# Rotating fermions on anti-de Sitter space

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VEA, EW, PLB **749** (2015) 597; CQG **34** (2017) 145010; Symmetry **2021**, 13, 2019.

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# Outline

# Motivation

Vortical effects relevant for:

- Rotating black holes (frame dragging) — M. Casals *et al.*, PRD **87** (2013) 064027
- Neutron stars (anomalous transport) — M. Kaminski *et al.*, PLB **760** (2016) 170
- QGP (spin polarisation) — STAR Collaboration, Nature **548** (2017) 62; VEA, M.N. Chernodub, EPJC **82** (2022) 61

Why AdS?

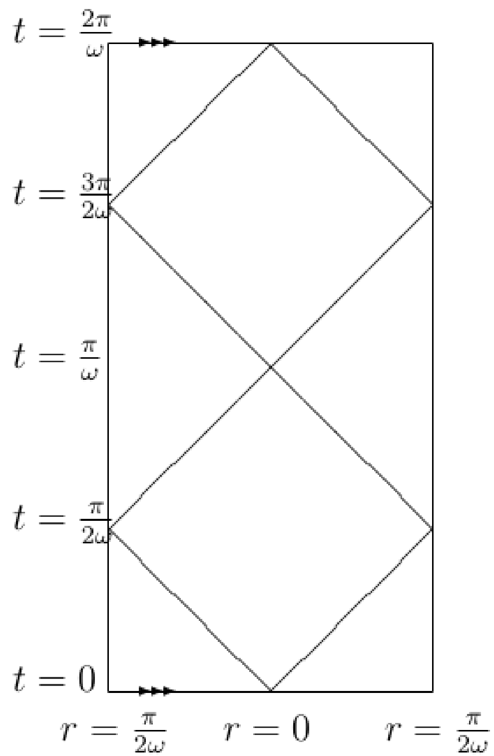
- AdS/CFT duality — Aharony, Gubser, Maldacena, Ooguri, Oz, Phys. Rep. **323** (2000) 183; Son, Starinets, JHEP**03** (2006) 052.
- AdS has timelike boundary $\rightarrow$ no SLS for "mild" $\Omega$. — R. Panerai, PRD **93** (2016) 104021
- Maximal symmetry $\Rightarrow$ analytically tractable.

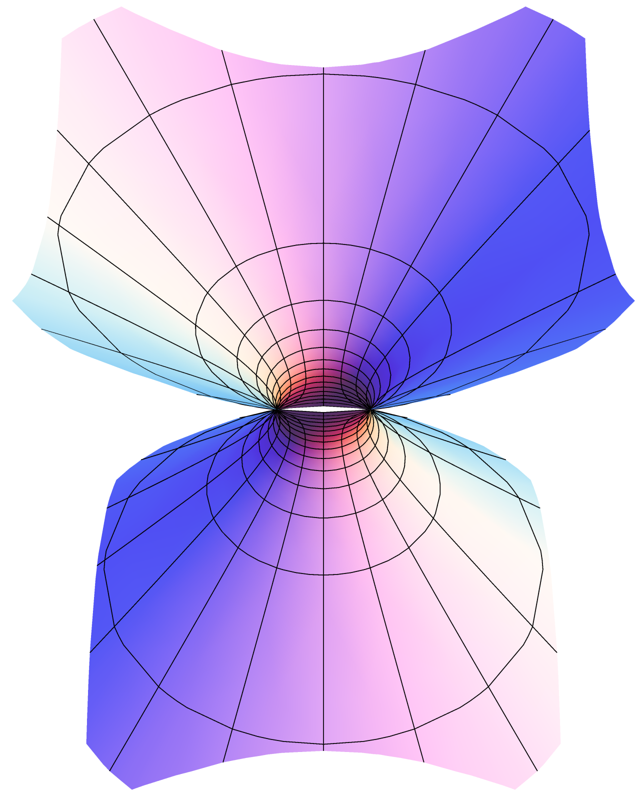
# Anti-de Sitter space

- Maximal symmetry:
  $R_{\mu\nu} = \frac{1}{4} R g_{\mu\nu}$,
  $R = 4\Lambda = -12\ell^{-2}$.

- Boundary at $r = \pi\ell/2$ reached in a finite time only by null geodesics.

- Closed timelike loops
  $\Rightarrow$ CAdS ($-\infty < t < \infty$).



- Line element:

$$ds^2 = \frac{\ell^2}{\cos^2 \overline{r}} \left[ -d\overline{t}^2 + d\overline{r}^2 + \sin^2 \overline{r} \left( d\theta^2 + \sin^2 \theta d\varphi^2 \right) \right],$$

  where $\overline{t} = t/\ell$, $\overline{r} = r/\ell$.

# Vacuum propagator

- For the maximally symmetric vacuum state, $S_F$ can be written as:

W. Mück, J. Phys. A **33** (2000) 2000 3021

$$iS^F_{\rm vac}(x,x') = [\mathcal{A}(s) + \mathcal{B}(s)\not{n}]\Lambda(x,x').$$

- The geodesic interval $s = \ell\overline{s}$ can be given through:

$$\cos\overline{s} = \frac{\cos\Delta\overline{t}}{\cos\overline{r}\cos\omega\overline{r}'} - \cos\gamma\tan\overline{r}\tan\overline{r}',$$

where $\cos\gamma = \cos\theta\cos\theta' + \sin\theta\sin\theta'\cos\Delta\varphi$.

- $n_\mu = \nabla_\mu s(x,x')$ is the normalised tangent to the geodesic at $x$.
- $\mathcal{A}$ and $\mathcal{B}$ depend only on $\overline{s}$ and satisfy:

$$i\frac{d}{d\overline{s}}\begin{pmatrix}\mathcal{A}\\ \mathcal{B}\end{pmatrix} + \frac{3i}{2}\begin{pmatrix}-\mathcal{A}\tan(\overline{s}/2)\\ \mathcal{B}\cot(\overline{s}/2)\end{pmatrix} - \ell M\begin{pmatrix}\mathcal{B}\\ \mathcal{A}\end{pmatrix} = \begin{pmatrix}0\\ i(-g)^{-1/2}\delta(x,x')\end{pmatrix}.$$

# Vacuum propagator

- The equations can be solved exactly:

$$\mathcal{A}_F = \frac{\Gamma_k}{16\pi^2\ell^3}\frac{\cos(\overline{s}/2)}{[-\sin^2(\overline{s}/2)]^{2+k}}\,{}_2F_1\left(1+k, 2+k; 1+2k; \mathrm{cosec}^2\left(\frac{\overline{s}}{2}\right)\right),$$
$$\mathcal{B}_F = \frac{i\Gamma_k}{16\pi^2\ell^3}\frac{\sin(\overline{s}/2)}{[-\sin^2(\overline{s}/2)]^{2+k}}\,{}_2F_1\left(k, 2+k; 1+2k; \mathrm{cosec}^2\left(\frac{\overline{s}}{2}\right)\right),$$

where $k = \ell m$, while the normalisation constant $\Gamma_k$ is given by

$$\Gamma_k = \frac{\Gamma\left(2+k\right)\Gamma\left(\frac{1}{2}\right)}{4^k\Gamma\left(\frac{1}{2}+k\right)}.$$

- In the limit $k \to 0$, we have $\Gamma_k \to 1$ and

$$\lim_{k\to 0}\mathcal{A}_F = \frac{1}{16\pi^2\ell^3}\left(\cos\frac{\overline{s}}{2}\right)^{-3}, \qquad \lim_{k\to 0}\mathcal{B}_F = \frac{i}{16\pi^2\ell^3}\left(\sin\frac{\overline{s}}{2}\right)^{-3}.$$

# Bi-spinor of parallel transport

- The bi-spinor of parallel transport $\Lambda(x,x')$ satisfies: W. Mück, J. Phys. A **33** (2000) 3021

$$D_\mu \Lambda(x,x') = -i\omega S_{\mu\nu} n^\nu \Lambda(x,x') \tan\left(\frac{\overline{s}}{2}\right).$$

- Employing the Cartesian gauge for the tetrad: I. I. Cotăescu, Rom. J. Phys. **52** (2007) 895

$$e_{\hat{t}} = \ell^{-1} \cos\overline{r}\, \partial_{\overline{t}}, \qquad e_{\hat{\imath}} = \cos\overline{r} \left[\frac{\overline{r}}{\sin\overline{r}}\left(\delta_{ij} - \frac{x^i x^j}{r^2}\right) + \frac{x^i x^j}{r^2}\right] \partial_j, \tag{1}$$

allows $\Lambda(x,x')$ to be expressed as: VEA, EW, CQG **34** (2017) 145010.

$$\Lambda(x,x') = \frac{\sec(\overline{s}/2)}{\sqrt{\cos\overline{r}\cos\overline{r}'}} \left[\cos\frac{\Delta\overline{t}}{2}\left(\cos\frac{\overline{r}}{2}\cos\frac{\overline{r}'}{2} + \sin\frac{\overline{r}}{2}\sin\frac{\overline{r}'}{2}\frac{\boldsymbol{x}\cdot\boldsymbol{\gamma}}{r}\frac{\boldsymbol{x}'\cdot\boldsymbol{\gamma}}{r'}\right)\right.$$
$$\left. + \sin\frac{\Delta\overline{t}}{2}\left(\sin\frac{\overline{r}}{2}\cos\frac{\overline{r}'}{2}\frac{\boldsymbol{x}\cdot\boldsymbol{\gamma}}{r}\gamma^{\hat{t}} + \sin\frac{\overline{r}'}{2}\cos\frac{\overline{r}}{2}\frac{\boldsymbol{x}'\cdot\boldsymbol{\gamma}}{r'}\gamma^{\hat{t}}\right)\right]. \tag{2}$$

# Schwinger–de Witt regularisation

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- Using the ansatz $S_F = (\alpha + \beta \not{n})\Lambda$, the v.e.v. of $T_{\mu\nu}$ can be computed as

$$\langle 0|T_{\mu\nu}|0\rangle = 8i \lim_{x'\to x} \left[g_{\mu\nu} \frac{\omega}{2} \cot \frac{\omega s}{2} \beta - n_\mu n_\nu \left(\frac{\partial}{\partial s} - \frac{\omega}{2} \cot \frac{\omega s}{2}\right) \beta\right].$$

- Due to maximal symmetry, $T_{\mu\nu} \sim g_{\mu\nu}$.

$$\langle 0|T_{\mu\nu}|0\rangle = \frac{i}{2} \lim_{x'\to x} \text{tr} \left\{\gamma_{(\nu}(x) D_{\mu)}(x) S_F(x,x') - S_F(x,x') \overleftarrow{D_{(\mu}}(x') \gamma_{\nu)}(x')\right\}.$$

- The renormalisation of $T_{\mu\nu}$ is performed by subtracting renormalisation counter-terms:

S. M. Christensen, Phys. Rev. D **17**, 946–963 (1978)

$$T_{\text{ren;SdW}} = -\frac{\omega^4}{4\pi^2}\left\{\tfrac{11}{60} + k - \tfrac{k^2}{6} - k^3 + 2k^2(k^2-1)\left[\ln \tfrac{\nu}{\omega} - \psi(k) + \ln 2\right]\right\}.$$

- The result is agrees with the Pauli–Villars regularisation method.

R. Camporesi and A. Higuchi, Phys. Rev. D **45**, 3591–3603 (1992)

# Hadamard regularisation of $S_F$

- Hadamard regularisation uses the auxiliary function $\mathcal{G}_F$:

$$S_F(x,x') = (i\gamma^\lambda D_\lambda + \mu)\mathcal{G}_F(x,x').$$

- The Hadamard theorem states that:

$$\mathcal{G}_F(x,x') = \frac{1}{8\pi^2}\Big(\underbrace{\frac{u}{\sigma} + v\ln|\nu^2\sigma|}_{\text{Hadamard form } \mathcal{G}_H} + w\Big)\Lambda(x,x'),$$

where $\sigma = -\frac{1}{2}s^2$; $u$, $v$ and $w$ are finite; and $\nu$ is a renormalisation mass scale.

- $u$ and $v$ can be determined from the Dirac equation:

$$u = \left(\frac{\omega s}{\sin\omega s}\right)^{\frac{3}{2}}, \qquad v = \frac{\mu^2-\omega^2}{2}\cos\frac{\omega s}{2}\, {}_2F_1\left(2-k, 2+k; 2; \sin^2\frac{\omega s}{2}\right),$$

leading to

VEA, EW, PLB **749** (2015) 597

$$T_{\text{ren}} = -\frac{\omega^4}{4\pi^2}\left\{\frac{11}{60} + k - \frac{7k^2}{6} - k^3 + \frac{3k^4}{2} + 2k^2(k^2-1)\left[\ln\frac{\widetilde{\nu}}{\omega} - \psi(k)\right]\right\},$$

where $k = \mu/\omega$, $\widetilde{\nu} = \nu e^{-\gamma}\sqrt{2}$ and $\gamma$ is Euler's constant.

# Rigid rotation in spherically-symmetric space-times

- SSST can be parametrised using central charts: VEA, I. I. Cotăescu, PRD **94 (2016) 085022**

$$ds^2 = W^2 \left[ -dt^2 + \frac{dr^2}{U^2} + \frac{r^2}{V^2}(d\theta^2 + \sin^2\theta d\varphi^2) \right], \tag{3}$$

- $W = U = V = 1$ for Mink, while on adS,

$$W = \frac{1}{\cos \overline{r}}, \qquad U = 1, \qquad V = \frac{r}{\sin \overline{r}}. \tag{4}$$

- The relativistic Boltzmann equation, $k^\mu \partial_\mu f = C[f]$, is satisfied when

$$f = f^{(\text{eq})}(-\beta \cdot k - \alpha), \qquad \nabla_\mu \beta_\nu + \nabla_\nu \beta_\mu = 0, \qquad \nabla_\mu \alpha = 0, \tag{5}$$

where $\beta^\mu = T^{-1} u^\mu$ and $\alpha = \mu/T$ ($= 0$ in this talk).

- Rigid rotation $\beta = \beta_0(\partial_t + \Omega \partial_\varphi)$ is a solution of the Killing equation, with

$$u = \frac{\Gamma}{W(r)}(\partial_t + \Omega \partial_\varphi), \qquad \begin{pmatrix} T \\ \mu \end{pmatrix} = \Gamma \begin{pmatrix} T_0 \\ \mu_0 \end{pmatrix}, \qquad \Gamma = (1 - \rho^2 \Omega^2)^{-1/2}, \tag{6}$$

where $\rho = r \sin\theta / V$ is $r \sin\theta$ on Minkowski and $\ell \sin \overline{r} \sin\theta$ on adS.

Becattini, Grossi, PRD **92** (2015) 045037

Velocity :

$$u = \Gamma(\partial_t + \Omega\partial_\varphi),$$

Acceleration :

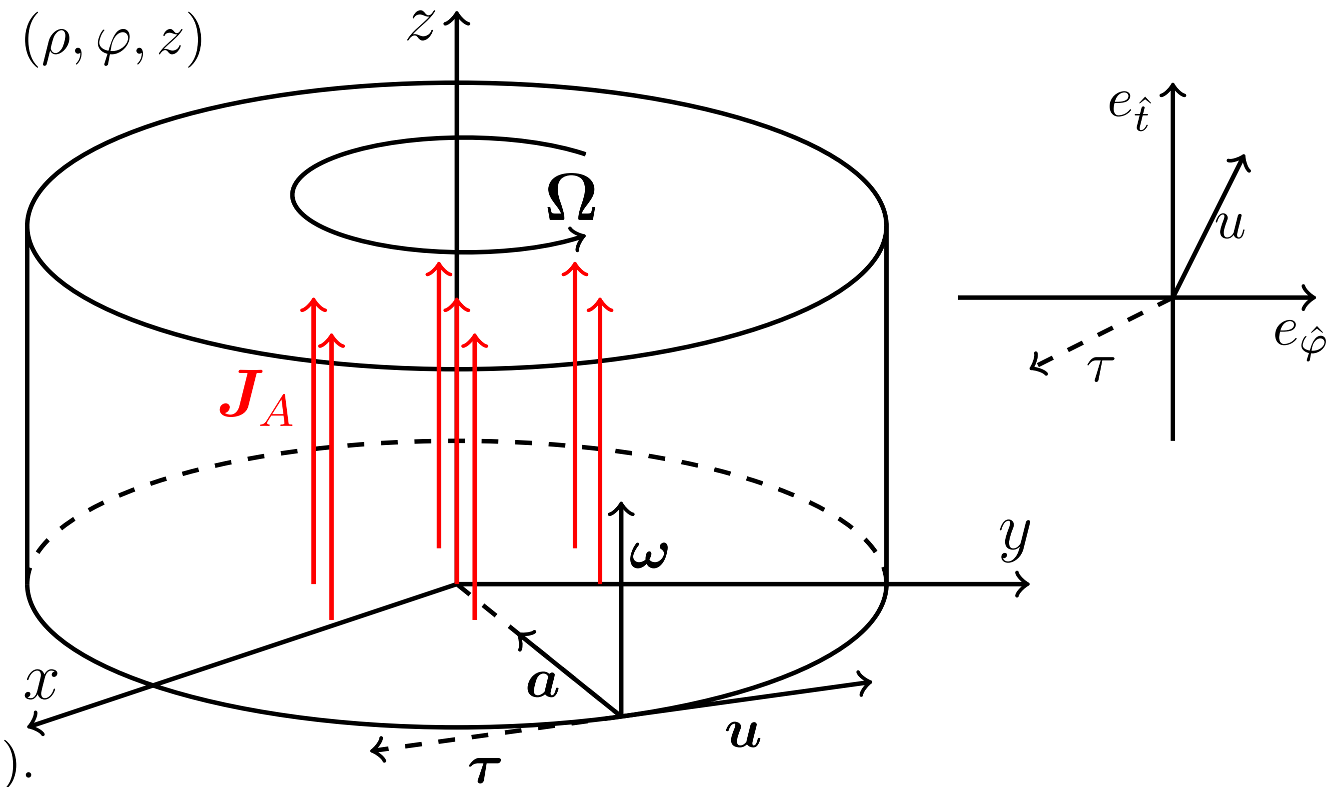
$$a = \nabla_u u = -\rho\Omega^2\Gamma^2\partial_\rho,$$

Vorticity :

$$\omega = \frac{1}{2}\varepsilon^{\hat\alpha\hat\beta\hat\gamma\hat\sigma} e_{\hat\alpha} u_{\hat\beta} (\nabla_{\hat\gamma} u_{\hat\sigma}) = \Gamma^2\Omega\partial_z,$$

Fourth vector :

$$\tau = -\varepsilon^{\hat\alpha\hat\beta\hat\gamma\hat\sigma} e_{\hat\alpha} \omega_{\hat\beta} a_{\hat\gamma} u_{\hat\sigma} = -\Omega^3\Gamma^5(\rho^2\Omega\partial_t + \partial_\varphi).$$




Axial vortical effect:

Kharzeev, Liao, Voloshin, Wang, PPNP **88** (2016) 1

$$J_A^\mu = \sigma_A^\omega \omega^\mu, \qquad \sigma_A^\omega = \frac{T^2}{6} + \frac{1}{24\pi^2}\left(\boldsymbol{\omega}^2 + 3\boldsymbol{a}^2 - 6M^2\right) + O(M^4). \qquad (7)$$

# Kinematic tetrad for adS.

Velocity : $u = \Gamma(e_{\hat{t}} + \rho\Omega e_{\hat{\varphi}})$,

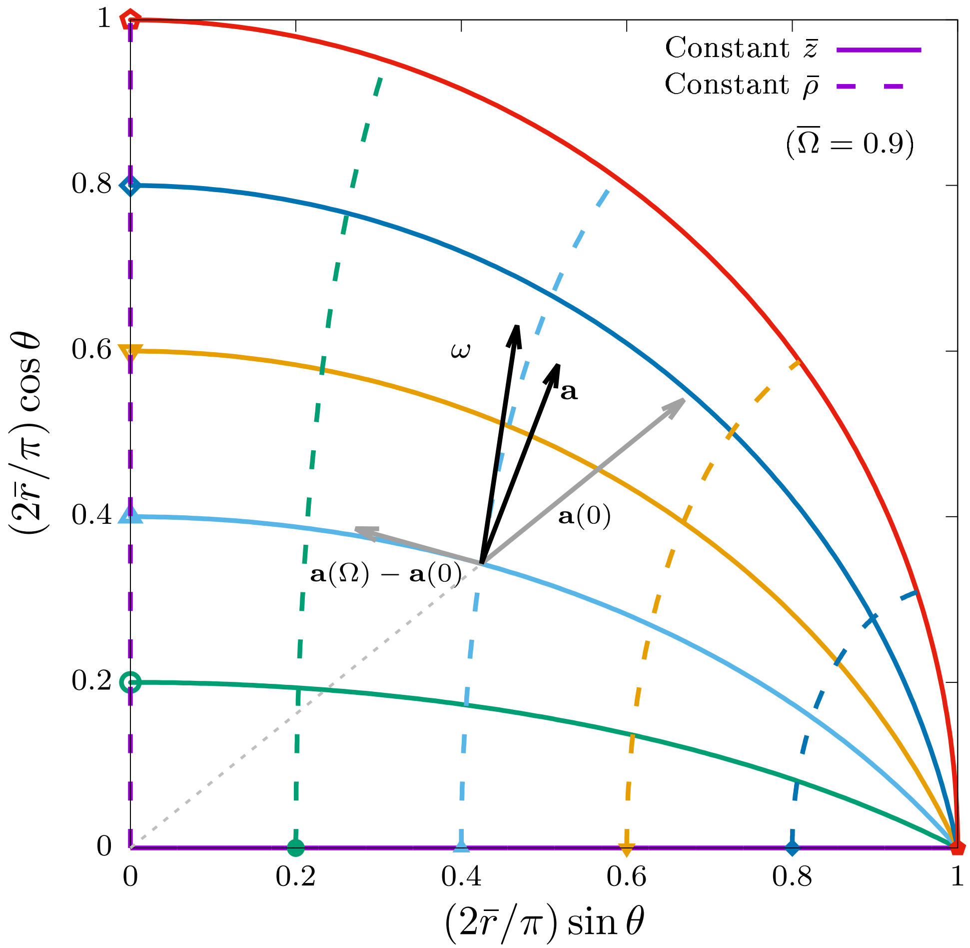
Acceleration : $a = a(0) + [a(\Omega) - a(0)]$,

$$a(0) = \frac{\sin \overline{r} \cos \overline{r}}{\ell^2} \partial_{\overline{r}},$$

$$a(\Omega) - a(0) = -\rho\Omega^2\Gamma^2 \frac{1-\overline{\rho}^2}{\sec^2 \overline{r}} \partial_\rho,$$

Vorticity : $\omega = \Omega\Gamma^2(1-\overline{\rho}^2)\partial_z$,

Fourth vector : $\tau = -\rho\Omega(1-\overline{\Omega}^2)\Gamma^5 \cos^3 \overline{r} \times (\rho\Omega\partial_t + \rho^{-1}\partial_\varphi)$.




Axial vortical effect correction:

$$J_A^\mu = \sigma_A^\omega \omega^\mu, \qquad \sigma_A^\omega = \frac{T^2}{6} + \frac{1}{24\pi^2}\left(\boldsymbol{\omega}^2 + 3\boldsymbol{a}^2 - 6M^2 + \frac{R}{4}\right) + O(M^4). \qquad (8)$$

# Geometric approach: point-splitting

- The t.e.v. of an operator $\widehat{A}$ is

$$\langle\widehat{A}\rangle_{\beta_0,\Omega} = Z^{-1}\mathrm{tr}(\hat{\varrho}\widehat{A}), \qquad \hat{\varrho} = e^{-\beta_0(\widehat{H}-\Omega\widehat{M}^z)}, \qquad Z = \mathrm{tr}(\hat{\varrho}). \tag{9}$$

- T.e.v.s can be obtained via point-splitting, e.g.

$$\langle\widehat{J}^\mu_A\rangle_{\beta_0,\Omega} = \lim_{x'\to x}\mathrm{tr}[\gamma^\mu\gamma^5 S^F_{\beta_0,\Omega}(x,x')\Lambda(x',x)], \tag{10}$$

where $\Lambda(x,x')$ is the bispinor of parallel transport.

- The Feynman propagator can be obtained from the Wightman functions,

$$\begin{aligned} S^F_{\beta_0,\Omega}(x,x') =& \Theta_c(\tau-\tau')S^+_{\beta_0,\Omega}(x,x') + \Theta_c(\tau'-\tau)S^-_{\beta_0,\Omega}(x,x'),\\ iS^+_{\beta_0,\Omega}(x,x') =& \langle\widehat{\Psi}(x)\widehat{\overline{\Psi}}(x')\rangle_{\beta_0,\Omega},\\ iS^-_{\beta_0,\Omega}(x,x') =& -\langle\widehat{\overline{\Psi}}(x')\widehat{\Psi}(x)\rangle_{\beta_0,\Omega}, \end{aligned} \tag{11}$$

where $\Theta_c(\tau-\tau')$ satisfies ($\varepsilon > 0$):

Mallik, Sarkar, *Hadrons at finite temperature* (CUP, 2016)

$$\Theta_c(t-t'-i\varepsilon) = 1, \qquad \Theta_c(t-t'+i\varepsilon) = 0. \tag{12}$$

# Wightman functions and KMS relation

- Taking into account $\hat{\varrho} = e^{-\beta_0(\widehat{H} - \Omega \widehat{M}^z)}$, $\widehat{\Psi}$ obeys

$$\hat{\varrho}\widehat{\Psi}(t,\varphi)\hat{\varrho}^{-1} = e^{-\beta_0 \Omega S^z}\widehat{\Psi}(t + i\beta_0, \varphi + i\beta_0\Omega). \tag{13}$$

- The KMS relation reads

VEA, EW, Symmetry **2021**, 13, 2019

$$\begin{aligned} S^{-}_{\beta_0,\Omega}(\tau,\varphi;x') =& iZ^{-1}\mathrm{tr}[\hat{\varrho}\widehat{\overline{\Psi}}(\tau',\varphi')\widehat{\Psi}(\tau,\varphi)] \\ =& iZ^{-1}\mathrm{tr}[\hat{\varrho}e^{\beta_0\Omega S^z}\widehat{\Psi}(\tau - i\beta_0, \varphi - i\beta_0\Omega)\widehat{\overline{\Psi}}(\tau',\varphi')] \\ =& -e^{\beta_0\Omega S^z} S^{+}_{\beta_0,\Omega}(\tau - i\beta_0, \varphi - i\beta_0\Omega; x'). \end{aligned} \tag{14}$$

- The KMS relation above allows $S^F_{\beta_0,\Omega}$ to be written in terms of $S^F_{\infty,\Omega}$:

$$\boxed{S^F_{\beta_0,\Omega}(x,x') = \sum_j (-1)^j e^{-j\beta_0\Omega S^z} S^F_{\infty,\Omega}(t + ij\beta_0, \varphi + ij\beta_0\Omega; x').} \tag{15}$$

- In bounded systems, $S^F_{\infty,\Omega} = S^F_{\infty,0}$ if no SLS forms inside the boundary.

VEA, EW, PRD **93** (2016) 104014

- A similar argument holds on adS, if $\Omega < \ell^{-1}$ (subcritical rotation).

# Thermal expectation values

- Introducing the notation

$$\zeta_j = -\frac{1}{\sin^2 \frac{\overline{s}_j}{2}} = \frac{\cos^2 \overline{r}}{\sinh^2 \frac{j\beta_0}{2\ell} - \overline{\rho}^2 \sinh^2 \frac{\Omega j \beta_0}{2}}, \tag{16}$$

we find $J^{\hat{\alpha}}_A = \sigma^\omega_A \omega^{\hat{\alpha}}$, where $[\Gamma_k = \Gamma(2+k)\sqrt{\pi}/4^k\Gamma(\frac{1}{2}+k)]$

$$\begin{aligned}\sigma^\omega_A =& \frac{\Gamma_k}{2\pi^2\ell^3\Omega\Gamma^2\cos^2\overline{r}} \sum_{j=1}^\infty (-1)^{j+1}\zeta_j^{2+k} \sinh\frac{j\beta_0}{2\ell} \sinh\frac{\Omega j\beta_0}{2} \\ & \times {}_2F_1(k, 2+k; 1+2k; -\zeta_j) \\ =& \frac{T^2}{6} + \frac{1}{24\pi^2}\left(\boldsymbol{\omega}^2 + 3\boldsymbol{a}^2 - 6M^2 + \frac{R}{4}\right) + O(T^{-1}).\end{aligned} \tag{17}$$

- At high temperature, the Minkowski expression is reproduced, including the temperature-independent terms.
- When $T \to 0$, $\sigma^\omega_A \to 0$ and the $O(T^0)$ are also suppressed.
- The curvature correction appears in the $O(T^0)$ term and is proportional to $R = -12\ell^{-2}$.

# Critical rotation: $\overline{\Omega} = 1$

- At critical rotation, $\Gamma = 1/\sqrt{1-\overline{\rho}^2}$ and

$$\zeta_j = \frac{\Gamma^2 \cos^2 \overline{r}}{\sinh^2 \frac{j\beta_0}{2\ell}}. \tag{18}$$

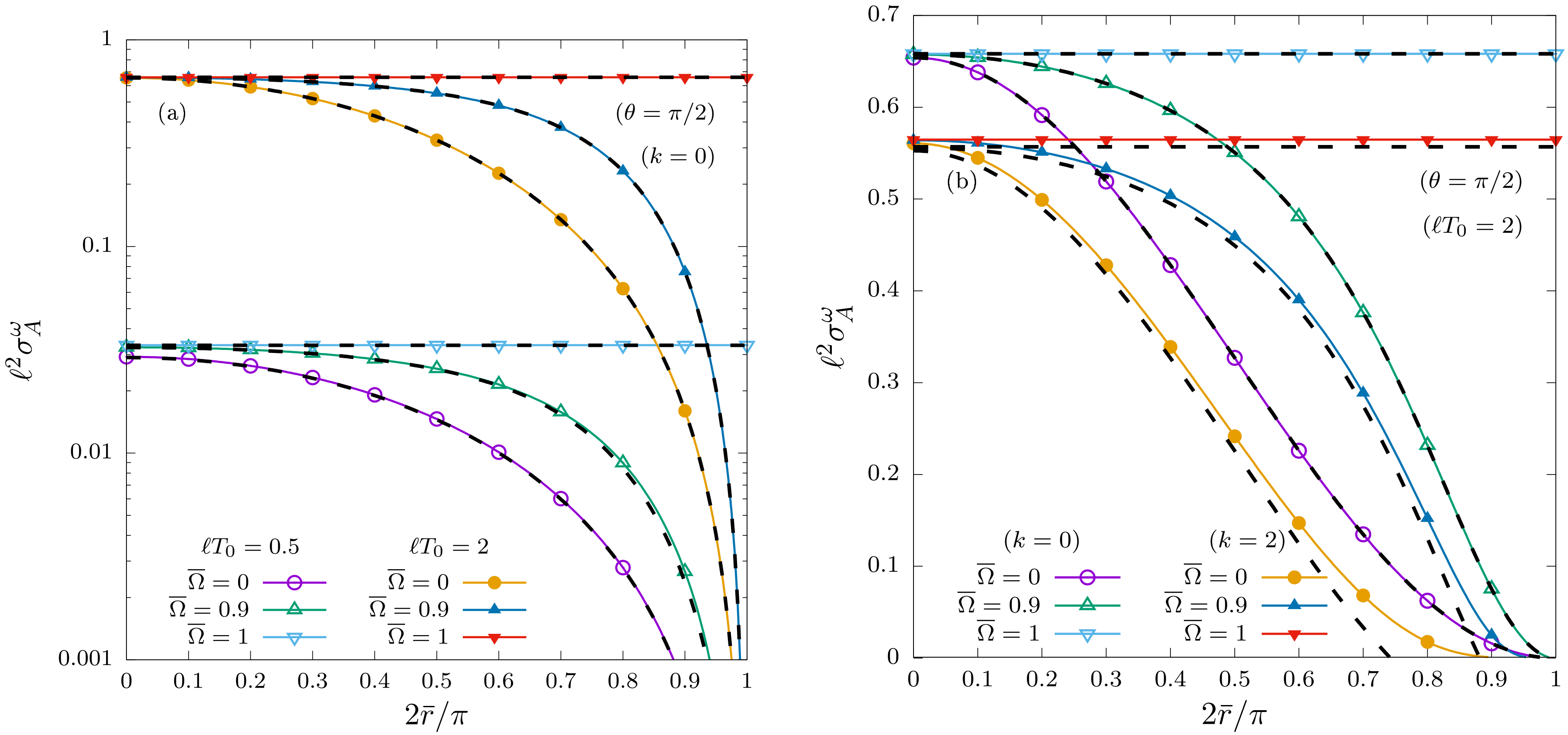
- In the equatorial plane ($\theta = \frac{\pi}{2}$),

$$\Gamma = \frac{1}{\cos \overline{r}} \qquad \Rightarrow \qquad T = T_0 \Gamma \cos \overline{r} = T_0, \qquad \zeta_j = \mathrm{cosech}^2 \frac{j\beta_0}{2\ell}. \tag{19}$$

- When $\overline{\Omega} = 1$, rotation balances "clumping" in the equatorial plane.
- In particular, $\sigma_A^\omega$ is constant in the EP:

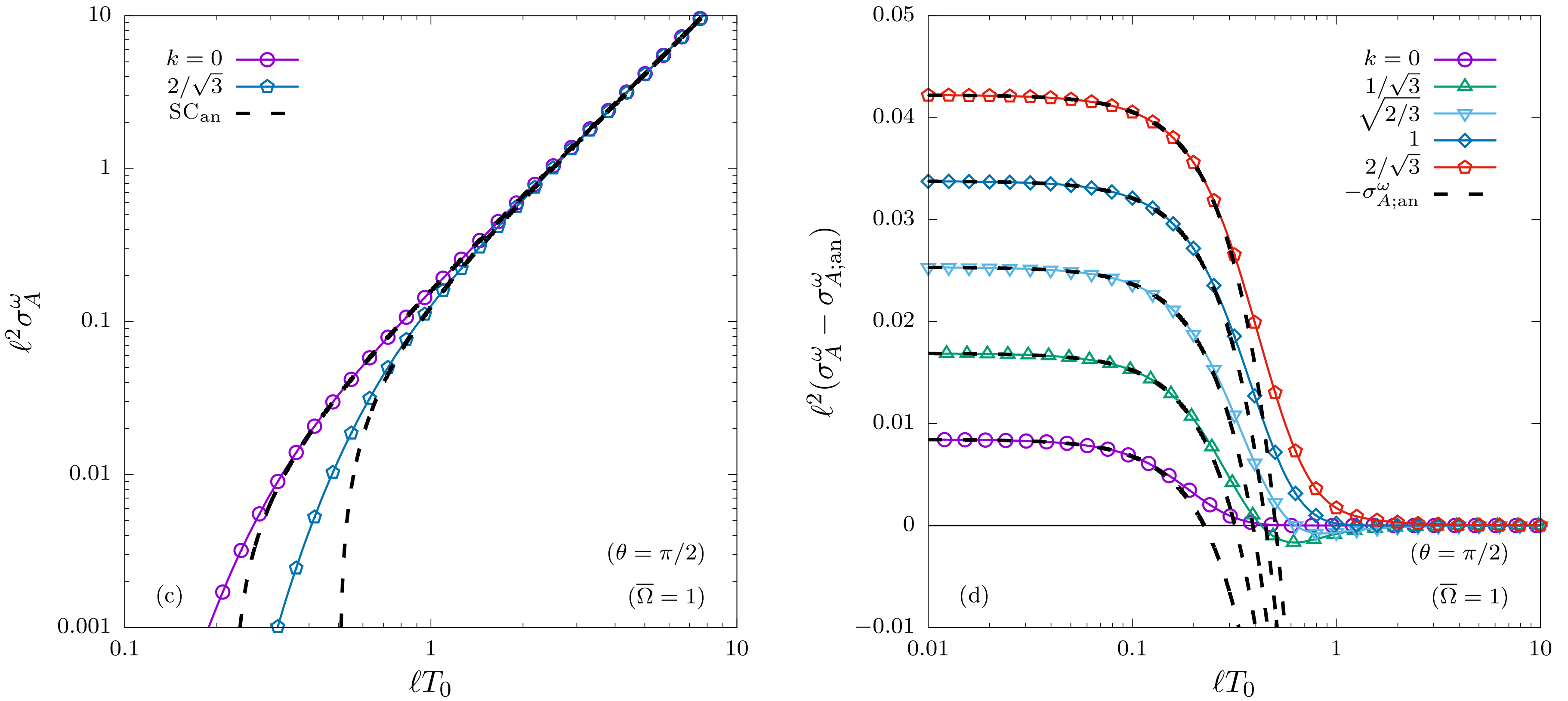
$$\lim_{\overline{\Omega}\to 1} \sigma_A^\omega \left(\theta = \frac{\pi}{2}\right) = \frac{\Gamma_k}{2\pi^2\ell^2} \sum_{j=1}^\infty \frac{(-1)^{j+1}}{(\sinh \frac{j\beta_0}{2\ell})^{2+2k}} \times {}_2F_1\left(k, 2+k; 1+2k; -\mathrm{cosech}^2 \frac{j\beta_0}{2\ell}\right). \tag{20}$$

# Equatorial plane (EP) profiles




- $T = T_0 \Gamma \cos \overline{r}$ decreases with $\overline{r}$ slower at larger $\overline{\Omega}$.
- Deviations from asymptotic formula appear at low $T_0$, high $k = \ell M$ and/or high $\overline{r}$.
- $\sigma_A^\omega$ in EP independent of $\overline{r}$ when $\overline{\Omega} = 1$ (critical rotation).

# Critical rotation $\overline{\Omega} = 1$: temperature dependence




- Asympyotic behaviour reached later for higher $k$.
- All terms including $O(T^0)$ confirmed by comparison with numerical results.

# Axial flux

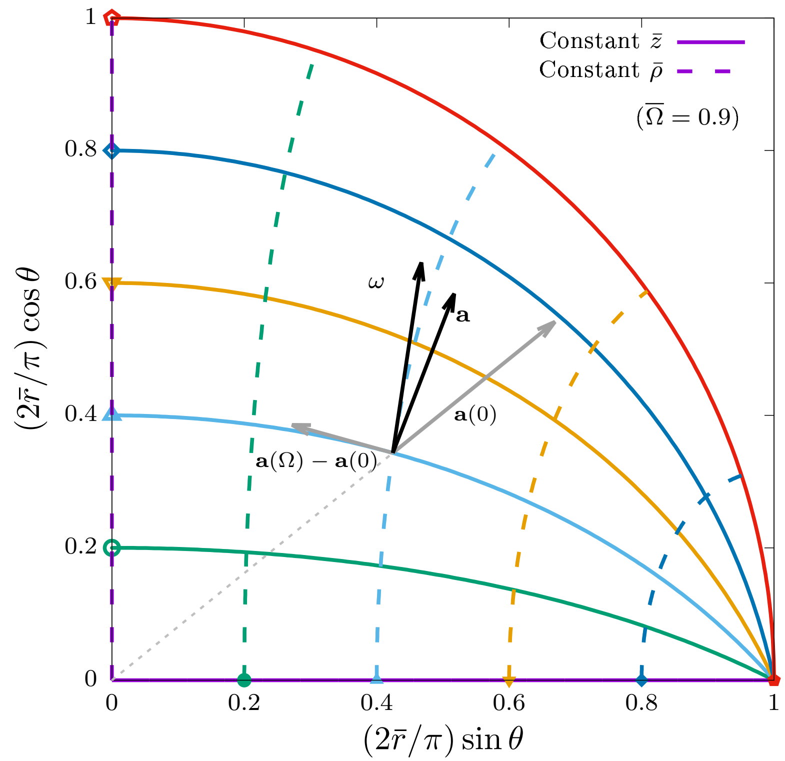
- $J_A^\mu$ obeys the conservation equation

$$\nabla_\mu J_A^\mu = \frac{1}{\sqrt{-g}} \frac{\partial(J_A^{\overline{z}}\sqrt{-g})}{\partial \overline{z}} = -2MPC,$$

where $PC = -i\overline{\psi}\gamma^5\psi$ and

$$\overline{z} = \tan\overline{r}\cos\theta, \qquad \overline{\rho} = \sin\overline{r}\sin\theta.$$

- AVE $\Rightarrow$ non-vanishing axial flux $F_A$:




$$F_A = \ell^{-1} \int_0^1 d\overline{\rho} \int_0^{2\pi} d\varphi \sqrt{-g} J_A^{\overline{z}}(\overline{z})$$

$$= \frac{\Gamma_k}{2\pi\ell(1+k)} \sum_{j=1}^{\infty} \frac{(-1)^{j+1} \sinh\frac{\Omega j\beta_0}{2} \left(\sinh\frac{j\beta_0}{2\ell}\right)^{-1-2k}}{\left(\sinh^2\frac{j\beta_0}{2\ell} - \sinh^2\frac{\Omega j\beta_0}{2}\right)(1+\overline{z}^2)^k}$$

$$\times {}_2F_1\left(k, 1+k; 1+2k; -\frac{(1+\overline{z}^2)^{-1}}{\sinh^2\frac{j\beta_0}{2\ell}}\right). \quad (21)$$

# Transparency to chiral particles

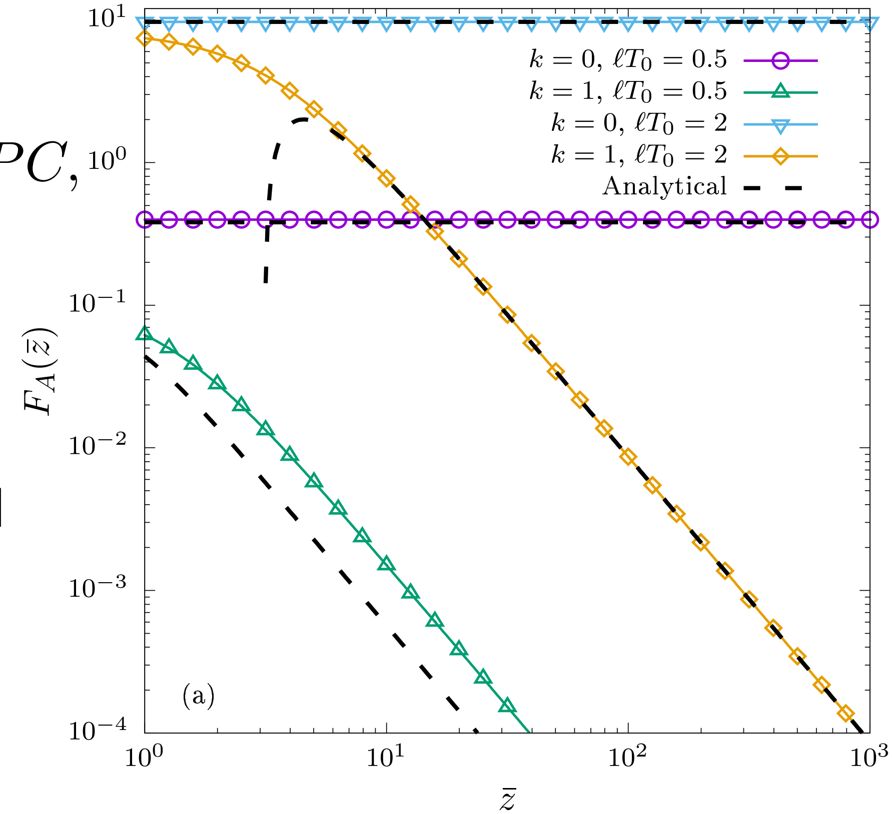
CRC-TR 211

- The conservation equation entails

$$F_A(\overline{z}) - F_A(0) = -\frac{2M}{\ell}\int_V d^3x\sqrt{-g}\,PC,$$

showing that $F_A(\pm\infty) = F_A(0)$ when $k = 0 \Rightarrow$ adS is transparent to chiral particles.

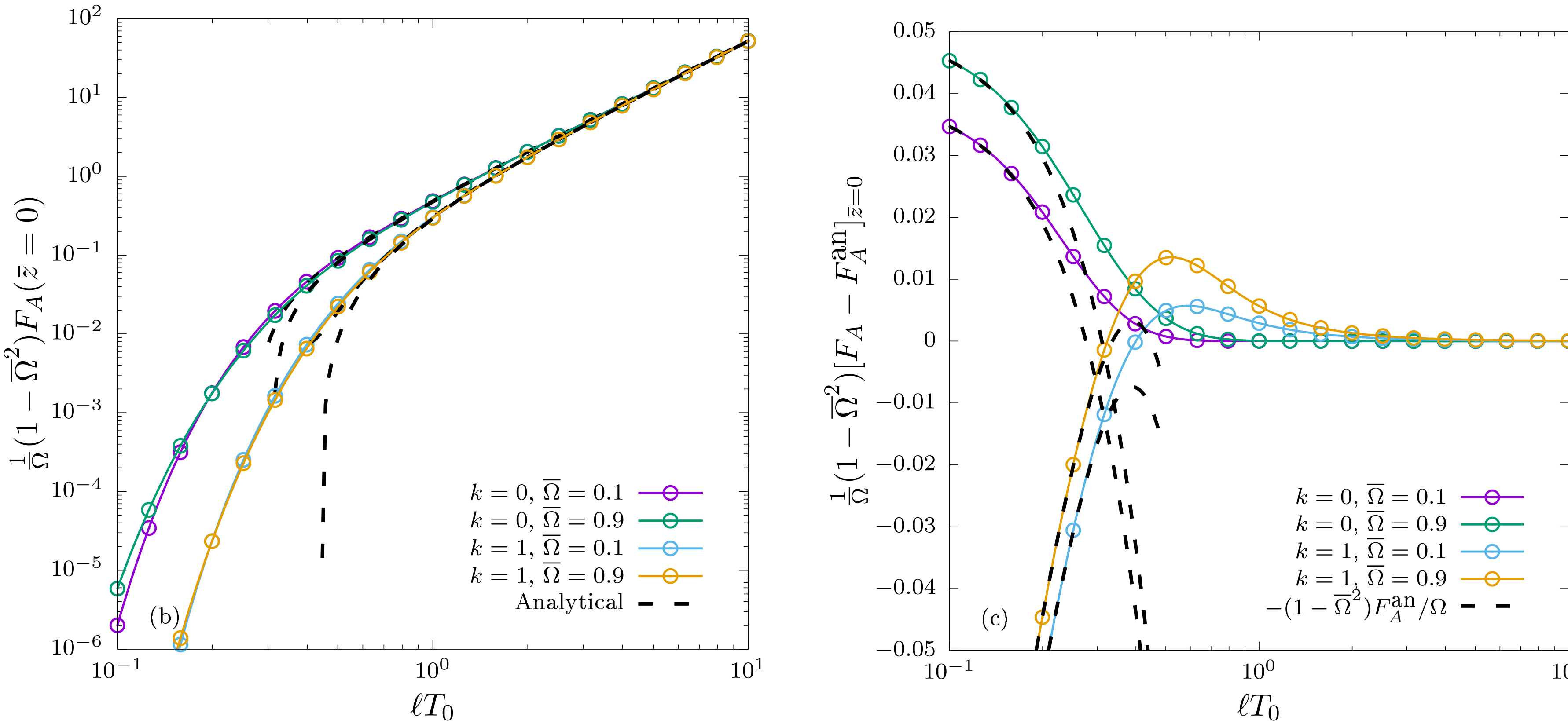
- At finite mass, $F_A(\overline{z})$ can be expanded for large $\overline{z}$ as




$$F_A(\overline{z}) = \frac{\Omega}{1-\overline{\Omega}^2}\frac{\Gamma(1+k)}{2\sqrt{\pi}\,\Gamma(\frac{1}{2}+k)} \times \left(\frac{\ell T_0}{\sqrt{1+\overline{z}^2}}\right)^{2k}\left\{\left[\zeta(2+2k)\left(1-\frac{1}{2^{1+2k}}\right)(2\ell T_0)^2 - \frac{\zeta(2k)}{6}\left(1-\frac{1}{2^{2k-1}}\right)(3+2k+\overline{\Omega}^2) + O(T_0^{-2})\right] + O(z^{-2})\right\} \quad (22)$$

- Since $F_A(\overline{z}) \simeq z^{-2k}$, a finite mass always ensures that $F(\pm\infty) = 0 \Rightarrow$ adS is opaque to non-chiral particles.

# High temperature asymptotics




$$F_A(\overline{z}) = \frac{\Omega}{1-\overline{\Omega}^2}\left\{\frac{\pi\ell^2T_0^2}{6} - \frac{3+\overline{\Omega}^2}{24\pi} - \frac{k^2(1+\overline{z}^2)}{2\pi}\ln\frac{\pi\ell T_0}{\sqrt{1+\overline{z}^2}} - \frac{k(1+\overline{z}^2)}{4\pi}[1+k-2k\mathcal{C}-2k\psi(1+k)] + O(T_0^{-1})\right\}. \qquad (23)$$

All terms are confirmed by comparison with numerical results.

# Conclusion

- New KMS relation allows $S^F_{\beta_0,\Omega}$ to be written in terms of $S^F_{\rm vac}$.
- Exact expression for $S^F_{\rm vac}$ allows vortical effects to be studied on adS.
- Axial vortical effect confirmed and curvature correction revealed.
- Axial flux $F_A$ of massless particles originates from southern hemisphere and escapes through northern hemisphere.
- For massive particles, $F_A \sim \overline{z}^{-2k}$ and no flux penetrates the boundary of adS.
- Possible extensions: supercritical rotation, finite chemical potential, thermodynamics...
- This work was supported through a grant of the Ministry of Research, Innovation and Digitization, CNCS - UEFISCDI, project number PN-III-P1-1.1-TE-2021-1707, within PNCDI III.

# Appendix

# We are hiring!

- FORQ: Facets of rotating quark-gluon plasma (July 2023 - June 2026)
- Project led by Maxim N. Chernodub in Timișoara
- 2 postdocs in July 2023 (adverts will be out soon...)
- One PhD student in October 2023
- 3 more postdocs in 2024
- 2 more PhD students in 2024
- 6 MSc students in 2024, 2025...