First Principles Lattice QCD Calculations of nEDMs

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T2 Seminar

Electric Dipole Moments (EDMs) are very sensitive probes of CP violation ($\not{CP}$)

- **Baryon Asymmetry of the Universe**
  - Sakharov's conditions[1] (post-inflation):
    1. baryon-number violation
    2. evolution has to occur out of equilibrium
    3. C and CP violation
- **Source of CP Violation**
  - The strength of the $\not{CP}$ in the CKM matrix is much too small to explain baryogenesis.
  - Also no CP violation in the lepton sector has been observed so far.
  - In the Standard Model:
    - ★ The SM has an additional source of CP violation arising from the effect of QCD instantons.

$$\mathcal{L}_{\rm QCD} \longrightarrow \mathcal{L}^{\not{CP}}_{\rm QCD} = \mathcal{L}_{\rm QCD} + i\Theta \frac{G^a_{\mu\nu}\widetilde{G}^a_{\mu\nu}}{32\pi^2} \tag{1}$$

    - ★ Because of asymptotic freedom, all nonperturbative configurations including instantons are strongly suppressed at high temperatures where baryon number violating processes occur.

[1] Pisma Zh.Eksp.Teor.Fiz. 5 (1967) 32.

nEDM as a Probe of $\not{CP}$

- **Additional much larger $\not{CP}$ is needed from physics beyond the SM (BSM)**
  - Using tools of EFT, one can organize all possible effective $\not{CP}$ interactions of quarks and gluons based on **symmetry** and **dimension**, and independent of specific BSM theory

$$
\begin{aligned}
\mathcal{L}_{d\leq 6}^{\not{CP}} = & \mathcal{L}_{QCD} \\
& + i\Theta \frac{G^a_{\mu\nu}\widetilde{G}^a_{\mu\nu}}{32\pi^2} && \text{dim=4 QCD } \Theta\text{-term} \\
& + i\sum_q d_q^\gamma \bar{q}\sigma^{\mu\nu}\tilde{F}_{\mu\nu}q && \text{dim=5 quark EDMs} \\
& + i\sum_q d_q^G \bar{q}\sigma^{\mu\nu}\tilde{G}_{\mu\nu}q && \text{dim=5 quark chromo-EDMs} \\
& + d_G f^{abc}G^a_{\mu\nu}\tilde{G}^{\nu\beta,b}G^{\mu,c}_{\beta} && \text{dim=6 Weinberg's 3g operator} \\
& + \cdots
\end{aligned}
$$

Connecting to nEDMs

For Θ-term:

- Under a chiral transformation, one can rotate Θ into a complex phase of the quark matrix and vice versa.
- $\overline{\Theta} = \Theta + \text{Arg Det} M_q$
- If the overall $\overline{\Theta}$ is nonzero, then this operator would induce an nEDM d_n of size

$$d_n = \overline{\Theta}\; X = \overline{\Theta}\; \lim_{q^2 \to 0} \frac{F_3(q^2)}{2M_N \overline{\Theta}} \,. \tag{2}$$

Here X is obtained from the $\not{CP}$ part of the matrix element of the electromagnetic vector current within the neutron state in the presence of the Θ-term and F_3 is the $\not{CP}$ violating form factor.

Expansion about the small Θ, $d_q^{\gamma,g}$ and d_G

- The approach that works for lattice QCD is to treat the small Θ, $d_q^{\gamma,g}$ and d_G as perturbations and expand the theory about the normal CP conserving action, such as the Wilson-clover action

$$\langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{\Theta} \approx \langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{\Theta=0} - i\Theta \left\langle N \left| J_\mu^{\text{EM}} \int d^4x \, \frac{G_{\mu\nu}^a \widetilde{G}_{\mu\nu}^a}{32\pi^2} \right| N \right\rangle , \tag{3}$$

$$\langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{cEDM} \approx \langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{d_q^G=0} - i d_q^G \left\langle N \left| J_\mu^{\text{EM}} \int d^4x \, \sum_q \bar{q}\sigma^{\mu\nu}\tilde{G}_{\mu\nu} q \right| N \right\rangle , \tag{4}$$

$$\langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{G} \approx \langle N \mid J_\mu^{\text{EM}} \mid N\rangle\Big|^{d_G=0} - d_G \left\langle N \left| J_\mu^{\text{EM}} \int d^4x \, f^{abc} G_{\mu\nu}^a \tilde{G}^{\nu\beta,b} G_\beta^{\mu,c} \right| N \right\rangle . \tag{5}$$

Path Integrals in QFT

Consider gluon field $\phi(x)$ and an action $S[A_\mu(x)] = \int d^4x \mathcal{L}[A_\mu(x)]$

$$\mathcal{Z} = \int \mathcal{D}A_\mu(x)\ e^{-iS[A_\mu(x)]} \tag{6}$$

then

- Continuum $\longrightarrow$

$$\int \mathcal{D}A_\mu(x) = \cdots \int dA_{\mu 1} \cdots \int dA_{\mu 0.000001} \cdots \int dA_{\mu 1} \cdots \int dA_{\mu \infty} \cdots \tag{7}$$

- Finite Volume$\longrightarrow$

$$\int \mathcal{D}A_\mu(x) = \int dA_{\mu 0} \int dA_{\mu \frac{1L}{2\pi}} \int dA_{\mu \frac{2L}{2\pi}} \int dA_{\mu \frac{3L}{2\pi}} \int dA_{\mu \frac{4L}{2\pi}} \cdots \tag{8}$$

- Discretized $\longrightarrow$

$$\int \mathcal{D}A_\mu(x) = \prod_x \int dA_{\mu x} = \int dA_{\mu 0} \int dA_{\mu 1} \int dA_{\mu 2} \int dA_{\mu 3} \int dA_{\mu 4} \cdots \tag{9}$$

Euclidean Path Integral in QFT

Even in the discretized lattice, we have pratical problem in

$$\mathcal{Z} = \int \mathcal{D}\psi(x)\ e^{-iS[A_\mu(x)]} \tag{10}$$

make a Wick rotation: $\boxed{t \longrightarrow -it}$ then

$$\boxed{-iS = -i\int d^3x dt \mathcal{L}} \longrightarrow \boxed{-\int d^3x dt \mathcal{L}_E = -S_E} \tag{11}$$

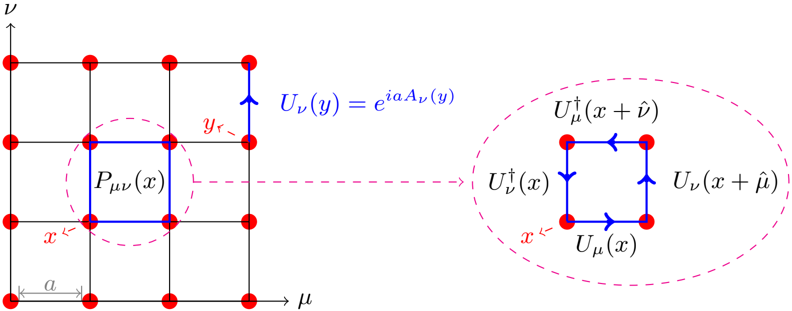
Euclidean path integral

$$\boxed{\mathcal{Z}_E = \int \mathcal{D}A_\mu(x)\ e^{-S_E[A_\mu(x)]}} \tag{12}$$

Lattice QCD[2]

$\Lambda_4 = \{n_\mu = (n_1, n_2, n_3, n_4) | n_i \in a[0, 1, \cdots, L_i - 1]\}$

$$\boxed{S_{gauge} = \frac{2}{g^2} \sum_{x \in \Lambda_4} \sum_{\mu<\nu} (1 - Re\ Tr\ [P_{\mu\nu}(x)])} \tag{13}$$



where the elementary plaquette, $\boxed{P_{\mu\nu}(x) = U_\mu(x) U_\nu(x+\hat{\mu}) U_\mu^\dagger(x+\hat{\nu}) U_\nu^\dagger(x)}$.

[2] K. G. Wilson, Confinement of Quarks, Phys. Rev. D10 (1974) 2445.

Lattice QCD

$$\boxed{S_{gauge} = \frac{2}{g^2}\sum_{x\in\Lambda_4}\sum_{\mu<\nu}(1 - Re\ Tr\ [P_{\mu\nu}(x)]) = \frac{a^4}{2g^2}\sum_{x\in\Lambda_4}\sum_{\mu<\nu} tr\ [F_{\mu\nu}(x)^2] + \mathcal{O}(a^2)}.$$

Physical observables $\mathcal{O}$ are evaluated as an expectation value over the relevant degrees of freedom

$$\boxed{\langle\mathcal{O}\rangle = \frac{\int \mathcal{D}[U]\mathcal{D}[\psi]\mathcal{D}[\bar{\psi}]\ \mathcal{O}\ e^{-[S_{gauge}+\int dx\bar{\psi}\not{D}\psi]}}{\int \mathcal{D}[U]\mathcal{D}[\psi]\mathcal{D}[\bar{\psi}]e^{-[S_{gauge}+\int dx\bar{\psi}\not{D}\psi]}}}. \tag{14}$$

The quark fields ψ & $\bar{\psi}$ are Grassmann variables:

$$\boxed{\langle\mathcal{O}\rangle = \frac{\int \mathcal{D}[U]\ \widetilde{\mathcal{O}}\ det\not{D}^2 e^{-[S_{gauge}]}}{\int \mathcal{D}[U] det\not{D}^2 e^{-[S_{gauge}]}}}. \tag{15}$$

This integration results in the "contraction" of fermion–anti-fermion pairs in all possible ways (Wick's theorem), replacing them with quark propagators $\not{D}^{-1}$.

UV Divergences & Renormalization

- **Lattice as a UV cutoff:**

$$O_N \quad (\text{defined on } N \text{ sites, } a \sim N^{-1/4})$$

acts as a hard cutoff at momentum scale $p_{\max} \sim \pi/a$. As $N \to \infty$ (so $a \to 0$), loop integrals pick up large-momentum modes:

$$\langle O_N \rangle \;\sim\; c_1\, a^{-\Delta} \;+\; (\text{log terms}) \;+\; \text{finite}, \quad \Delta \geq 0.$$

In particular, $\lim_{a\to 0} \langle O_N \rangle = \infty$ unless counterterms are introduced.

- **Multiplicative renormalization:**

$$O_R \;=\; \lim_{a\to 0} Z_O(a)\, O_N, \quad Z_O(a) \sim a^{\Delta}\Big[1 + O(g_0^2 \ln a)\Big],$$

ensures $\langle O_R \rangle$ stays finite in the continuum limit.

Operator Mixing

- **Operator mixing on the lattice:**

$$O_N \;=\; \sum_i Z_{O,i}(a)\, O_i^{(\mathrm{cont})}(\mu), \qquad \dim O_i \le \dim O.$$

  - Each $O_i^{(\mathrm{cont})}$ is a continuum operator of equal or lower dimension.
  - Determining all coefficients $Z_{O,i}(a)$ is nontrivial—mixing may involve identity operators, mass terms, or higher-derivative operators.
  - In general, $\langle O_N \rangle$ picks up contributions from every lower dimension $O_i^{(\mathrm{cont})}$, each with its own divergent coefficient in a.

Gradient Flow as a Renormalization-Smearing

The gradient flow method introduces an extra fictitious coordinate t [flow-time] (with dimensions $(\text{length})^2$) and defines the gauge filed $B_\mu(x,t)$ at positive flow-time t by the following ordinary differential equations[3]:

$$\frac{d}{dt}B_\mu(x,t) = D_\nu G_{\nu\mu}(x,\tau_{gf}) = \frac{i}{gT^a}D_\nu\left[D_\nu, D_\mu\right]; \tag{16}$$

$$B_\mu(x,t)\Big|_{(t=0)} = A_\mu(x); \tag{17}$$

$$D_\mu = \partial_\mu + \left[B_\mu(x,t), \frac{d}{dt}\right]; \tag{18}$$

$$G_{\mu\nu}(x,t) = \partial_\mu B_\nu(x,t) - \partial_\nu B_\mu(x,t) + \left[B_\mu(x,t), B_\nu(x,t)\right]. \tag{19}$$

[3] Lüscher, M. Properties and uses of the Wilson flow in lattice QCD. J. High Energ. Phys. 2010, 71 (2010). https://doi.org/10.1007/JHEP08(2010)071

Lattice gradient flow

In the lattice QCD, the simplest choice of the gauge action is the Wilson action,

$$S_W(U) = \frac{2}{g_0^2} \sum_{x \in \Lambda_4} \sum_{\mu<\nu} \left(1 - Re\ Tr\ \left[U_\mu(x) U_\nu(x+\hat{\mu}) U_\mu^\dagger(x+\hat{\nu}) U_\nu^\dagger(x)\right]\right) \tag{20}$$

then the associated flow $V_t(x,\mu)$ of lattice gauge fields (the "Wilson flow" is defined by the equations)

$$\frac{d}{dt} V_t(x,\mu) = -g_0^2 \{\partial_{x,\mu} S_W(V_t)\} V_t(x,\mu), \qquad V_t(x,\mu)\Big|_{t=0} = U(x,\mu) \tag{21}$$

The $\partial_{x,\mu}$ ($\mathfrak{su}(3)$-valued differential operator) acting on a differential function $f(U)$ of the gauge filed reads

$$\partial^a_{x,\mu} f(U) = \frac{d}{ds} f\left(e^{sX} U\right)\Big|_{s=0}, \qquad X(y,\nu) = \begin{cases} T^a & \text{if } (y,\nu) = (x,\mu), \\ 0 & \text{otherwise.} \end{cases} \tag{22}$$

While these depend on the choice of the generators T^a, the combination

$$\partial_{x,\mu} f(U) = T^a \partial^a_{x,\mu} f(U) \tag{23}$$

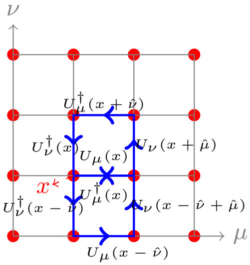
can be shown to be basis-independent.

Lattice gradient flow

The lattice gradient flow equation in the Wilson gauge action is given as

$$\frac{dV_t(x,\mu)}{dt} = -i\left[\frac{1}{2}\left(Y_\mu(x,t) - Y_\mu^\dagger(x,t)\right) - \frac{1}{6}\;Tr\left(Y_\mu(x,t) - Y_\mu^\dagger(x,t)\right)\right]V_t(x,\mu) \qquad (24)$$

where, $Y_\mu(x,t) = U_\mu(x)R_{\mu\nu}(x)$.



Automatic UV Finiteness at $t > 0$

- At positive flow time $t > 0$, the gauge field $B_\mu(t,x)$ (or lattice $V_t(x,\mu)$) is smeared over a radius $\sqrt{8\,t}$. Concretely, in the continuum:

$$B_\mu(t,x) \;=\; \int d^4y\; K_t(x-y)\; A_\mu(y) \;+\; \mathcal{O}(g_0^2),$$

  where $K_t(x-y)$ is a heat-kernel of width $\sqrt{8\,t}$. On the lattice, a similar smearing holds to leading order in t/a^2.

- Because UV fluctuations (momenta $p \gtrsim 1/\sqrt{t}$) are exponentially suppressed, any composite operator constructed from V_t at fixed physical $t > 0$ does not require additional multiplicative or additive renormalization as $a \to 0$. Equivalently,

$$\lim_{a\to 0} \big\langle\, O_t[\,V_t\,] \big\rangle \;=\; \text{finite},$$

Implementing Gradient Flow: Gauge Fields

For Wilson flow of gauge field, we have

$$Z(U,\mu) = \sum_{\mu \neq \nu} \left(U_\mu(x) * U_\nu(x+\hat{\mu}) * U_\mu^\dagger(x+\hat{\nu}) * U_\nu^\dagger + U_\mu(x) * U_\nu^\dagger(x+\hat{\mu}-\hat{\nu}) * U_\mu^\dagger(x-\hat{\nu}) * U_\nu(x-\hat{\nu}) \right) \quad (25)$$

then

$$Z(U,\mu) = \frac{1}{2}\left(Z(U,\mu) - Z^\dagger(U,\mu) \right) - \frac{1}{6}\mathrm{Tr}\left[Z(U,\mu) - Z^\dagger(U,\mu) \right] \quad (26)$$

Assuming V_t is known at some flow time t, an approximation to the exact solution $V_{t+\epsilon}$ at time $t+\epsilon$ is obtained by computing the fields

$$W_0 = V_t, \quad (27)$$

$$W_1 = \exp\{\frac{1}{4} Z_0\} W_0, \quad (28)$$

$$W_2 = \exp\{\frac{8}{9} Z_1 - \frac{17}{36} Z_0\} W_1, \quad (29)$$

$$W_3 = \exp\{\frac{3}{4} Z_2 - \frac{8}{9} Z_1 + \frac{17}{36} Z_0\} W_2, \quad (30)$$

where

$$Z_i = \epsilon Z(W_i), \qquad i = 0, 1, 2. \quad (31)$$

the ϵ flowed gauge field is now W_3.

Implementing Gradient Flow: Fermions

For propagators $\texttt{prop}$, given V_t and $\texttt{prop}_t$ at some flow time t, and assuming the gauge fields W_0, W_1 and W_2 are as above, we have

$$\phi_1 = \texttt{prop} + \frac{1}{4}\Delta(W_0, \texttt{prop}), \tag{32}$$

$$\phi_2 = \phi_0 + \frac{8}{9}\Delta(W_1, \phi_1) - \frac{2}{9}\Delta(W_0, \texttt{prop}), \tag{33}$$

$$\phi_3 = \phi_1 + \frac{3}{4}\Delta(W_2, \phi_2), \tag{34}$$

where

$$\Delta(U, \texttt{prop}) = D_\mu D_\mu(U, \texttt{prop}) = \epsilon(U_\mu(x)\texttt{prop}(x+\mu) - U_\mu^\dagger(x-\mu)\texttt{prop}(x-mu))/2 \tag{35}$$

the ϵ flowed propagator is now ϕ_3.

How do we do Lattice calculations?

The proton two-point correlation function:

$$C_{\alpha\beta}(t,\mathbf{p}) = a^3 \sum_{\mathbf{x}} e^{-i\mathbf{p}.\mathbf{x}} C_{\alpha\beta}(t,\mathbf{x}) = a^3 \sum_{\mathbf{x}} e^{-i\mathbf{p}.\mathbf{x}} \langle 0| \chi_\alpha(\mathbf{x},t)\widetilde{\chi}_\beta(\mathbf{0},0)|0\rangle \tag{36}$$

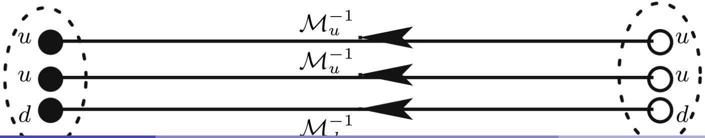
where

$$\chi_\alpha(\mathbf{x},t) = e^{i\alpha\gamma^5} \epsilon^{ijk} u^i_\alpha(\mathbf{x},t) u^j_\gamma(\mathbf{x},t)[C^{-1}\gamma_5]_{\gamma\delta} d^k_\delta(\mathbf{x},t) \tag{37}$$

- need to determine α to make it parity even

After integrating out the quark fields in the path integral formulation, this correlator is expressed in terms of products of the inverse of the Dirac operator

$$C_{\alpha\beta}(t,\mathbf{p}) = -\, a^3 \sum_{\mathbf{x}} e^{-i\mathbf{p}.\mathbf{x}} \epsilon^{ijk}\epsilon^{i'j'k'} [C^{-1}\gamma_5]_{\alpha'\alpha''}[\gamma_5 C]_{\beta'\beta''}$$
$$\times \left\langle [\mathcal{M}_d^{-1}]^{ki'}_{\alpha''\beta'} \left\{ [\mathcal{M}_u^{-1}]^{jj'}_{\alpha'\beta''} [\mathcal{M}_u^{-1}]^{ik'}_{\alpha\beta} - [\mathcal{M}_u^{-1}]^{ij'}_{\alpha\beta''} [\mathcal{M}_d^{-1}]^{ik'}_{\alpha'\beta} \right\} \right\rangle$$



Proton Two Point Function

On inserting a complete set of states between the interpolating operators[4]

$$C_{\alpha\beta}(t,\mathbf{p}) = a^3 \sum_{\mathbf{x}} e^{-i\mathbf{p}.\mathbf{x}} \langle 0| \chi_\alpha(\mathbf{x},t)\widetilde{\chi}_\beta(\mathbf{0},0) |0\rangle \tag{38}$$

$$= a^3 \sum_{n,\sigma} \frac{e^{-E_n(\mathbf{p})t}}{2E_n(\mathbf{p})} \langle 0| \chi_\alpha |n;\mathbf{p},\sigma\rangle \langle n;\mathbf{p},\sigma| \widetilde{\chi}_\beta |0\rangle \tag{39}$$

$$= a^3 Z(\mathbf{p}) \sum_{\sigma} u_\alpha(n=0,\mathbf{p},\sigma)\widetilde{u}_\beta(n=0,\mathbf{p},\sigma)\frac{e^{-E_n(\mathbf{p})t}}{2E_n(\mathbf{p})} + \cdots \tag{40}$$

Tracing this correlator against a given Dirac structure, often chosen to be $\Gamma^+ = \frac{1}{2}(1+\gamma_4)$, leads to

$$\boxed{C_{\Gamma^+}(t,\mathbf{p})} = \Gamma^+_{\beta\alpha} C_{\alpha\beta}(t,\mathbf{p}) \xrightarrow{t\to\infty} \boxed{Ce^{-E_n(\mathbf{p})t}} \tag{41}$$

where C is a time-independent constant.

[4]Refer: Boussarie, Renaud, et al. "TMD Handbook." arXiv:2304.03302 [hep-ph].

Proton Two Point Function: Gradient Flow

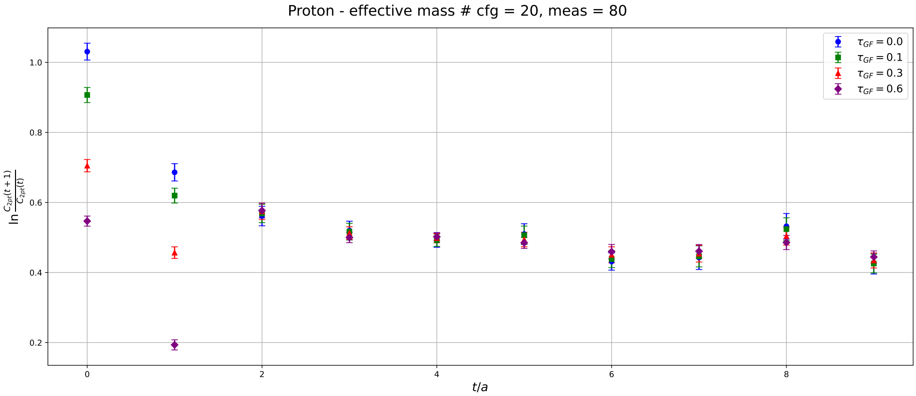


Figure 1: Proton 2-pt for different flow time ($\tau_{\mathrm{gf}} \equiv \sqrt{8t}$ for the flow time, where t is the parameter in the flow equations)

Behavior of the Correlation Functions ($\tau_{gf} = 0$)

Using spectral decomposition, the behavior of two- and three-point functions is given by the expansion:

$$
\begin{aligned}
C^{2\text{pt}}(t_f, t_i) = \\
& |\mathcal{A}_0|^2 e^{-aM_0(t_f - t_i)} + |\mathcal{A}_1|^2 e^{-aM_1(t_f - t_i)} + \\
& |\mathcal{A}_2|^2 e^{-aM_2(t_f - t_i)} + |\mathcal{A}_3|^2 e^{-aM_3(t_f - t_i)} + \ldots ,
\end{aligned}
\tag{42}
$$

$$
\begin{aligned}
C_\Gamma^{3\text{pt}}(t_f, \tau, t_i) = \\
& |\mathcal{A}_0|^2 \langle 0|\mathcal{O}_\Gamma|0\rangle e^{-aM_0(t_f - t_i)} + \\
& |\mathcal{A}_1|^2 \langle 1|\mathcal{O}_\Gamma|1\rangle e^{-aM_1(t_f - t_i)} + \\
& |\mathcal{A}_2|^2 \langle 2|\mathcal{O}_\Gamma|2\rangle e^{-aM_2(t_f - t_i)} + \\
& \mathcal{A}_1 \mathcal{A}_0^* \langle 1|\mathcal{O}_\Gamma|0\rangle e^{-aM_1(t_f - \tau)} e^{-aM_0(\tau - t_i)} + \\
& \mathcal{A}_0 \mathcal{A}_1^* \langle 0|\mathcal{O}_\Gamma|1\rangle e^{-aM_0(t_f - \tau)} e^{-aM_1(\tau - t_i)} + \\
& \mathcal{A}_2 \mathcal{A}_0^* \langle 2|\mathcal{O}_\Gamma|0\rangle e^{-aM_2(t_f - \tau)} e^{-aM_0(\tau - t_i)} + \\
& \mathcal{A}_0 \mathcal{A}_2^* \langle 0|\mathcal{O}_\Gamma|2\rangle e^{-aM_0(t_f - \tau)} e^{-aM_2(\tau - t_i)} + \\
& \mathcal{A}_1 \mathcal{A}_2^* \langle 1|\mathcal{O}_\Gamma|2\rangle e^{-aM_1(t_f - \tau)} e^{-aM_2(\tau - t_i)} + \\
& \mathcal{A}_2 \mathcal{A}_1^* \langle 2|\mathcal{O}_\Gamma|1\rangle e^{-aM_2(t_f - \tau)} e^{-aM_1(\tau - t_i)} + \ldots ,
\end{aligned}
\tag{43}
$$

Fit With Different t_{sep} ($\tau_{GF} = 0$)

we analyze the correlation functions with the following combinations $\{N_{2pt}, N_{3pt}\} = \{2,2\}$, $\{3,2\}$, $\{4,2\}$, $\{3,3\}$ and $\{4,3\}$ where the first (second) value is the number of states included in fits to the two-point (three-point) functions.
Axial Charge:

$\{N_{2pt}, N_{3pt}\}$	aM_0	aM_1	A_0^2	A_1^2	A_1^2/A_0^2	χ^2/d.o.f	AIC
$\{2,2\}$	0.4727(0083)	0.8509(0078)	5.40(41)e-10	3.61(80)e-10	1.4980(4383)	1.1534(4509)	55.522
$\{3,2\}$	0.4757(0094)	0.8481(0122)	5.63(52)e-10	2.59(151)e-10	2.1801(14704)	1.1575(5267)	57.354
$\{3,3\}$	0.4761(0084)	0.8479(0130)	5.66(46)e-10	2.53(139)e-10	2.2365(13971)	1.2371(5559)	62.349
$\{4,2\}$	0.4757(0090)	0.8481(0127)	5.63(51)e-10	2.59(148)e-10	2.1796(14387)	1.2289(5932)	61.324
$\{4,3\}$	0.4761(0081)	0.8480(0130)	5.66(44)e-10	2.53(132)e-10	2.2401(13408)	1.3183(5323)	66.230

Table 1: Estimates of the nucleon masses M_0 and M_1 and the amplitudes A_0 and A_1 extracted from simultaneous fit: smearing parameter $\sigma = 5$

Effective Mass: Fit With Different t_{sep} ($\tau_{GF} = 0$)

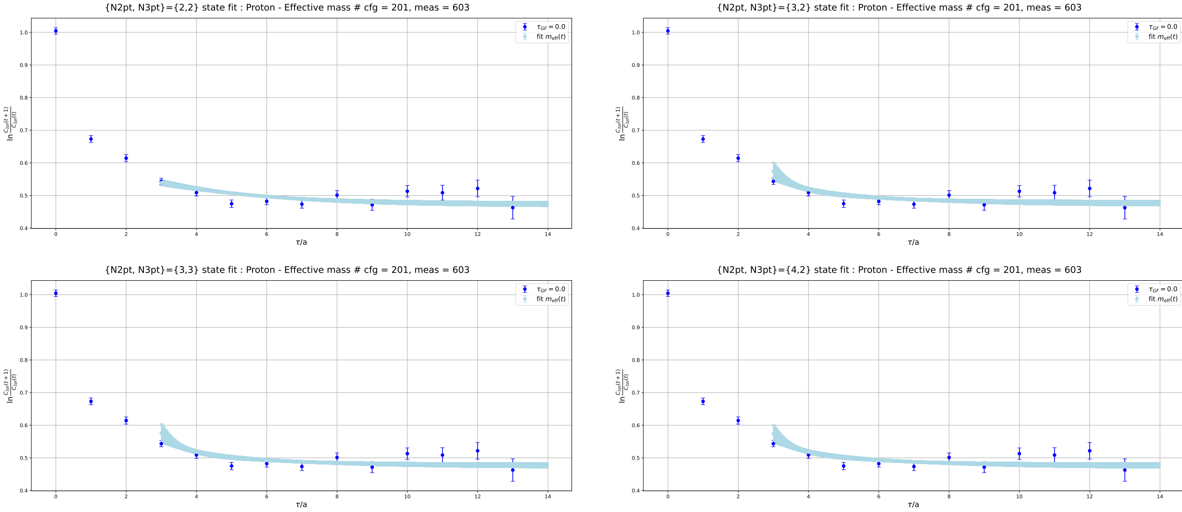


Figure 2: Effective Mass fit from $\{2,2\}$, $\{3,2\}$, $\{3,3\}$, $\{4,2\}$

Axial Charge: Fit With Different t_{sep} ($\tau_{GF} = 0$)

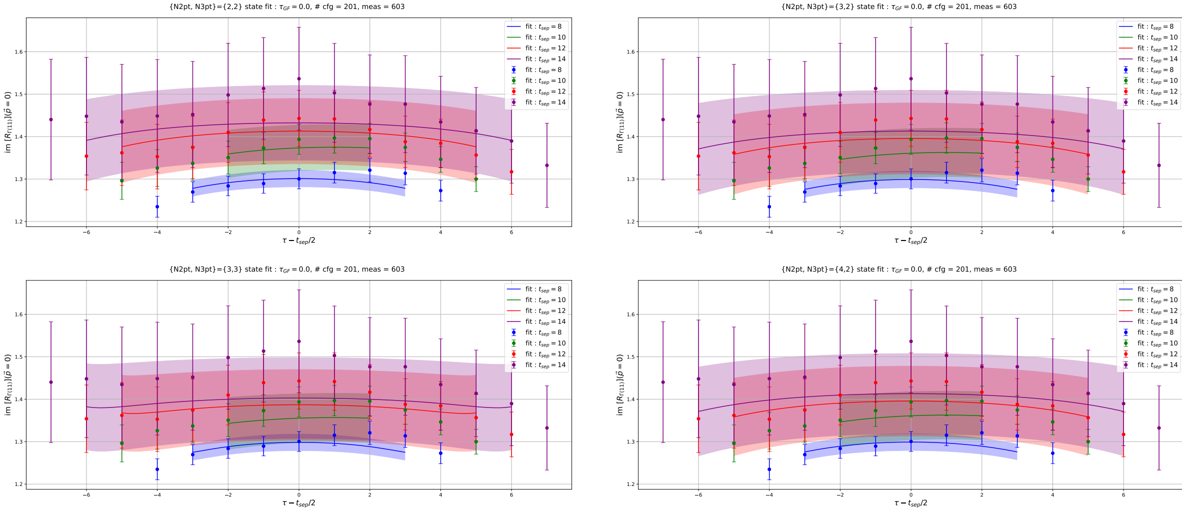


Figure 3: Axial Charge fit from $\{2, 2\}$, $\{3, 2\}$, $\{3, 3\}$, $\{4, 2\}$

Summery

Work in Progress$\cdots$

- Charges & Form Factors in QCD
- EDM due to Θ and ggg
- cEDM