

Strange quark electric dipole moment with topological anomalies

Chao-Qiang Geng, Xiang-Nan Jin, Chia-Wei Liu[✉], and Bin Wu[✉]

*School of Fundamental Physics and Mathematical Sciences, Hangzhou Institute for Advanced Study,
UCAS, Hangzhou 310024, China*



(Received 20 February 2026; accepted 5 June 2026; published 22 June 2026)

We introduce an anomaly-based framework to probe quark electric dipole moments in exclusive hadronic final states produced in e^+e^- annihilation. Chern-Simons-induced anomalous couplings yield a clean T-odd asymmetry A_T from interference between the Standard-Model amplitude and dipole-moment contributions. As an application, $\gamma^* \rightarrow K^+K^-\pi^0$ provides direct sensitivity to the strange-quark electric dipole moment (EDM) d_s in the chiral limit. The precision on d_s can reach $\mathcal{O}(10^{-16})$ e cm with existing CMD-3 data and $\mathcal{O}(10^{-18})$ e cm using current J/ψ samples at BESIII, improving the current direct constraint from hyperon EDM by two to three orders.

DOI: 10.1103/rvzl-q4t1

I. INTRODUCTION

Electric dipole moments (EDMs) of fermions are powerful null tests of CP violation [1,2]. In particular, nucleon EDM searches provide the most stringent constraints on CP violation in QCD, implying $|\bar{\theta}| \lesssim 10^{-10}$ for the QCD vacuum angle [3–5]. Complementary bounds on hadronic CP violation also arise from atomic EDM searches [6]. The electron EDM currently provides the tightest constraint on many classes of physics beyond the Standard Model (SM) [7–18].

Nevertheless, EDMs of quarks beyond the first generation remain subject to little direct experimental constraint [19–21]. In the lepton sector, the muon EDM provides a direct probe of CP violation beyond the first generation from storage-ring measurements [22]. Very recently, the BESIII collaboration reported the first direct bound on the Λ hyperon EDM from $J/\psi \rightarrow \Lambda \bar{\Lambda}$ [23]

$$|d_\Lambda(m_{J/\psi}^2)| < 6.5 \times 10^{-19} \text{ e cm}, \quad (1)$$

where $\sqrt{s}$ is the virtual-photon invariant mass. Because QCD is nonperturbative in this kinematic regime, a direct relation between $d_\Lambda(s)$ and the strange-quark EDM d_s can only be established at the zero-recoil point with an on-shell photon; $d_\Lambda(0) = d_s$ [3] in the quark model. Once the momentum dependence of $d_\Lambda(s)$ is taken into account, the resulting constraint on d_s becomes weaker, since

$d_\Lambda(m_{J/\psi}^2) = 5.29 \times 10^{-4} d_s \text{ e cm}$ [24]. On the other hand, the neutron EDM provides a significantly stronger indirect constraint [25,26]

$$d_n = -(1.5 \pm 0.7) \times 10^{-16} \bar{\theta} \text{ e cm} - (0.20 \pm 0.01) d_u + (0.78 \pm 0.03) d_d + (2.7 \pm 1.6) \times 10^{-3} d_s, \quad (2)$$

with the experimental bound $|d_n| < 1.8 \times 10^{-26} \text{ e cm}$ [5]. This relation makes it clear that d_n alone cannot disentangle the individual quark EDM contributions. In particular, cancellations among different terms may substantially weaken the direct sensitivity to d_s , thereby leaving room for the strange-quark EDM to be probed more effectively in complementary observables.

In this work we propose a new way to probe quark EDMs using CP -odd angular correlations in exclusive hadronic final states produced in e^+e^- annihilation. In particular, J/ψ may appear as the intermediate state as shown in Fig. 1. The key theoretical handle is the Wess-Zumino-Witten (WZW) term [27,28] in chiral perturbation theory (χ PT) [29–32], which encodes the QCD chiral anomaly and fixes the structure and normalization of odd-intrinsic-parity amplitudes [33–36]. This anomaly protection provides a controlled SM baseline against which dipole-induced CP violation can be isolated. As an application, we focus on the strange-quark EDM d_s in $e^+e^- \rightarrow K^+K^-\pi^0$, where the low-energy $\gamma^* \rightarrow K^+K^-\pi^0$ amplitude is governed by the WZW interaction and remains fixed in the chiral limit even with resonances included [37–40]. Compared with hyperon channels, this mode has a larger branching fraction and avoids the loss of efficiency from subsequent decays, which would otherwise reduce the effective statistics by $(\alpha_{\text{sub}} \mathcal{B}_{\text{sub}})^2$ [41,42], where $\mathcal{B}_{\text{sub}}$ and

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International license](#). Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

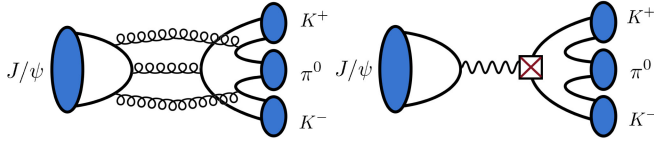


FIG. 1. Schematic structure of the amplitudes for $J/\psi \rightarrow \gamma^* \rightarrow K^+ K^- \pi^0$. The left panel shows the SM contribution mediated by three-gluon annihilation. The right panel shows the EDM-induced contribution, where the boxed cross denotes the strange-quark EDM.

α_{sub} denote the branching fraction and decay parameter of the subdecays.

This paper is organized as follows. In Sec. II, we construct the CP -odd observable. In Sec. III, we match d_s onto χ PT and present the WZW prediction. In Sec. IV, we estimate the attainable experimental precision. We conclude in Sec. V.

II. CP-ODD OBSERVABLES

As a concrete example, consider $\gamma^*(q) \rightarrow K^+(p_+) K^-(p_-) \pi^0(p_0)$. We employ the symmetric set of kinetic variables $s = q^2$, and $s_{\pm} = (p_{\pm} + p_0)^2$. The total amplitude is parametrized as

$$\mathcal{M}(\gamma^* \rightarrow K^+ K^- \pi^0) = \epsilon_{\mu}(q) J^{\mu}, \quad (3)$$

where ϵ is the polarization vector of the virtual photon. Using a gauge-invariant decomposition, the hadronic current is parametrized as

$$J^{\mu} = F_V(s, s_+, s_-) V^{\mu} + F_A(s, s_+, s_-) A^{\mu}, \quad (4)$$

with the vector and axial vectors are

$$V^{\mu} = \epsilon^{\mu\nu\alpha\beta} p_{+\nu} p_{-\alpha} p_{0\beta},$$

$$A^{\mu} = (q \cdot p_0)(p_+ - p_-)^{\mu} - (q \cdot (p_+ - p_-)) p_0^{\mu}, \quad (5)$$

respectively, satisfying $q_{\mu} V^{\mu} = q_{\mu} A^{\mu} = 0$. Invariance under C -parity demands F_V and F_A to be invariant under exchanging s_+ and s_- ; see the figures in Ref. [43]. Since the final state has negative intrinsic parity, one sees that F_V and F_A are CP -even and CP -odd, respectively.

In the γ^* center-of-mass frame, $q^{\mu} = (\sqrt{s}, \vec{0})$ and $p_i^{\mu} = (E_i, \vec{p}_i)$. The amplitude can be written as

$$\mathcal{M} = -\vec{\epsilon} \cdot (F_V \vec{V} + F_A \vec{A}), \quad (6)$$

with $\vec{A} = \sqrt{s}[E_0(\vec{p}_+ - \vec{p}_-) + (E_+ - E_-)(\vec{p}_+ + \vec{p}_-)]$ and $\vec{V} = \sqrt{s}(\vec{p}_- \times \vec{p}_+)$. Taking the unpolarized $e^+ e^- \rightarrow \gamma^*$ case, we sum over the two transverse photon polarizations and obtain

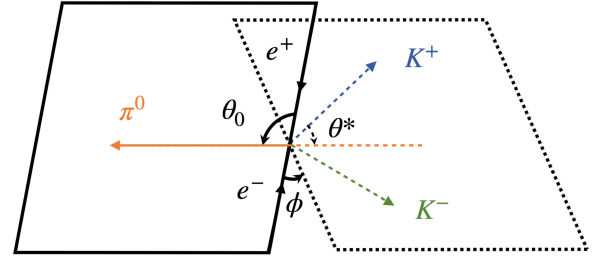


FIG. 2. The definitions of the angles in Eq. (8). The angle θ^0 is the angle between the momenta of e^+ and π^0 in the center-of-mass frame. The angle θ^* is the angle between π^0 and K^+ in the $K^+ K^-$ rest frame. The angle ϕ is the dihedral angle between the two planes spanned by (π^0, e^+, e^-) and (π^0, K^+, K^-) .

$$|\overline{\mathcal{M}}|^2 = \frac{1}{2} \sum_{\lambda=\pm} |\mathcal{M}|^2 = \frac{1}{2} |(F_V \vec{V} + F_A \vec{A})_{\perp}|^2, \quad (7)$$

where $\vec{X}_{\perp} \equiv \vec{X} - (\vec{X} \cdot \hat{z}) \hat{z}$ denotes the component transverse to the beam axis $\hat{z}$. The interference term proportional to $\Re(F_V F_A^*) \vec{V}_{\perp} \cdot \vec{A}_{\perp}$ is C -even and CP -odd, and will be shown to be induced by d_s .

The fully differential decay distribution in the γ^* rest frame is given by

$$\frac{d\Gamma}{d\Phi} = \frac{1}{(2\pi)^4} \frac{|\vec{p}_0| |\vec{p}_+^*|}{16s} |\overline{\mathcal{M}}|^2, \quad (8)$$

where $d\Phi \equiv dm_{KK} d\cos\theta_0 d\Omega_+^*$. Here, θ_0 is the inclined angle of the three-momentum of π^0 and γ^* polarization in the γ^* rest frame, while $\Omega_+^* = (\theta^*, \phi)$ and $\vec{p}_+^*$ are those of K^+ in the $K^+ K^-$ rest frame; see Fig. 2.

To isolate the contribution linear in F_A without assuming a specific kinematic dependence, we define the signed interference observable $\vec{V}_{\perp} \cdot \vec{A}_{\perp}$ and construct the counting asymmetry

$$A_T \equiv \frac{N(\vec{V}_{\perp} \cdot \vec{A}_{\perp} > 0) - N(\vec{V}_{\perp} \cdot \vec{A}_{\perp} < 0)}{N(\vec{V}_{\perp} \cdot \vec{A}_{\perp} > 0) + N(\vec{V}_{\perp} \cdot \vec{A}_{\perp} < 0)}. \quad (9)$$

Only the interference term $\Re(F_V F_A^*) (\vec{V}_{\perp} \cdot \vec{A}_{\perp})$ contributes, providing a robust probe of $\Re(F_V F_A^*)$. We note that A_T is C -even and P -odd, and $\gamma^* \rightarrow K^+ K^- \pi^0$ is self-conjugate under charge conjugation; hence A_T is a genuinely CP -odd observable. In particular, final-state interactions (with or without weak interactions at the leading order) cannot induce a nonzero A_T .

Finally, we note that the framework here can be straightforwardly extended to the one with J/ψ as the initial state, as its J^{PC} is identical to γ^* .

III. MAPPING d_s TO χ PT

We start from the flavor-conserving CP -violating strange-quark dipole operator

$$\mathcal{L}_{\text{eff}} = \frac{d_s}{2} \tilde{F}^{\mu\nu} \bar{s} \sigma_{\mu\nu} s, \quad (10)$$

where d_s parametrizes the strange-quark electric dipole operator induced by new physics. Although the d_n bound yields a much tighter indirect constraint on d_s , it is tied essentially to the $q^2 = 0$ limit. The process considered here instead probes the timelike region, where the relation between the measured hadronic observable and the underlying strange-quark EDM is intrinsically q^2 -dependent. An observable effect in this channel would therefore provide genuinely new information, while at the same time testing whether nontrivial cancellations among CP -odd contributions to d_n [44].

The QCD Lagrangian enjoys a global $SU(3)_L \otimes SU(3)_R$ flavor symmetry, with $q_L \rightarrow g_L q_L$, $q_R \rightarrow g_R q_R$ and $q = (u, d, s)^T$. To make chiral symmetry manifest in the presence of d_s , we introduce spurion fields $t_\pm$ that transforming as $t \rightarrow g_L t g_R^\dagger$. With these spurions, the effective Lagrangian can be rewritten as

$$\mathcal{L}_{\text{eff}} = \frac{d_s}{2} \tilde{F}^{\mu\nu} (\bar{q}_L \sigma_{\mu\nu} t q_R + \bar{q}_R \sigma_{\mu\nu} t^\dagger q_L). \quad (11)$$

Matching to Eq. (10), one sets $t = \text{diag}(0, 0, 1)$ in flavor space. With the spurion transformation rules, $\mathcal{L}_{\text{eff}}$ is formally invariant under $SU(3)_L \otimes SU(3)_R$.

To match onto chiral fields, we note that $U = \exp(2iP/f_P)$ with

$$P = \begin{bmatrix} \frac{\eta_0}{\sqrt{3}} + \frac{\eta_8}{\sqrt{6}} + \frac{\pi^0}{\sqrt{2}} & \pi^+ & K^+ \\ \pi^- & \frac{\eta_0}{\sqrt{3}} + \frac{\eta_8}{\sqrt{6}} - \frac{\pi^0}{\sqrt{2}} & K^0 \\ K^- & \bar{K}^0 & \frac{\eta_0}{\sqrt{3}} - \frac{2\eta_8}{\sqrt{6}} \end{bmatrix}, \quad (12)$$

which transforms as $U \rightarrow g_L U g_R^\dagger$. The lowest-order chiral realization of the strange-quark tensor currents is

$$\begin{aligned} \mathcal{L} &= \frac{d_s}{2} \tilde{F}^{\mu\nu} \frac{\Lambda_2}{8} T_{\mu\nu}(t) \\ T_{\mu\nu}(t) &= i\text{Tr}[(tU^\dagger + Ut^\dagger)[\alpha_\mu, \alpha_\nu]] \\ &\quad - \epsilon_{\mu\nu\lambda\sigma} \text{Tr}[(tU^\dagger - Ut^\dagger)\alpha^\lambda \alpha^\sigma], \end{aligned} \quad (13)$$

where $\alpha_\mu = -i\partial_\mu U U^\dagger$ and Λ_2 is a low-energy constant. The tensor $T_{\mu\nu}$ is constructed to satisfy $\frac{1}{2}\epsilon^{\mu\nu\alpha\beta} T_{\alpha\beta}(t) = T^{\mu\nu}(-it)$, mirroring $\frac{1}{2}\epsilon^{\mu\nu\alpha\beta} \sigma_{\alpha\beta} = -i\sigma^{\mu\nu}\gamma_5$ [45].

Ideally, one searches for channels in which the chiral approximation is reliable and all incoming and outgoing states are on shell. Imposing flavor conservation, the simplest such radiative channel is $\eta^{(\prime)} \rightarrow \pi^+ \pi^- \gamma$ [46–48] at leading order. Nevertheless, C -parity symmetry forces the $\pi^+ \pi^-$ system to be an isovector, and hence d_s cannot induce this decay in the isospin limit. We therefore turn to $\gamma^* \rightarrow K^+ K^- \pi^0$ instead. The relevant piece reads

$$\mathcal{L}^{(3)} = \frac{-id_s \Lambda_2}{\sqrt{2} f_K^2 f_\pi} F^{\mu\nu} [K^+ (\partial_\mu K^-) (\partial_\nu \pi^0)] + \text{H.c.}, \quad (14)$$

where $F^{\mu\nu} = (\partial^\mu A^\nu - \partial^\nu A^\mu)$. We focus on this channel because $K^\pm$ is on shell and π^0 can be reconstructed from two photons in relevant experiments.

In the SM, the leading contribution is fixed by the WZW term. Since this term is the descendant of a five-dimensional Chern-Simons form, it is naturally written in the language of differential forms [27,28]

$$\mathcal{L}_{\chi\text{PT}}^{(\text{WZW})} = \frac{N_c}{48\pi^2} \text{Tr}[eQA \wedge \alpha^3]. \quad (15)$$

Here A stands for the photon field. The overall normalization is fixed by the underlying chiral anomaly. Equivalently, it is topologically quantized by the five-dimensional Chern-Simons construction, so the coefficient is an integer, which is N_c in QCD. Expanding to $\mathcal{O}(f_P^{-3})$ yields the familiar local $\gamma^* \rightarrow K^+ K^- \pi^0$ vertex

$$\mathcal{L}_{\chi\text{PT}}^{\gamma KK\pi^0} = \frac{\sqrt{2} N_c e}{12\pi^2 f_K^2 f_\pi} A \wedge d\pi^0 \wedge dK^+ \wedge dK^-. \quad (16)$$

The WZW-induced amplitude therefore provides the irreducible odd-parity baseline for $\gamma^* \rightarrow K^+ K^- \pi^0$. The interference between the WZW term and the EDM naturally occurs because both induce odd-intrinsic-parity processes. In contrast, ordinary χ PT conserves intrinsic parity and does not contribute.

In the chiral limit, χ PT yields

$$F_V(0, 0, 0) = \frac{\sqrt{2} N_c e}{12\pi^2 f_K^2 f_\pi}, \quad F_A(0, 0, 0) = \frac{-d_s \Lambda_2}{\sqrt{2} f_K^2 f_\pi}. \quad (17)$$

Ideally, one can test this mechanism via scattering processes such as $K^+ K^- \rightarrow \pi^0 \gamma$. However, to construct a CP -odd observable, one must measure the photon polarization, and the required initial state is difficult to produce experimentally. Hence, we consider the process $\gamma^* \rightarrow K^+ K^- \pi^0$ with $\sqrt{s} > 2m_K + m_\pi$. Point-by-point extractions of F_V and F_A can be performed experimentally. Here we approximate

them using a vector-meson-dominance (VMD) model [49] based on the sequential decays $\gamma^* \rightarrow K^+ K^{*-}(892) + \text{c.c.}$, given by

$$F_{V,A}(s, s_+, s_-) = F_{V,A}(s, 0, 0) V(s_+, s_-),$$

$$V(s_+, s_-) = \frac{1}{2} \left(\frac{m_{K^*}^2}{m_{K^*}^2 - s_+ - i\Gamma_{K^*} \sqrt{s_+}} + \frac{m_{K^*}^2}{m_{K^*}^2 - s_- - i\Gamma_{K^*} \sqrt{s_-}} \right), \quad (18)$$

where $F_V(s, 0, 0)$ is a nonperturbative input to be fitted from the experiments. The determination of the ratio $F_A(s, 0, 0)/F_V(s, 0, 0)$ needs to await lattice input, and here we take it as a constant in $\gamma^* \rightarrow K^+ K^- \pi^0$.

For $J/\psi \rightarrow K^+ K^- \pi^0$, the leading form factor F_V is induced by three-gluon annihilation, and the ratio F_A/F_V is no longer given by Eq. (17). The leading form factor $F_V(s, 0, 0)$ is extracted from the experimental branching fraction, while we adopt the resonant structure

$$F_A(m_{J/\psi}^2, 0, 0) = \frac{e Q_c f_{J/\psi}}{m_{J/\psi}} \frac{m_A^2}{m_A^2 - m_{J/\psi}^2} \frac{d_s \Lambda_2}{\sqrt{2} f_K^2 f_\pi}, \quad (19)$$

where $Q_c = 2/3$ is the charge of the charm quark, $f_{J/\psi} = (410.4 \pm 1.7)$ MeV is the decay constant of J/ψ [50], and m_A is the nominal axial-vector meson mass and we vary it from 1 GeV to 2 GeV as an uncertainty of our framework. The dependencies on s_+ and s_- can be well described by the VMD in Eq. (18). In particular, the BESIII collaboration found that more than 95% of the $\mathcal{B}(J/\psi \rightarrow K^+ K^- \pi^0)$ is made of the sequential decay $J/\psi \rightarrow K^+ K^{*-} + \text{c.c.}$ [43].

IV. NUMERICAL RESULT

We propose to probe the EDM-induced interference in $e^+ e^- \rightarrow K^+ K^- \pi^0$. For a sample of N selected signal events, the purely statistical uncertainty is

$$\delta A_T \simeq \sqrt{\frac{1 - A_T^2}{N}} \approx \frac{1}{\sqrt{N}} \quad (|A_T| \ll 1). \quad (20)$$

The relation between A_T and d_s is denoted as

$$d_s = \frac{-N_c e}{6\pi^2 \Lambda_2} \frac{A_T}{R(s)}. \quad (21)$$

Here, $R(s)$ is a phase factor that gauges our ability to extract d_s from A_T , and it can be computed numerically from Eq. (9) once the form factors are provided. Taking F_A and F_V to be saturated by VMD, we have

$$R(s) = r_V \frac{\int d\Phi |\vec{p}_0| |\vec{p}_+^*| |V(s_+, s_-)|^2 |\vec{V}_\perp \cdot \vec{A}_\perp|}{\int d\Phi |\vec{p}_0| |\vec{p}_+^*| |V(s_+, s_-)|^2 |\vec{V}_\perp \cdot \vec{V}_\perp|}. \quad (22)$$

Here, r_V denotes the factor associated with $V \rightarrow K^+ K^- \pi^0$, where $V = \gamma^*, J/\psi$, given by

$$r_{\gamma^*} = 2,$$

$$r_{J/\psi} = \frac{-\sqrt{2} N_c e^2 Q_c f_{J/\psi} m_A^2}{6\pi^2 f_K^2 f_\pi m_{J/\psi} F_V(m_{J/\psi}^2, 0, 0) (m_A^2 - m_{J/\psi}^2)}. \quad (23)$$

We take $\Lambda_2 = (18.20_{-0.99}^{+1.17}) \times 10^{-3}$ GeV [51].

For CMD-3, the analyzed sample ($L \simeq 34$ pb $^{-1}$) [52,53] and the measured Born cross section $\sigma(e^+ e^- \rightarrow K^+ K^- \pi^0) \sim 0.05\text{--}1.1$ nb in $\sqrt{s} = 1.4\text{--}2.0$ GeV [53] imply $N_{\text{prod}} \sim (1.7 \times 10^3 - 2.7 \times 10^4)$. The VMD behavior in Eq. (18) gives $R(s) \approx 1.56$ at 1.6 GeV. Taking the accumulated events over q^2 with $N_{\text{tot}} \sim 10^3\text{--}10^4$, we obtain

$$\delta d_s \sim (3.5 - 11.1) \times 10^{-16} \text{ e cm}. \quad (24)$$

A prospective VEPP-2000 dataset of $L \sim 1$ fb $^{-1}$ [54] would improved the precision by an order. The overall normalization largely cancels in asymmetries.

Similar measurements can be performed at BESIII at center-of-mass energies from 2.00 to 3.08 GeV, where the dominant resonance is expected to be the K_2^* [55], with an event yield of $\mathcal{O}(10^5)$. Here we focus on the J/ψ resonance, which is dominated by the $K^*(892)$ contribution [43]. From the branching fraction of $\mathcal{B}(J/\psi \rightarrow K^+ K^- \pi^0) = (2.88 \pm 0.01 \pm 0.12) \times 10^{-3}$, we find $F_V(m_{J/\psi}^2, 0, 0) = -3.7 \times 10^{-2} \text{ GeV}^{-3}$ which gives $R(s) = 0.8 - 4.7$ from Eq. (19). At BESIII, 182,972 candidate events were observed in 2019 [43], and the J/ψ sample size has increased by about 45 times since 2009 [56]. The expected statistical precision of A_T is around 3.5×10^{-4} , resulting in

$$\delta d_s \simeq (4.2 - 26.9) \times 10^{-18} \text{ e cm}, \quad (25)$$

which can achieve better constraint than the one from d_A in Eq. (1) and ready to be tested at BESIII. The uncertainties of d_s quoted here mainly comes from the unknown of F_A . A solid determination can be achieved once $F_A(s, 0, 0)$ is determined from a nonperturbative method. In general, all CP violating operators may contribute to F_A and may affect the constraint for d_s here. Other CP -odd operators without explicit photon fields (e.g., chromo-EDMs and four-quark operators) mainly affect purely hadronic observables; here we set them to zero and focus on the EDM as a minimal benchmark, leaving a systematic EFT treatment to future work.

The Super τ Charm Facility will increase the production of J/ψ by two orders of magnitude [57,58] and δd_s by an order of magnitude, comparing to BESIII. In addition, a border region of $\sqrt{s}$ range from 4.2 GeV to 7 GeV may be explored. Complementary information may also come from Belle II, which aims at an integrated luminosity

of $\sim 50 \text{ ab}^{-1}$ [59], providing enormous statistics for exclusive hadronic final states with excellent charged-hadron identification. Although Belle II operates near the $\Upsilon(4S)$, the initial-state-radiation program enables contiguous measurements of $e^+e^- \rightarrow K^+K^-\pi^0$ over a wide effective $\sqrt{s}$ range [60].

V. CONCLUSION

We develop a framework to probe quark EDMs in $\gamma^* \rightarrow PPP$ using a gauge-invariant basis of transverse vectors. We identify the CP -odd and T -odd interference observable $\vec{V}_\perp \cdot \vec{A}_\perp$. As a concrete example, we study $\gamma^* \rightarrow K^+K^-\pi^0$ in χ PT: the vector current is fixed by the WZW term, while the EDM induces an axial current through the chiral realization of d_s . We match the quark-level operator onto chiral fields and obtain the form factors in the chiral limit. The resulting interference between the topological anomaly and d_s is expected, since both violate intrinsic parity.

We define a signed interference asymmetry A_T to isolate EDM effects and compute the statistical sensitivity δd_s . Numerical estimates for $e^+e^- \rightarrow K^+K^-\pi^0$ at CMD-3

indicate sensitivities at the 10^{-16} e cm level with current data, and improved reach with higher luminosity. We extend the analysis to $J/\psi \rightarrow K^+K^-\pi^0$ and note that existing BESIII samples could reach $\delta d_s \sim 10^{-18} \text{ e cm}$. The asymmetry construction cancels the overall normalization and mitigates many detector effects, providing a robust strategy to probe CP -violating strange-quark dipole moments.

ACKNOWLEDGMENTS

This work is supported in part by the National Natural Science Foundation of China (NSFC) under Grants No. 12547104 and No. 12575096; and the China Postdoctoral Science Foundation (CPSF) under Grant No. 2025M773361.

DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

-
- [1] J. H. Smith, E. M. Purcell, and N. F. Ramsey, *Phys. Rev.* **108**, 120 (1957).
 - [2] W. Bernreuther and M. Suzuki, *Rev. Mod. Phys.* **63**, 313 (1991); **64**, 633(E) (1992).
 - [3] X. G. He, B. H. J. McKellar, and S. Pakvasa, *Int. J. Mod. Phys. A* **04**, 5011 (1989); **06**, 1063(E) (1991).
 - [4] T. E. Chupp, P. Fierlinger, M. J. Ramsey-Musolf, and J. T. Singh, *Rev. Mod. Phys.* **91**, 015001 (2019).
 - [5] C. Abel, S. Afach, N. J. Ayres, C. A. Baker, G. Ban, G. Bison, K. Bodek, V. Bondar, M. Burghoff, E. Chandel *et al.*, *Phys. Rev. Lett.* **124**, 081803 (2020).
 - [6] S. K. Lamoreaux, J. P. Jacobs, B. R. Heckel, F. J. Raab, and N. Fortson, *Phys. Rev. Lett.* **59**, 2275 (1987).
 - [7] S. M. Barr and A. Zee, *Phys. Rev. Lett.* **65**, 21 (1990); **65**, 2920(E) (1990).
 - [8] X. G. He, B. H. J. McKellar, and S. Pakvasa, *Phys. Lett. B* **283**, 348 (1992).
 - [9] M. J. Ramsey-Musolf and S. Su, *Phys. Rep.* **456**, 1 (2008).
 - [10] Y. Li, S. Profumo, and M. Ramsey-Musolf, *J. High Energy Phys.* **08** (2010) 062.
 - [11] J. Shu and Y. Zhang, *Phys. Rev. Lett.* **111**, 091801 (2013).
 - [12] M. Jung and A. Pich, *J. High Energy Phys.* **04** (2014) 076.
 - [13] S. Inoue, M. J. Ramsey-Musolf, and Y. Zhang, *Phys. Rev. D* **89**, 115023 (2014).
 - [14] I. Doršner, S. Fajfer, A. Greljo, J. F. Kamenik, and N. Košnik, *Phys. Rep.* **641**, 1 (2016).
 - [15] K. Fuyuto, M. Ramsey-Musolf, and T. Shen, *Phys. Lett. B* **788**, 52 (2019).
 - [16] W. Dekens, J. de Vries, M. Jung, and K. K. Vos, *J. High Energy Phys.* **01** (2019) 069.
 - [17] S. Li, Y. Xiao, and J. M. Yang, *Nucl. Phys.* **B974**, 115629 (2022).
 - [18] T. S. Roussy, L. Caldwell, T. Wright, W. B. Cairncross, Y. Shagam, K. B. Ng, N. Schlossberger, S. Y. Park, A. Wang, J. Ye *et al.*, *Science* **381**, adg4084 (2023).
 - [19] K. Inami *et al.* (Belle Collaboration), *J. High Energy Phys.* **04** (2022) 110.
 - [20] L. Vecchi, *J. High Energy Phys.* **09** (2025) 064.
 - [21] Z. L. Huang, X. Y. Du, X. G. He, C. W. Liu, and Z. Y. Zou, *Chin. Phys. Lett.* **43**, 030201 (2026).
 - [22] G. W. Bennett *et al.* (Muon (g-2) Collaboration), *Phys. Rev. D* **80**, 052008 (2009).
 - [23] M. Ablikim *et al.* (BESIII Collaboration), *arXiv:2506.19180*.
 - [24] K. B. Chen, X. G. He, J. P. Ma, and X. B. Tong, *Phys. Rev. Lett.* **136**, 051902 (2026).
 - [25] R. Gupta, B. Yoon, T. Bhattacharya, V. Cirigliano, Y. C. Jang, and H. W. Lin, *Phys. Rev. D* **98**, 091501(R) (2018).
 - [26] J. Dragos, T. Luu, A. Shindler, J. de Vries, and A. Yousif, *Phys. Rev. C* **103**, 015202 (2021).
 - [27] J. Wess and B. Zumino, *Phys. Lett.* **37B**, 95 (1971).
 - [28] E. Witten, *Nucl. Phys.* **B223**, 422 (1983).
 - [29] S. Weinberg, *Physica A* **96**, 327 (1979).
 - [30] J. Gasser and H. Leutwyler, *Ann. Phys. (N.Y.)* **158**, 142 (1984).
 - [31] J. Gasser and H. Leutwyler, *Nucl. Phys.* **B250**, 465 (1985).
 - [32] H. Leutwyler, *Ann. Phys. (N.Y.)* **235**, 165 (1994).
 - [33] O. Kaymakçalan, S. Rajeev, and J. Schechter, *Phys. Rev. D* **30**, 594 (1984).
 - [34] K. c. Chou, H. y. Guo, K. Wu, and X. c. Song, *Phys. Lett. B* **134**, 67 (1984).

- [35] T. Fujiwara, T. Kugo, H. Terao, S. Uehara, and K. Yamawaki, *Prog. Theor. Phys.* **73**, 926 (1985).
- [36] Y. L. Wu, Y. B. Xie, and G. Z. Zhou, *Mod. Phys. Lett. A* **01**, 23 (1986).
- [37] M. Harada and K. Yamawaki, *Phys. Rep.* **381**, 1 (2003).
- [38] M. Dax, D. Stamen, and B. Kubis, *Eur. Phys. J. C* **81**, 221 (2021).
- [39] C. Q. Geng, C. W. Liu, and Y. L. Wu, *J. High Energy Phys.* **07** (2025) 119.
- [40] C. Q. Geng, C. W. Liu, and Y. L. Wu, *arXiv:2601.03740*.
- [41] J. Fu, H. B. Li, J. P. Wang, F. S. Yu, and J. Zhang, *Phys. Rev. D* **108**, L091301 (2023).
- [42] Y. Du, X. G. He, J. P. Ma, and X. Y. Du, *Phys. Rev. D* **110**, 076019 (2024).
- [43] M. Ablikim *et al.* (BESIII Collaboration), *Phys. Rev. D* **100**, 032004 (2019).
- [44] T. Bhattacharya, V. Cirigliano, R. Gupta, H. W. Lin, and B. Yoon, *Phys. Rev. Lett.* **115**, 212002 (2015).
- [45] O. Cata and V. Mateu, *J. High Energy Phys.* **09** (2007) 078.
- [46] C. Q. Geng, J. N. Ng, and T. H. Wu, *Mod. Phys. Lett. A* **17**, 1489 (2002).
- [47] P. Sanchez-Puertas, *J. High Energy Phys.* **01** (2019) 031.
- [48] L. Gan, B. Kubis, E. Passemar, and S. Tulin, *Phys. Rep.* **945**, 1 (2022).
- [49] J. J. Sakurai, *Ann. Phys. (N.Y.)* **11**, 1 (1960).
- [50] D. Hatton *et al.* (HPQCD Collaboration), *Phys. Rev. D* **102**, 054511 (2020).
- [51] S. Z. Jiang, Q. Wang, and Y. Zhang, *Phys. Rev. D* **87**, 094014 (2013).
- [52] B. Khazin, *Nucl. Phys. B, Proc. Suppl.* **181–182**, 376 (2008).
- [53] A. L. Erofeev (CMD–3 Collaboration), *Proc. Sci. EPS-HEP2017* (2017) 252.
- [54] S. A. Rastigeev, Y. S. Aktershev, E. V. Bedarev, O. V. Belikov, D. E. Berkaev, V. S. Denisov, E. A. Eminov, L. B. Fomin, K. M. Gorchakov, G. V. Karpov *et al.*, *Instrum. Exp. Tech.* **67**, S39 (2024).
- [55] M. Ablikim *et al.* (BESIII Collaboration), *J. High Energy Phys.* **07** (2022) 045.
- [56] M. Ablikim *et al.* (BESIII Collaboration), *Chin. Phys. C* **46**, 074001 (2022).
- [57] Q. Zhang *et al.*, *Front. Phys. (Beijing)* **19**, 14701 (2024).
- [58] X. C. Ai *et al.*, *Nucl. Sci. Tech.* **36**, 242 (2025).
- [59] SuperKEKB/Belle II, Luminosity Projection (Dec 2024), Online PDF, available at https://public.belle2.org/downloads/LuminosityProjection_2024Dec.pdf.
- [60] B. Aubert *et al.* (BABAR Collaboration), *Phys. Rev. D* **77**, 092002 (2008).