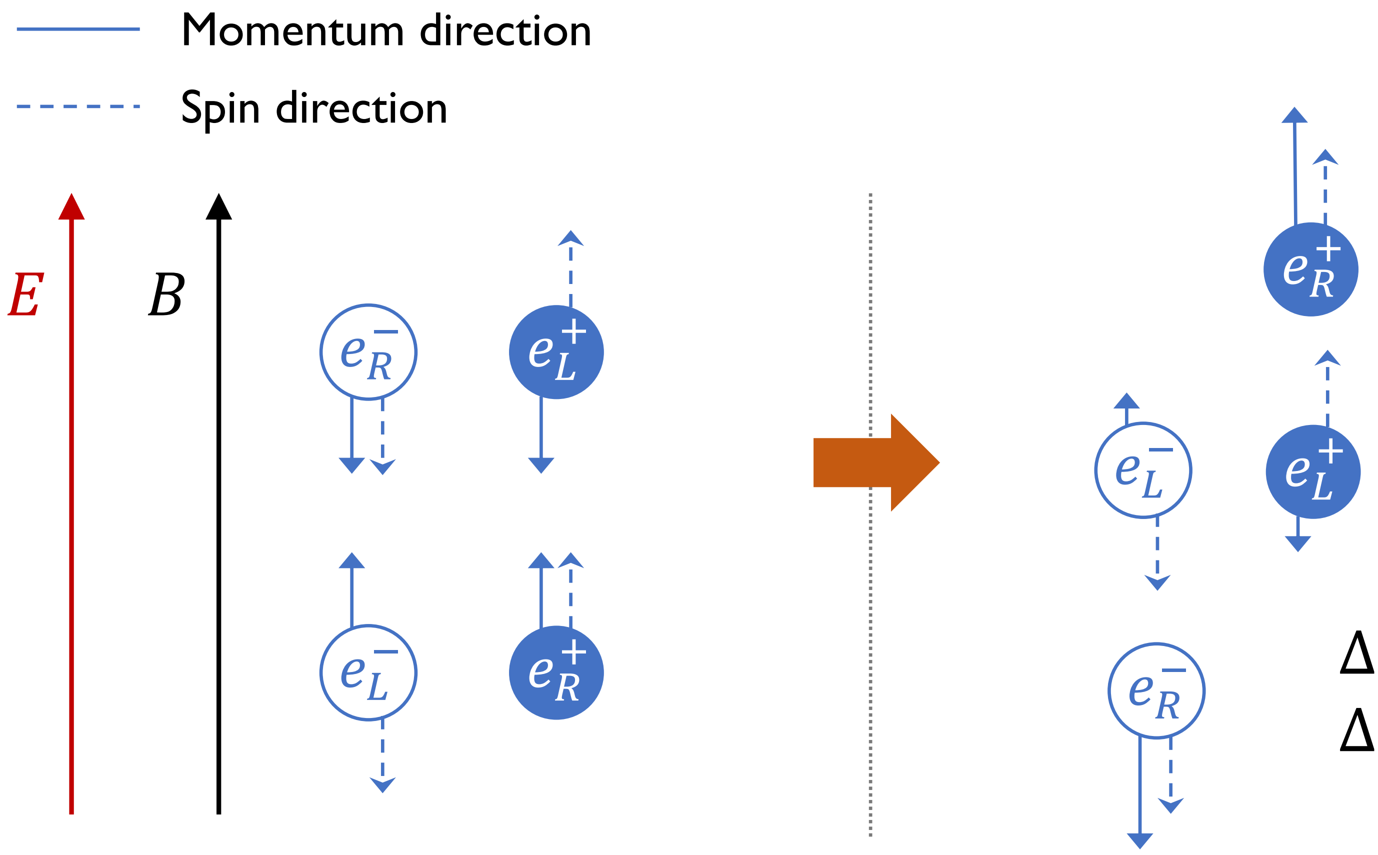
$$\mu_5(t) = c_\Psi \frac{\dot{a}}{f_a}$$

The axion dark matter background gives rise to an oscillating chiral chemical potential for fermions coupled to the axion.

CME-induced current

I. Energy balance argument

Nielsen and Ninomiya ‘83
Fukushima, Kharzeev, Warringa ‘08



$$\Delta P_F^R = eEt$$
$$\Delta P_F^L = -eEt$$

Density change of
right-handed fermion states

$$\underbrace{\frac{\Delta P_F^R}{2\pi}}_{\text{Longitudinal phase space density}} \cdot \underbrace{\frac{eB}{2\pi}}_{\text{The density of Landau levels in the transverse direction}} = \frac{e^2}{4\pi^2}\vec{E}\cdot\vec{B}t$$

Density change of
left-handed fermion states

$$\frac{\Delta P_F^L}{2\pi} \cdot \frac{eB}{2\pi} = -\frac{e^2}{4\pi^2}\vec{E}\cdot\vec{B}t$$

➡

$$\frac{d}{dt}(n_R - n_L) = \frac{e^2}{2\pi^2}\vec{E}\cdot\vec{B}$$

which partially reproduces the axial anomaly equation

$$\partial_\mu(\bar{\Psi}\gamma^\mu\gamma^5\Psi) = \frac{e^2}{2\pi^2}\vec{E}\cdot\vec{B} + 2m\bar{\Psi}\Psi$$

$$\frac{d}{dt}(n_R - n_L) = \frac{e^2}{2\pi^2}\vec{E}\cdot\vec{B}$$

➡ $$\underbrace{\mathcal{E}_{LR}\frac{d}{dt}(N_R - N_L)} = \frac{e^2}{2\pi^2}\mathcal{E}_{LR}\int d^3x\vec{E}\cdot\vec{B}$$

Energy cost per unit time needed for adding a right-handed fermion and removing a left-handed fermion

||

$$\int d^3x\,\vec{j}\cdot\vec{E}$$

This energy cost has to be supplied by the electric power.

➡ $$\boxed{\vec{j} = \frac{e^2}{2\pi^2}\mathcal{E}_{LR}\vec{B}}$$

$\mathcal{E}_{LR}$: Energy cost for converting a left mode to a right mode

Essentially

$$
\begin{aligned}
j = \rho v &= e(n_R - n_L)\, v_F \\
&= e\left(\frac{(p_F^R - p_F^L)}{2\pi} \, \frac{eB}{2\pi} \right) v_F \\
&= \frac{e^2}{2\pi^2} \cdot \underbrace{\frac{1}{2}(p_F^R - p_F^L) v_F}_{\varepsilon_{LR}} \cdot B
\end{aligned}
$$

Landau level quantization

Fermi surface ($\omega = \mu$)

$$\begin{aligned} \omega_R &= \sqrt{(|p_z| - \mu_5)^2 + m^2} \\ \omega_L &= \sqrt{(|p_z| + \mu_5)^2 + m^2} \end{aligned} \quad \Rightarrow \quad \begin{aligned} p_F^R &= \sqrt{\mu^2 - m^2} + \mu_5 \\ p_F^L &= \sqrt{\mu^2 - m^2} - \mu_5 \end{aligned}$$

$$\mathcal{E}_{LR} = \frac{1}{2}(p_F^R - p_F^L) v_F \quad \Rightarrow \quad \boxed{\mathcal{E}_{LR} = \mu_5 v_F}$$

Therefore,

$$\boxed{\vec{j} = \frac{e^2}{2\pi^2} \mu_5 v_F \vec{B}}$$

(Caveat) v_F dependence is missing in the original literature of Fukushima, Kharzeev, Warringa (0808.3382).

CME-induced current

II. Effective lagrangian approach for non-relativistic fermions

DK Hong, SHI, KS Jeong, D Yeom '22

$$\mathcal{L} = i\bar{\Psi}\gamma^{\mu}D_{\mu}\Psi - m\bar{\Psi}\Psi + \mu_5\bar{\Psi}\gamma^0\gamma^5\Psi$$

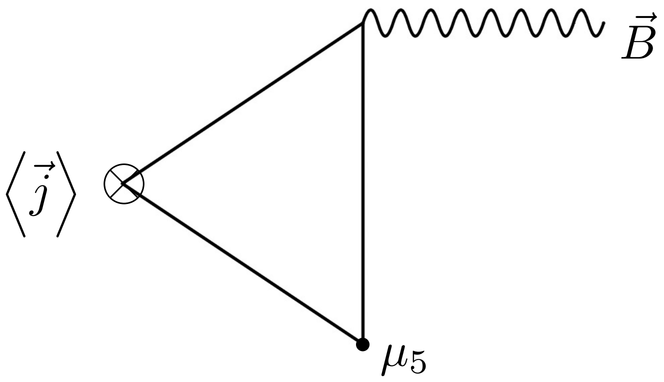
$$\Psi = e^{-imt}\tilde{\Psi} \qquad \tilde{\Psi} = \begin{pmatrix} \psi \\ \chi \end{pmatrix}$$

With the standard rep of the γ matrices, χ can be integrated out for the momentum scale below m.

➡ $$\mathcal{L}_{\text{eff}} = \psi^{\dagger}\left[i\partial_t + \frac{1}{2m}(\vec{\sigma}\cdot\vec{D})^2 + \mu_5\frac{(-i\vec{\nabla})}{m}\cdot\vec{\sigma}\right]\psi$$

(Caveat) This part was not noticed in our draft v1.

(to be corrected : thanks to K. Choi)



$$\vec{j}_L(\mu_L, B) \equiv \left\langle \frac{\delta S}{\delta \vec{A}} \right\rangle_{\vec{A}=0} = e \left\langle \Psi_L^\dagger \vec{\sigma} v_F \Psi_L \right\rangle = -ie \int_p \mathrm{Tr_L} \left(\frac{\sigma^i v_F}{(1+i\epsilon)p^0 - \frac{\vec{p}^2}{2m} + \mu_B \vec{\sigma} \cdot \vec{B}_{\mathrm{ext}} + \mu_L} \right)$$

➡
$$j_L^i(\mu_L, B) = \frac{e^2}{2\pi^2} \mu_L B \, \delta^{i3}$$

$$j_R^i(\mu_R, B) = -\frac{e^2}{2\pi^2} \mu_R B \delta^{i3}$$

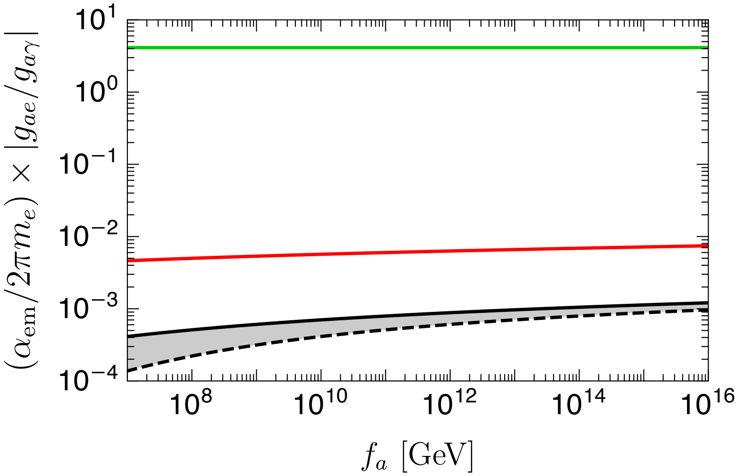
➡
$$\boxed{j^3 = j_L^3 + j_R^3 = \frac{e^2}{2\pi^2}(\mu_L - \mu_R) B}$$

$\mu_L - \mu_R = \mu_5 v_F \neq \mu_5$ **(to be corrected)**

The axion-electron coupling

$$\mathcal{L}_{\rm int} = C_e \frac{\partial_\mu a}{f} \bar{\psi}\gamma^\mu \gamma_5 \psi$$

$$C_e \simeq \begin{cases} \mathcal{O}(1) & \text{DFSZ-like models} \\ \mathcal{O}(10^{-4} \sim 10^{-3}) & \text{KSVZ-like models} \\ \mathcal{O}(10^{-3} \sim 10^{-2}) & \text{string-theoretic axions} \end{cases}$$



The measurement of the axion-electron coupling will give us an important clue for underlying high energy physics.

K Choi, SHI, HJ Kim, H Seong '21

Detecting axion dark matter via the CME

DK Hong, SHI, KS Jeong, D Yeom '22

$$C_e \frac{\partial_\mu a}{f} \bar{\psi} \gamma^\mu \gamma_5 \psi \quad \Rightarrow \quad \mu_5 = C_e \frac{\dot{a}}{f}$$

$$a(t) = \frac{\sqrt{2\rho_{\rm DM}}}{m_a} \sin(m_a t)$$

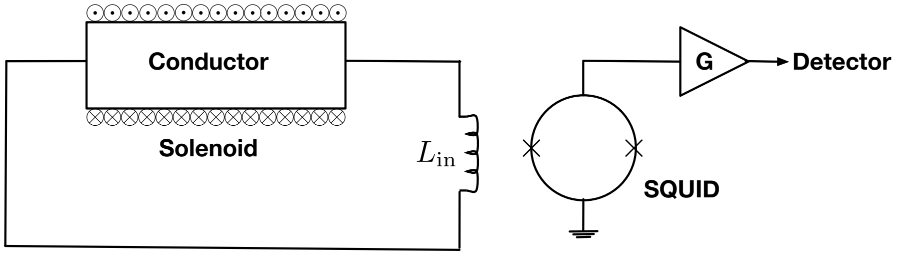
The axion dark matter field induces an oscillating chiral chemical potential if the axion couples to the electrons.

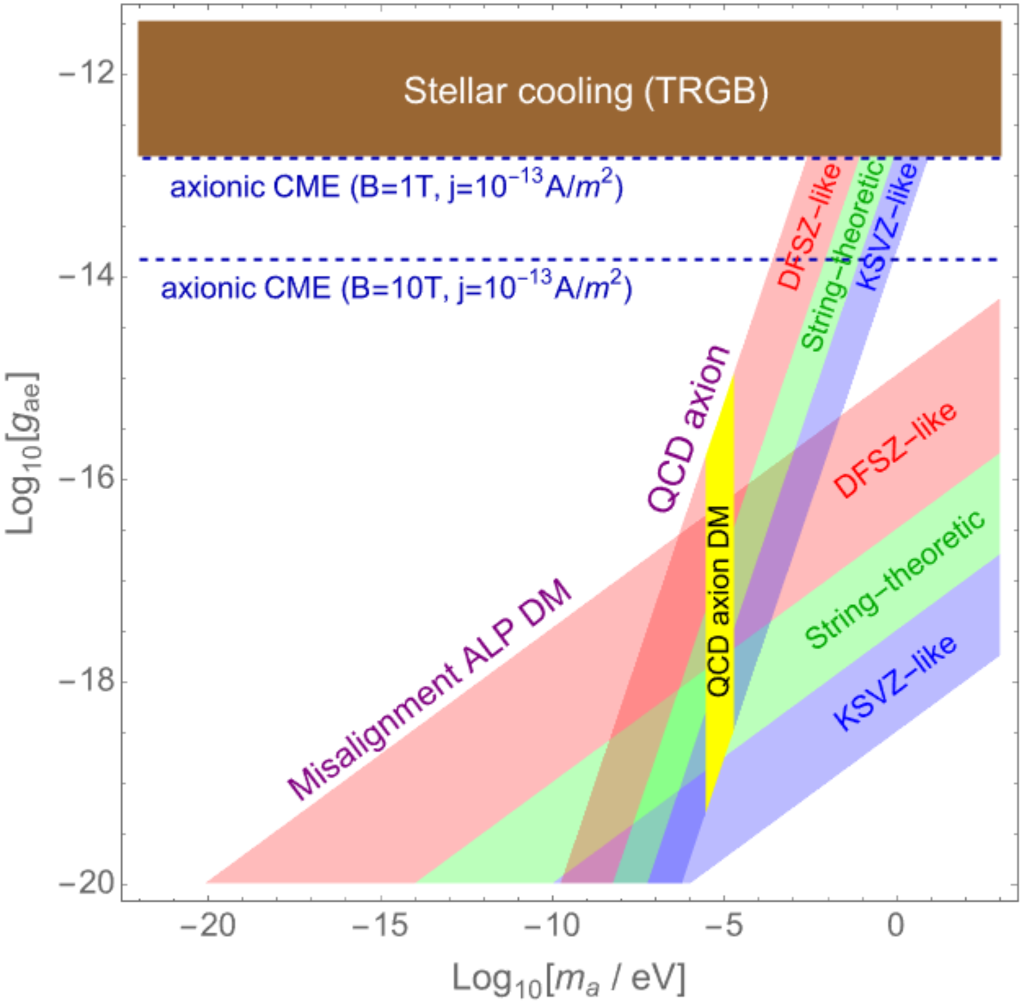
$$\mu_5 = C_e \frac{\sqrt{\rho_{\rm DM}}}{f} \cos(m_a t) \sim 0.25 \times 10^{-32}\ {\rm GeV} \cdot \left(\frac{\rho_{\rm DM}}{0.4\ {\rm GeV\,cm^{-3}}}\right)^{1/2} \cdot \left(\frac{10^{12}\ {\rm GeV}}{f/C_e}\right)$$

$$\Rightarrow \quad j^3 = \frac{e^2}{2\pi^2} \mu_5 v_F B = 6.8 \times 10^{-15}\,{\rm A\,m^{-2}} \cos(m_a t)$$

$$\times \left(\frac{v_F}{10^{-2}}\right) \cdot \left(\frac{\rho_{\rm DM}}{0.4\,{\rm GeV\,cm^{-3}}}\right)^{1/2} \cdot \left(\frac{10^{12}\,{\rm GeV}}{f/C_e}\right) \cdot \left(\frac{B}{10\,{\rm Tesla}}\right)$$

CME





Such a tiny CME-induced current may be measurable by exploiting a highly sensitive superconducting SQUID coil or a semiconductor diode.

Conclusions

- Axions are theoretically well-motivated new particles which may be an important clue for underlying UV physics when they are discovered.
- The underlying UV physics may be distinguishable by precision measurements of low energy axion couplings.
- The measurement of the axion-electron coupling is particularly important for pinning down the microscopic origin of the QCD axion.
- The chiral magnetic effect (CME) offers an intriguing possibility for the measurement of the axion-electron coupling, when the axion comprises a major fraction of dark matter.
- At the moment, the size of the induced current seems rather controversial for the case of massive fermions.

Back-up slides

String-theoretic models

$$C_{[m_1 m_2 .. m_p]}(x^\mu, y^m) = a(x^\mu)\Omega_{[m_1 m_2 .. m_p]}(y^m)$$

Ω : harmonic p-form on the compact internal space

4D axions identified as zero modes of higher-dimensional p-form gauge field

SUSY-preserving compactification

$$T = \tau + ia$$

Axion chiral superfield (τ : volume modulus of p-cycle dual to Ω)

$$U(1)_{PQ} : \quad a \rightarrow a + \mathrm{const}$$

: remnant of a higher-dimensional gauge symmetry

$$\delta C_{[m_1 m_2 .. m_p]} = \partial_{[m_1} \Lambda_{m_2,..,m_p]}$$

4D Low energy effective action

$$K = K_0(T + T^*) + Z_I(T + T^*)\Phi_I^* \Phi_I$$

$$\mathcal{F}_A = c_A T \qquad Z_I \propto (T + T^*)^{\omega_I} \quad \omega_I \sim \mathcal{O}(1)$$

scaling weight of Φ_I

$$c_A \sim \mathcal{O}(1)$$

Conlon, Cremades, Quevedo '06

$$K = K_0(T+T^*) + Z_I(T+T^*)\Phi_I^*\Phi_I$$

$$\mathcal{F}_A = c_A T \qquad\qquad Z_I \propto (T+T^*)^{\omega_I} \quad \omega_I \sim \mathcal{O}(1) \quad c_A \sim \mathcal{O}(1)$$

$$T = \tau + ia$$

$$\mathcal{L}_{\rm eff} = \frac{M_P^2}{4}\partial_\tau^2 K_0\,(\partial_\mu a)^2 + \frac{\omega_I}{2\tau}\partial_\mu a\left(\psi_I^\dagger\bar\sigma^\mu\psi_I + \phi_I^\dagger i\overset{\leftrightarrow}{D}{}^\mu\phi_I\right) - \frac{1}{4}c_A\tau F^{A\mu\nu}F^A_{\mu\nu} + \frac{c_A}{4}aF^{A\mu\nu}\widetilde{F}^A_{\mu\nu}$$

$\sim O(1)$ $\sim O(1)$

$$\tau = \frac{1}{c_A g_A^2} \sim \mathcal{O}(1)$$

String-theoretic axion couplings to matter fields and gauge fields are comparable to each other.

Canonical normalization

$$a \to \frac{a}{8\pi^2 f_a} \qquad f_a = \frac{M_P}{8\pi^2}\sqrt{\frac{\partial_\tau^2 K_0}{2}}$$

$$\mathcal{L}_{\rm eff} = -\frac{1}{4g_A^2}F^{A\mu\nu}F^A_{\mu\nu} + \frac{1}{2}(\partial_\mu a)^2 + \omega_I\frac{c_A g_A^2}{16\pi^2}\frac{\partial_\mu a}{f_a}\left(\psi_I^\dagger\bar\sigma^\mu\psi_I + \phi_I^\dagger i\overset{\leftrightarrow}{D}{}^\mu\phi_I\right) + \frac{c_A}{32\pi^2}\frac{a}{f_a}F^{A\mu\nu}\widetilde{F}^A_{\mu\nu}$$

$\sim O(g^2/16\pi^2)$

Laboratory searches for axion DM -photonic probes

$$\frac{g_{a\gamma}}{4} a F \tilde{F}$$ ➡ $$\nabla \times \vec{B} = \frac{\partial \vec{E}}{\partial t} \underbrace{- g_{a\gamma} \vec{B} \partial_t a}_{\vec{J}_{\rm eff}}$$

effective current

Background axion DM field

$$a \approx a_0 \cos\left[m_a (t - \vec{v}\cdot\vec{x})\right]$$

$$\rho_a = \frac{1}{2} m_a^2 a_0^2 \qquad |\vec{v}| \sim 10^{-3} c$$

➡ $$\vec{J}_{\rm eff} \approx g_{a\gamma} \sqrt{2\rho_a} \vec{B} \sin m_a t$$

The best experimental sensitivity on $g_{a\gamma}$ is obtained when $\rho_a = \rho_{DM}$.

Misalignment axion DM

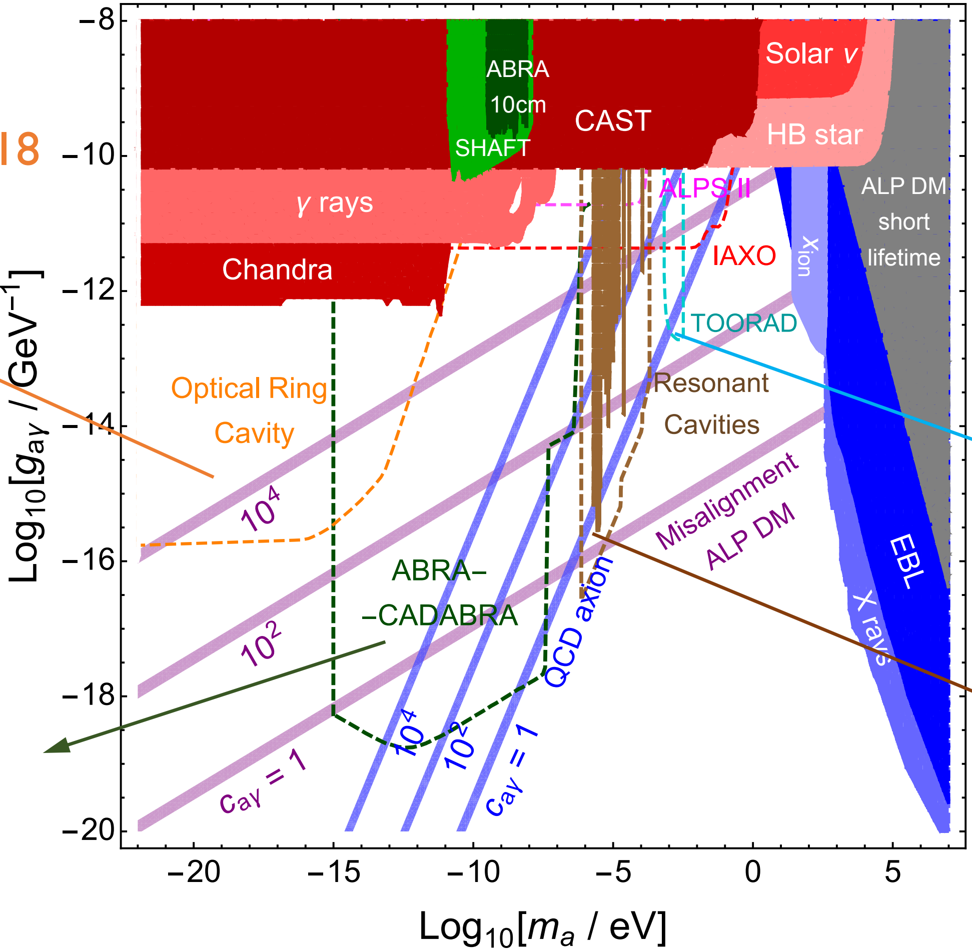
$$f_a \simeq 10^{17}\,{\rm GeV} \left(\frac{10^{-22}\,{\rm eV}}{m_a}\right)^{1/4} \sqrt{\frac{\rho_a}{\rho_{\rm DM}}}$$

➡ $$g_{a\gamma} = \frac{e^2}{8\pi^2} \frac{1}{f_a} c_{a\gamma}$$

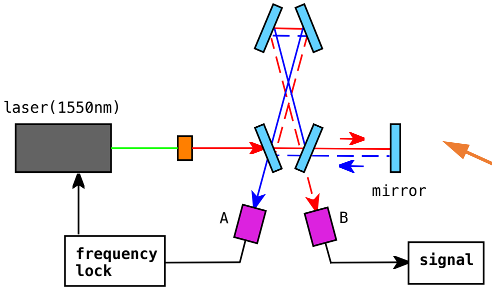
Given axion DM mass,
$g_{a\gamma}$ is determined for $c_{a\gamma} \sim O(1)$.

Current and future limits on $g_{a\gamma}$

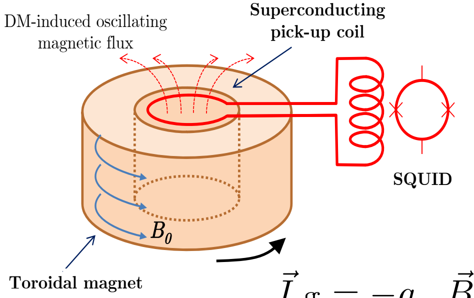
Choi, SHI, Shin '20



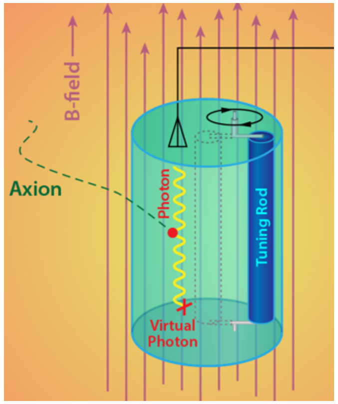
Obata, Fujita, Michimura '18



Kahn, Safdi, Thaler '16



$$\vec{J}_{\rm eff} = -g_{a\gamma}\vec{B}_0\partial_t a$$



Marsh, Fong, Lentz, Smejkal, Ali '18

ADMX, IBS-CAPP, MADMAX…

Laboratory searches for axion DM -nucleonic probes

$$g_{aN}\frac{\partial_\mu a}{2m_N}\bar{N}\gamma^\mu\gamma^5 N \quad \Rightarrow \quad \underbrace{g_{aN}\frac{\nabla a}{\gamma_N m_N}}_{\vec{B}_{\rm eff}}\cdot\gamma_N\vec{S}_N$$

γ_N : nucleon gyromagnetic ratio

Background axion DM field

$$a \approx a_0\cos\left[m_a(t-\vec{v}\cdot\vec{x})\right] \quad \Rightarrow \quad \vec{B}_{\rm eff}\approx g_{aN}\frac{\sqrt{2\rho_a}}{\gamma_N m_N}\vec{v}_a\sin m_a t$$

$$\rho_a = \frac{1}{2}m_a^2 a_0^2 \qquad |\vec{v}|\sim 10^{-3}c$$

The best experimental sensitivity on g_{aN} is obtained when $\rho_a = \rho_{DM}$.

Misalignment axion DM

$$f_a \simeq 10^{17}\,{\rm GeV}\left(\frac{10^{-22}\,{\rm eV}}{m_a}\right)^{1/4}\sqrt{\frac{\rho_a}{\rho_{\rm DM}}} \quad \Rightarrow \quad g_{aN} = \frac{m_N}{f_a}c_{aq}\times\mathcal{O}(1)$$

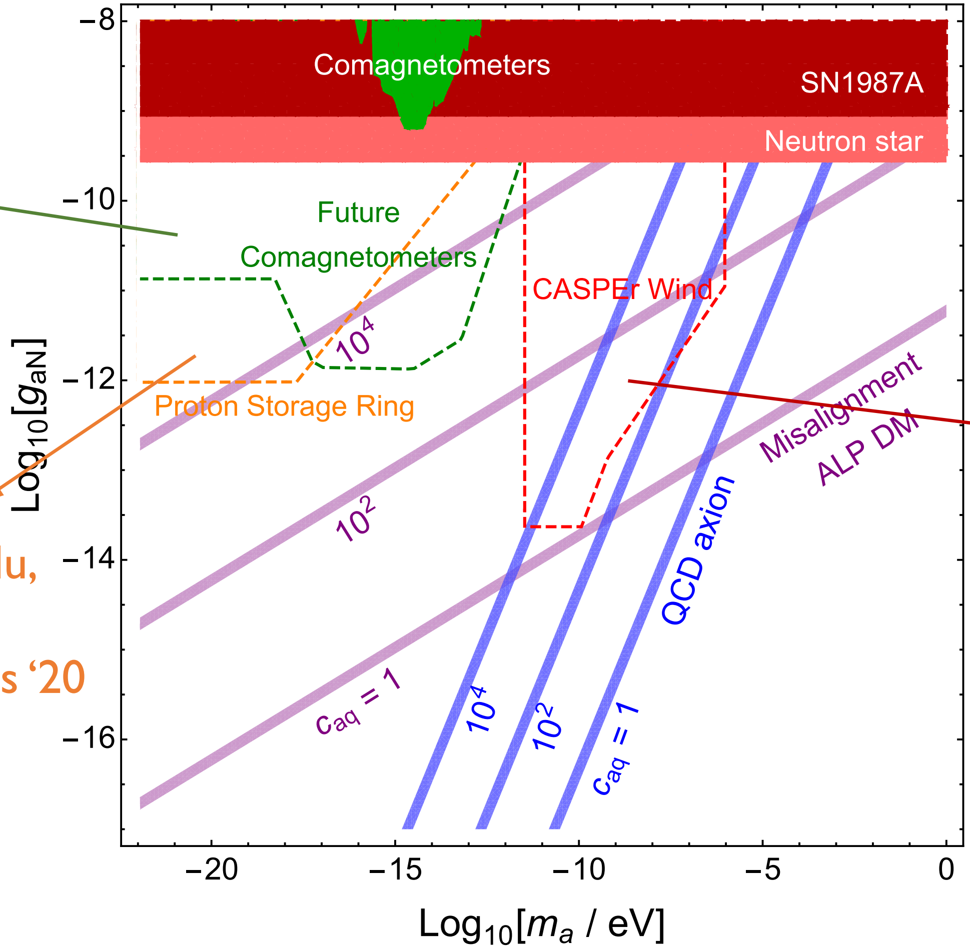
Given axion DM mass,
g_{aN} is determined for $c_{aq}\sim O(1)$.

Current and future limits on g_{aN}

Choi, SHI, Shin ’20

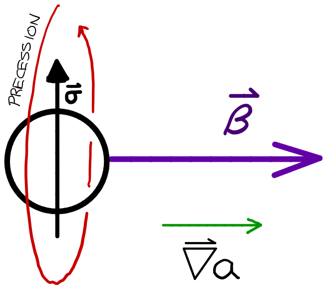
Bloch, Hochberg, Kuflik, Volansky ‘19

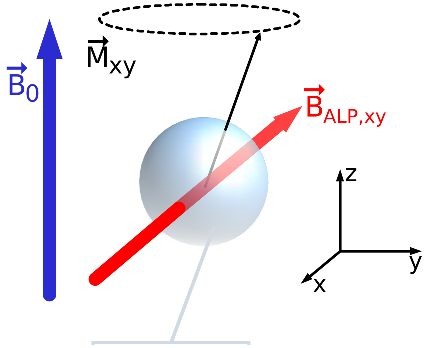
$$\frac{B^e_{\text{eff}}}{B^N_{\text{eff}}} \sim \frac{c_{ae} m_e}{c_{aN} m_N} \neq 1$$



Kimball et al ‘17

Graham, Haciomeroglu, Kaplan, Omarov, Rajendran, Semertzidis ‘20





NMR

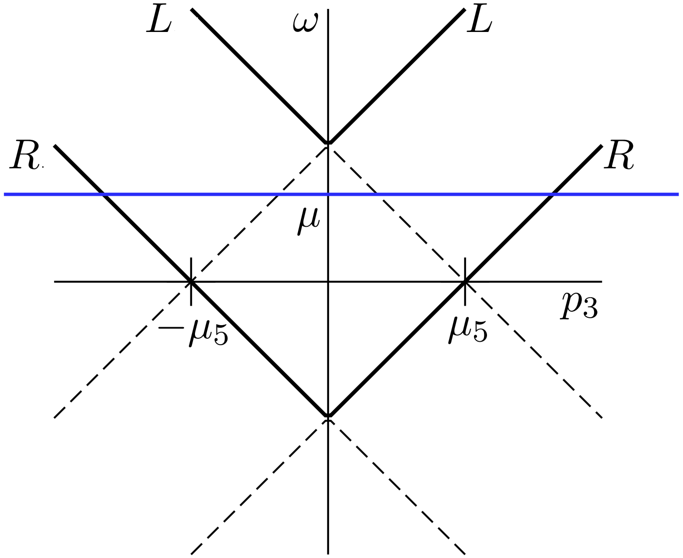
Dirac equation for a massless fermion

$$\begin{pmatrix} 0 & \omega - p_z\sigma^3 + \mu_5 \\ \omega + p_z\sigma^3 - \mu_5 & 0 \end{pmatrix} \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix} = 0$$

➡

$$\omega_R = |p_z| - \mu_5$$

$$\omega_L = |p_z| + \mu_5$$



Fermi surface

$$n_i = \frac{1}{e^{\beta(\omega_i - \mu)} + 1}$$

At low temperatures and low μ, only the right chiral modes are occupied for a positive μ_5.

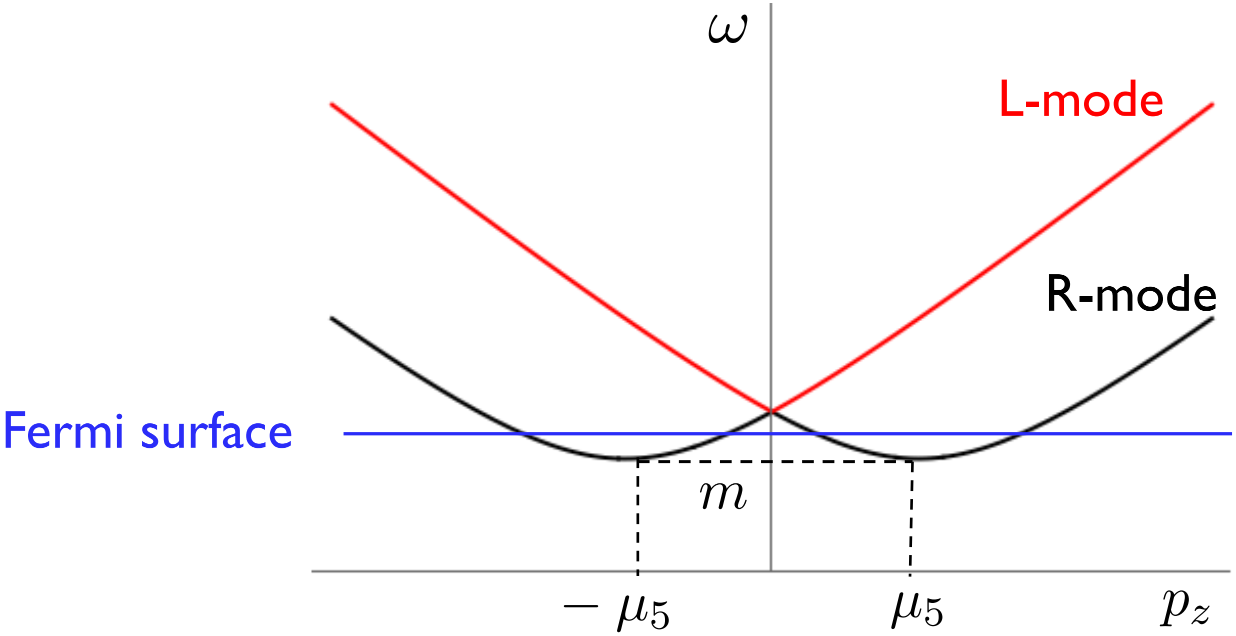
→ helicity imbalance

Dirac equation for a massive fermion

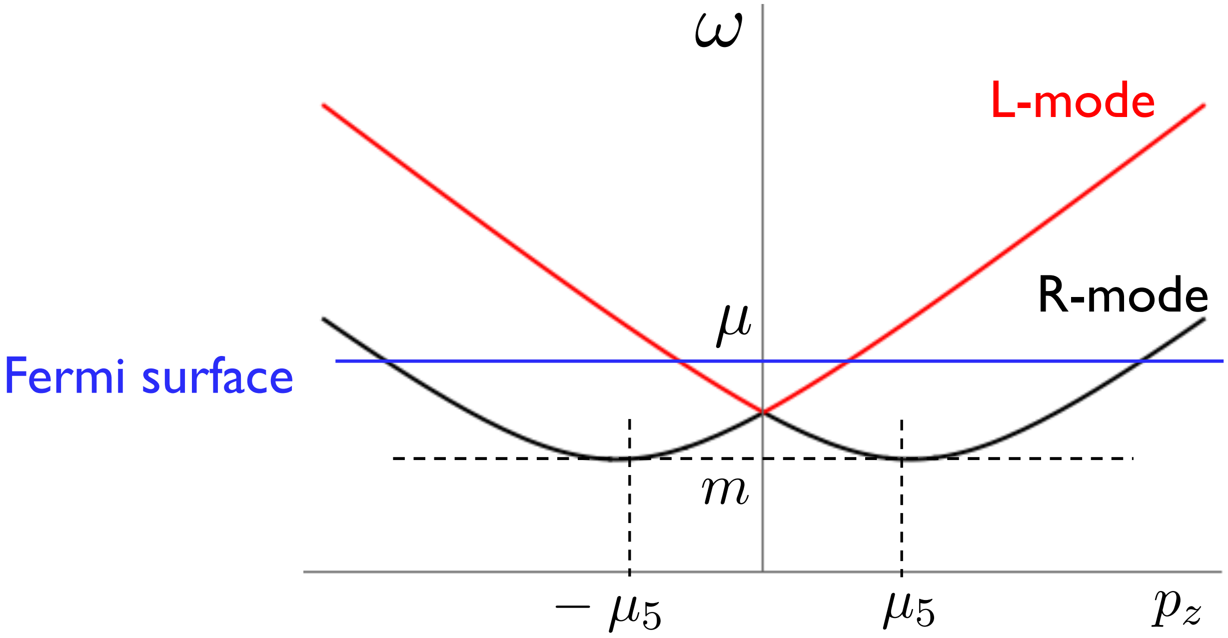
$$\begin{pmatrix} -m & \omega - p_z\sigma^3 + \mu_5 \\ \omega + p_z\sigma^3 - \mu_5 & -m \end{pmatrix} \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix} = 0$$

$$\omega_R = \sqrt{(|p_z| - \mu_5)^2 + m^2}$$
$$\omega_L = \sqrt{(|p_z| + \mu_5)^2 + m^2}$$

Helicity eigenstates



Qualitatively the same as the massless case



- The difference between the L-mode and the R-mode disappears for $\mu_5 = 0$.
- Without a magnetic field, the spin directions of particles will not be fixed.
- The occupation number difference depends on the chemical potential μ and the fermion mass m.

$$J(\mu_5, B, \mu, m, T) = \mu_5 B \times f(\mu, m, T)$$

+ *possible higher order terms* in (μ_5, B)