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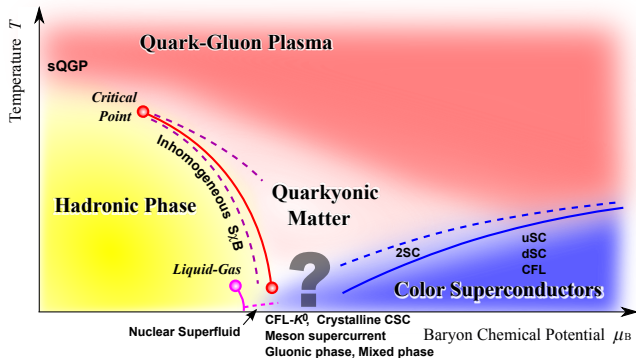
Rotation Effects on Gluon Mass and the Deconfinement Transition

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第二届“旋转与磁场下的 QCD 物质研讨会”

2026 年 9 月 14 日·淮南



K. Fukushima and T. Hatsuda, Rep. Prog. Phys. **74**, 014001 (2011).

Adding rotation

Heavy-ion collisions:

non-central collisions generate a highly vortical quark-gluon plasma, with a characteristic vorticity scale of $\Omega \sim 10^{21} \text{ s}^{-1}$.

STAR Collaboration, Nature **548**, 62 (2017).

Rapidly rotating neutron stars:

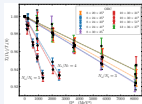
millisecond pulsars probe dense QCD matter under macroscopic rotation. For the fastest observed pulsars, $f \sim 10^3 \text{ Hz}$.

J. W. T. Hessels et al., Science **311**, 1901 (2006).

These systems motivate extending the conventional phase diagram as

$$(T, \mu_B) \longrightarrow (T, \mu_B, \Omega).$$

Lattice Yang-Mills

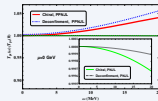


$T_c(\Omega_I) \downarrow$; analytic continuation gives

$$T_c(\Omega) \uparrow$$

V. V. Braguta *et al.*, Phys. Rev. D **103**, 094515 (2021).

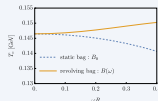
$T_c(\Omega) \uparrow$ – Polarized PNJL



Polarized Polyakov loop: $T_c(\Omega) \uparrow$.

F. Sun *et al.*, Phys. Rev. D **109**, 116017 (2024).

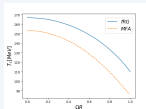
$T_c(\Omega) \uparrow$ – Revolving bag



Revolution effect: $T_c(\Omega) \uparrow$.

K. Mameda and K. Takizawa, Phys. Lett. B **847**, 138317 (2023).

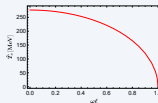
$T_\chi(\Omega) \downarrow$ – Quark-meson



Quark-meson/FRG: $T_\chi(\Omega) \downarrow$.

H.-L. Chen, Z.-B. Zhu and X.-G. Huang, Phys. Rev. D **108**, 054006 (2023).

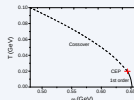
$T_c(\Omega) \downarrow$ – Holographic QCD



Holography: $T_c(\Omega) \downarrow$.

Y.-Q. Zhao *et al.*, JHEP **04**, 115 (2023).

$T_\chi(\Omega) \downarrow$ – NJL



NJL: rotation suppresses chiral condensation, $T_\chi \downarrow$.

Y. Jiang and J. Liao, Phys. Rev. Lett. **117**, 192302 (2016).

Model predictions are not universal: $T_c(\Omega) \uparrow$ or $T_c(\Omega) \downarrow$. **Motivation:** isolate the gluonic mechanism behind the lattice trend.

Lattice results

1. Deconfinement temperature

Rotating pure Yang–Mills lattice simulations find $T_c(\Omega)/T_c(0) = 1 + C_2\Omega^2 + \dots$ with $C_2 > 0$, namely rotation favors confinement and increases T_c .

V. V. Braguta *et al.*, Phys. Rev. D **103**, 094515 (2021).

2. Negative Barnett effect

Around $T \simeq (1 - 1.5)T_c$, the rotating gluon plasma can develop a negative moment of inertia, which has been interpreted as a negative gluonic Barnett effect associated with the spin response of the gluonic medium.

V. V. Braguta *et al.*, Phys. Rev. D **110**, 014511 (2024).

3. Spatially inhomogeneous phase

The local deconfinement temperature becomes radial dependent, $T_c = T_c(r)$, and increases away from the rotation axis. After analytic continuation to real rotation, this leads to an inhomogeneous equilibrium structure with a deconfined region near the axis and a confined outer region.

V. V. Braguta, M. N. Chernodub and A. A. Roenko, Phys. Lett. B **855**, 138783 (2024).

Effective-model studies

$T_c(\Omega) \uparrow$ / lattice-compatible

- **PNJL model**

F. Sun *et al.*, Phys. Rev. D **109**, 116017 (2024).

- **Holographic QCD**

Y. Chen *et al.*, Phys. Rev. D **111**, 046006 (2025). Z. Li, Y. Chen, and M. Huang, arXiv:2607.23199 [hep-ph] (2026).

- **Revolving bag model**

K. Mameda and K. Takizawa, Phys. Lett. B **847**, 138317 (2023).

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$T_c(\Omega) \downarrow$ / opposite trend

- **Quark–meson model**

H.-L. Chen, Z.-B. Zhu and X.-G. Huang, Phys. Rev. D **108**, 054006 (2023).

- **Holographic QCD**

X. Chen *et al.*, JHEP **07**, 132 (2021).
Y.-Q. Zhao *et al.*, JHEP **04**, 115 (2023).

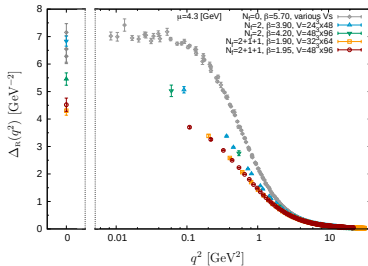
- **NJL model:**

Y. Jiang and J. Liao, Phys. Rev. Lett. **117**, 192302 (2016).
X. Wang *et al.*, Phys. Rev. D **99**, 016018 (2019).

...

- Landau-gauge lattice QCD shows that the gluon propagator is finite in the deep infrared.
- This is naturally described by a massive-type infrared propagator.
- The plotted data include the pure-gluon channel $N_f = 0$, which is the closest lattice input to our pure SU(3) Yang–Mills setup.

$$\Delta_R(q^2 \rightarrow 0) < \infty, \quad \Delta(q^2) \sim \frac{1}{q^2 + m_{\text{IR}}^2}.$$



A. Ayala *et al.*, Phys. Rev. D **86**, 074512 (2012).

Rotating Proca field
A massive spin-1 field obeys

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m^2 A_\mu A^\mu, \quad \partial_\mu A^\mu = 0,$$

leaving three physical polarizations.

In the corotating frame, $t' = t$ and $\phi' = \phi - \Omega t$. For a vector field,

$$i\partial_t \rightarrow i\partial_t - \Omega J_z, \quad J_z = L_z + S_z.$$

For a J_z eigenmode, $J_z A_\mu^{(j)} = j A_\mu^{(j)}$, the frequency shifts as $k_0 \rightarrow k_0 - j\Omega$.

Using $A_\mu \sim e^{-ik_0 t + ik_z z + ij\phi} \varepsilon_\mu^{(j,\lambda)}(\rho)$, the propagator is

$$D_{\mu\nu}(x, y) \sim \sum_{j,\lambda} \int dk_0 dk_z k_t dk_t \frac{e^{-ik_0 \Delta t} e^{ik_z \Delta z} e^{ij\Delta\phi} A_{\mu,\lambda}(\rho) A_{\nu,\lambda}^*(r)}{(k_0 - j\Omega)^2 - E_k^2 + i\eta},$$

with $E_k = \sqrt{k_t^2 + k_z^2 + m^2}$.

T. Xu and Y. Jiang, Chin. Phys. C **45**, 114102 (2021).

Proca field

$$\mathcal{L}_P = -\frac{1}{4}F_{\mu\nu}F_{\mu\nu} + \frac{1}{2}m^2 A_\mu A_\mu.$$

The equation of motion gives

$$\partial_\mu A_\mu = 0,$$

so

4 components $\longrightarrow$ 3 massive polarizations.

No ghost or auxiliary field is required.

Curci–Ferrari model

$$\mathcal{L}_{CF} = \frac{1}{4}F^2 + \frac{1}{2}m^2 A^2 + ih^a \partial_\mu A_\mu^a + \bar{c}^a \partial_\mu D_\mu^{ab} c^b.$$

The h field imposes $\partial_\mu A_\mu^a = 0$, while the ghost sector removes the unphysical gauge contribution. At one loop,

$$\frac{3}{2} \ln(Q^2 + m^2) - \frac{1}{2} \ln Q^2 \xrightarrow{Q^2 \gg m^2} \boxed{\ln Q^2},$$

recovering the two Yang–Mills polarizations.

CF gluons are not Proca particles: the ghost and auxiliary sectors preserve the Yang–Mills degree-of-freedom counting.

At quadratic order,

$$\mathcal{K}_{\mu\nu}(Q) = (Q^2 + m^2)\delta_{\mu\nu} - Q_\mu Q_\nu = (Q^2 + m^2)P_{\mu\nu}^\perp + m^2 P_{\mu\nu}^\parallel.$$

Since

$$\text{tr } P^\perp = 3, Q^2 = \omega_n^2 + q_t^2 + q_z^2 + m^2$$

the three transverse directions give

$$\Gamma_T^{(1)} = \frac{3}{2} \text{Tr} \ln(Q^2 + m^2).$$

For the longitudinal gluon and h field,

$$\det \begin{pmatrix} m^2 & Q \\ Q & 0 \end{pmatrix} \propto Q^2 \quad \Rightarrow \quad \Gamma_{Lh}^{(1)} = +\frac{1}{2} \text{Tr} \ln Q^2,$$

while the Grassmann ghost contributes

$$\Gamma_{\text{gh}}^{(1)} = -\text{Tr} \ln Q^2.$$

Therefore,

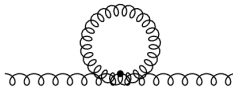
$$\boxed{\Gamma_{\text{CF}}^{(1)} = \frac{3}{2} \text{Tr} \ln(Q^2 + m^2) - \frac{1}{2} \text{Tr} \ln Q^2.}$$

For $Q^2 \gg m^2$,

$$\frac{3}{2} \ln(Q^2 + m^2) - \frac{1}{2} \ln Q^2 \longrightarrow \ln Q^2 = 2 \times \frac{1}{2} \ln Q^2.$$

Thus the CF model contains a gluon mass scale, but recovers the two Yang–Mills polarizations rather than three Proca degrees of freedom.

Local analysis at the rotation axis:: $\rho = r = 0$



From the Yang–Mills interaction

$$\mathcal{L}_4 = \frac{g^2}{4} f^{abe} f^{cde} A_\mu^a A_\nu^b A_\mu^c A_\nu^d \Rightarrow \Pi_{\mu\nu}^{4g,ab} = g^2 C_A \delta^{ab} (\delta_{\mu\nu} D_{\alpha\alpha} - D_{\mu\nu}).$$

Rotation enters only through the internal propagator

$$D_{\mu\nu} = P_{\mu\nu}^{[+1]} G_+ + P_{\mu\nu}^{[0]} G_0 + P_{\mu\nu}^{[-1]} G_-, \quad G_\pm = \frac{1}{(\omega_n \mp i\Omega)^2 + E_q^2}, \quad G_0 = \frac{1}{\omega_n^2 + E_q^2}.$$

Pure rotational part

$$\Delta_\Omega \Pi_{\mu\nu}^{4g} = \Pi_{\mu\nu}^{4g}(\Omega) - \Pi_{\mu\nu}^{4g}(0), \quad \Delta S = G_+ + G_- - 2G_0, \quad \Delta D = G_+ - G_-, \quad A(Q) = 1 - \frac{q_t^2}{2Q^2}.$$

$$\Delta_\Omega \Pi_{\mu\nu}^{4g} = g^2 C_A T \sum_n \int_q A(Q) \begin{pmatrix} \Delta S & 0 & 0 & 0 \\ 0 & \Delta S/2 & i\Delta D/2 & 0 \\ 0 & -i\Delta D/2 & \Delta S/2 & 0 \\ 0 & 0 & 0 & \Delta S \end{pmatrix}_{\mu\nu}$$

$$m^2 \rightarrow m_{\text{IR}}^2 + P_{\mu\nu}^{\pm,0}(K \rightarrow 0) \Delta_\Omega \Pi_{\mu\nu}^{4g}$$

Temporal color background

$$\bar{A}_0 = \bar{A}_0^3 t^3 + \bar{A}_0^8 t^8,$$

$$q_{0\kappa} = \omega_n + g(\kappa_3 \bar{A}_0^3 + \kappa_8 \bar{A}_0^8).$$

Rotation-induced mode splitting

$$\boxed{\omega_{n\kappa j} = q_{0\kappa} - j\Omega_I}, \quad j = +1, 0, -1.$$

Polarization weights

$$A_\kappa = 1 - \frac{q_t^2}{2Q_\kappa^2}, \quad B_\kappa = 1 + \frac{q_t^2}{Q_\kappa^2},$$

One-loop CF potential

$$\boxed{V_{\text{eff}}(\Omega_I, \bar{A}_0) = \frac{T}{2} \sum_{\kappa, n} \int_q \left[A_\kappa \ln \left((q_{0\kappa} - \Omega_I)^2 + q^2 + m^2 \right) \right. \\ \left. + B_\kappa \ln \left(q_{0\kappa}^2 + q^2 + m^2 \right) + A_\kappa \ln \left((q_{0\kappa} + \Omega_I)^2 + q^2 + m^2 \right) - \ln Q_\kappa^2 \right].}$$

Polyakov loop definition

For $SU(3)$,

$$\langle L \rangle = \frac{1}{3} \left[2e^{i \frac{\beta g \bar{A}_0^8}{2\sqrt{3}}} \cos\left(\frac{\beta g \bar{A}_0^3}{2}\right) + e^{-i \frac{\beta g \bar{A}_0^8}{\sqrt{3}}} \right]$$

Stationary background

For each T and Ω_I , solve

$$\frac{\partial V_{\text{eff}}}{\partial \bar{A}_0^3} = 0, \quad \frac{\partial V_{\text{eff}}}{\partial \bar{A}_0^8} = 0.$$

Polyakov loop as an order parameter

- **Confinement phase:**

$$\langle L \rangle \simeq 0$$

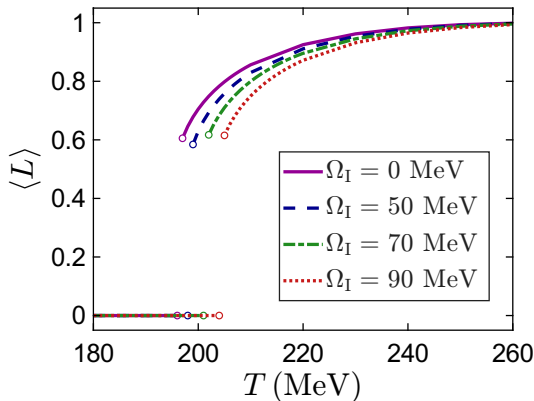
corresponding to an approximately center-symmetric background.

- **Deconfinement phase:**

$$\langle L \rangle > 0$$

and at sufficiently high temperature,

$$\langle L \rangle \rightarrow 1.$$



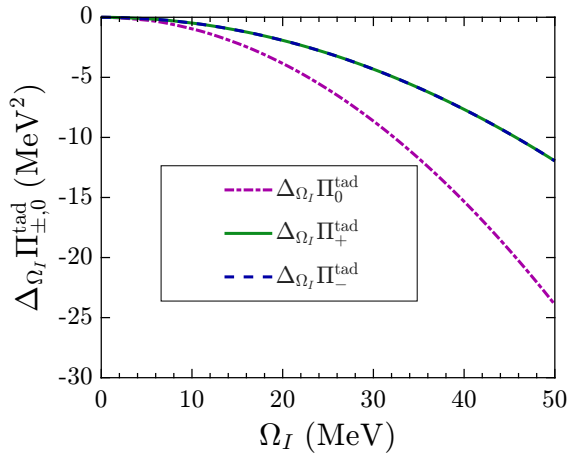
Polyakov loop vs temperature under imaginary rotation at
 $g = 0$

Key observations

- A discontinuous jump in $\langle L \rangle$ indicates a **first-order deconfinement transition**.
- Increasing Ω_I shifts the jump to higher temperature and raises T_c .
- Imaginary rotation therefore **stabilizes the confined phase** in the calculation at fixed gluon mass.

$$\Omega_I \uparrow \implies T_c(\Omega_I) \uparrow$$

$$m_{\text{IR}} = 0.54 \text{ GeV}$$



Four-gluon Tadpole under Imaginary Rotation ($T = 250$ MeV)

Rotation-induced tadpole correction

$$\varepsilon_{\mu}^j \Pi_{\mu\nu} \varepsilon_{\nu}^j.$$

- The four-gluon tadpole gives a **positive** contribution to the gluon screening mass.
- The correction grows approximately as

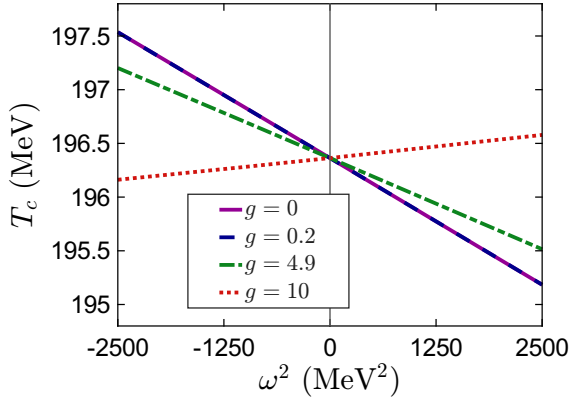
$$\Delta\Pi_s^{\text{tad}} \propto \Omega^2.$$

- The two circular transverse modes remain degenerate:

$$\Delta\Pi_+^{\text{tad}} \simeq \Delta\Pi_-^{\text{tad}}.$$

- The third transverse channel is more strongly screened:

$$\Delta\Pi_0^{\text{tad}} < \Delta\Pi_{\pm}^{\text{tad}} < 0$$



Coupling controls the rotational response

- For weak and moderate coupling,
 $g = 0, 0.2, 4.9$,
 the transition temperature decreases with
 increasing ω^2 :

$$\frac{dT_c}{d\omega^2} < 0.$$

- The slope becomes weaker as the interaction strength increases.
- At sufficiently strong coupling, e.g. $g = 10$, the trend is reversed:

$$\frac{dT_c}{d\omega^2} > 0$$

The effect of rotation on deconfinement is therefore **not purely kinematic**. It is strongly affected by the interaction-induced gluon screening.

rotation + gluon interaction $\implies$ different T_c responses

Summary

- We formulated rotating $SU(3)$ Yang–Mills thermodynamics in the massive Curci–Ferrari framework at $r = 0$, with gluon modes organized by $j = +, -, 0$.
- At fixed gluon mass, imaginary rotation shifts the **first-order deconfinement transition** to higher temperature.
- The four-gluon tadpole preserves the degeneracy of the $+$ and $-$ channels. In this calculation, the $s = 0$ correction has twice their magnitude.
- The response of T_c to ω^2 depends strongly on g , showing the importance of interaction-induced screening.

Outlook

- Add the three-gluon loop/ghost loop and complete the renormalized, momentum-dependent gluon self-energy.
- Extend the local analysis to finite radius and impose the causality bound $|\Omega|R < 1$.

Next test: Does the increase of T_c survive after all one-loop contributions are included self-consistently?

Thank you for your attention!