Generalized Global Symmetries and Nonperturbative Quantum Flavodynamics

Seth Koren

(they/them or whatever)

University of Notre Dame

Leptons: 2211.07639 with Clay Córdova, Sungwoo Hong, Kantaro Ohmori

Quarks: 2402.12453 with Clay & Sungwoo

Related ideas in my

Naturalness: 2009.11870
SM proton stability: 2204.01741
Discrete (B-L) for the lithium problem: 2204.01750
SM flavor 2-group: 2212.13193 with Clay
SM 1-form symmetry: 2406.17850 with A. Martin
Non-inv symmetry in PQWW: 24XX with A. Delgado
Non-inv symmetry in DSFZ: 24XX with Sungwoo & G. Choi
SM 2-form symmetry: 25XX with Sungwoo
Lots more coming

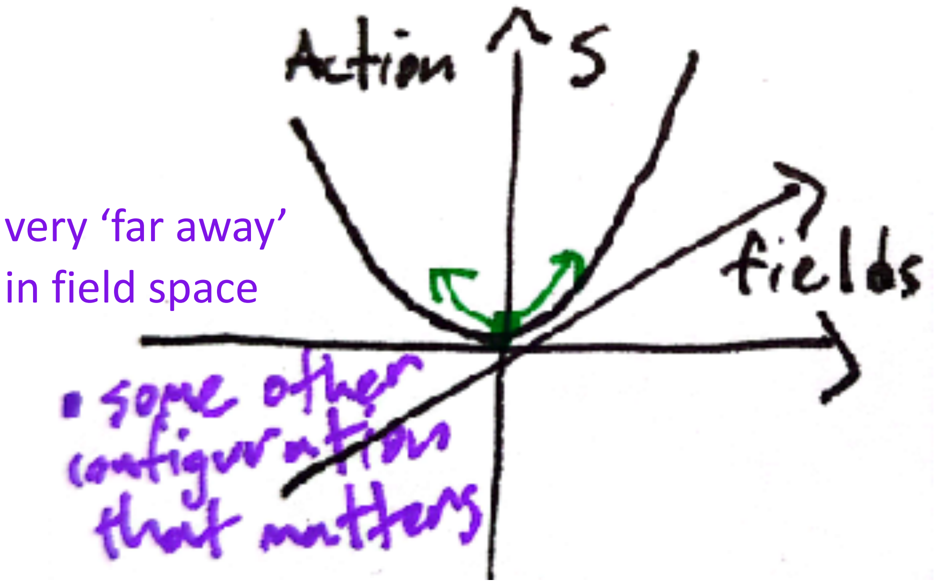
Talk outline

- Motivating generalized symmetries
  - Breaking generalized symmetries
  - Non-invertible symmetries
- Gauging lepton flavor
  - Non-invertible symmetry for Dirac neutrinos
- Gauging quark flavor
  - Non-invertible symmetry for strong CP

Nonperturbative QFT effects

Most of our time spent understanding *perturbative* QFT effects, sensibly

Expand around vacuum, calculate e.g. $\langle \psi_{\text{out}} | S | \psi_{\text{in}} \rangle$

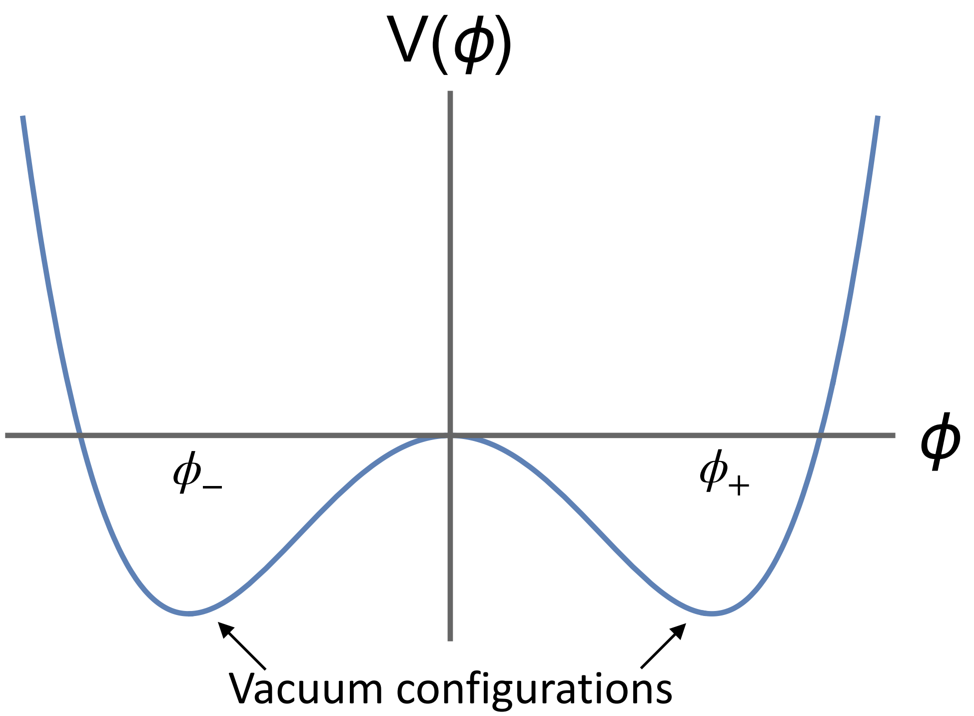


All great. But quantum field theory
is richer than perturbation theory!

Topology in field theory

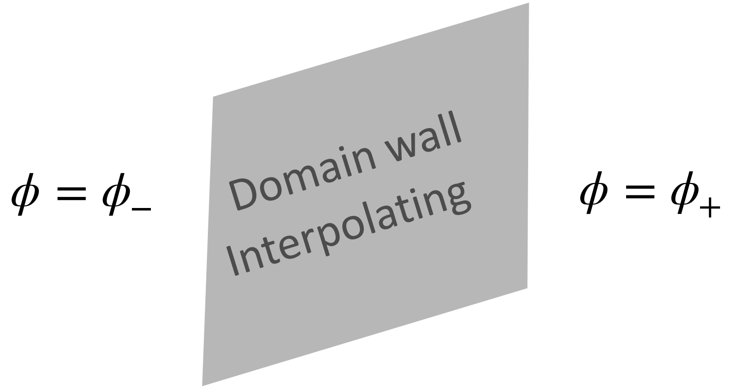
Often there are 'topological quantum numbers' that classify field space

A $\mathbb{Z}_2$-symmetric scalar breaking $\mathbb{Z}_2 \to \varnothing$, distinct vacua $\pi_0(\mathscr{M}_{\text{vac}}) = \mathbb{Z}_2$



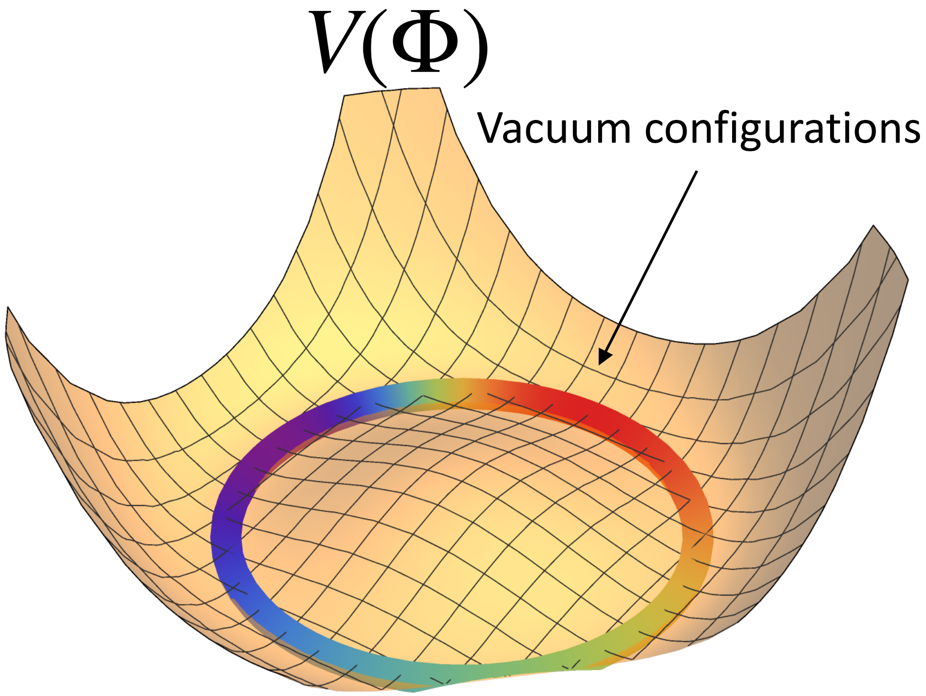
Local vacuum solutions

$\phi(x) : \mathbb{R}^4 \to \mathscr{M}_{\text{vac}}$ can have defects; domain walls separate different regions

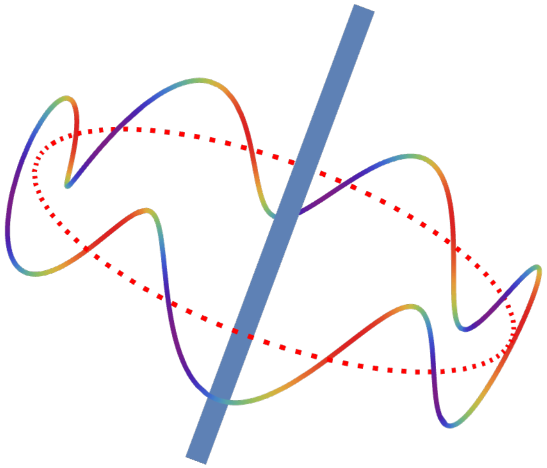


Higher-dimensional topology

A $U(1)$-symmetric scalar breaking $U(1) \to \varnothing$,
$\pi_0(\mathscr{M}_{\text{vac}}) = 1$, $\pi_1(\mathscr{M}_{\text{vac}}) = \mathbb{Z}$



$\Phi(x) : \mathbb{R}^4 \to \mathscr{M}_{\text{vac}}$ can have winding number which leads to cosmic strings



B-L strings for the lithium problem 2204.01750

I'll mention later π_2 for monopoles and π_3 for instantons but these are harder to visualize

Generalized Global Symmetries

Symmetries are important!

Usually look at Lagrangian data and consider transforming local operators

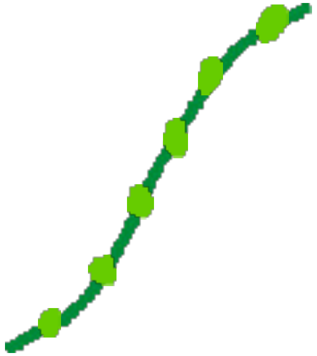
$$\psi^a(x) \to R^a{}_b \psi^b(x)$$

But what about these extended operators associated to this nonperturbative, topological data in our theory?

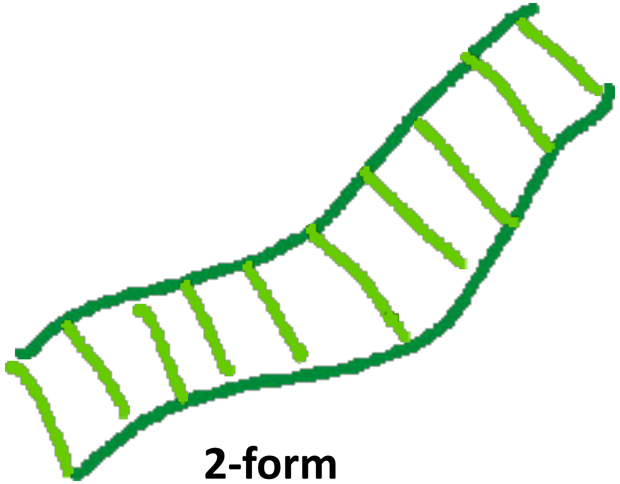
GGS Framework
Gaiotto, Kapustin, Seiberg, Willett
1412.5148

Higher-form symmetries

0-form symmetry
charged local operators
e.g. particles



1-form
line operators
e.g. Wilson line



2-form
surface operators
e.g. cosmic string

3-form
volume operators
e.g. domain wall

$$\partial_\mu J^\mu = 0$$

$$\partial_\mu J^{\mu\nu} = 0$$

Generally $\partial_\mu J^{\mu_1 \mu_2 \cdots \mu_{p+1}} = 0$ antisymmetric

Break by adding charged operator to Lagrangian e.g. $\delta\mathscr{L} = MNN$

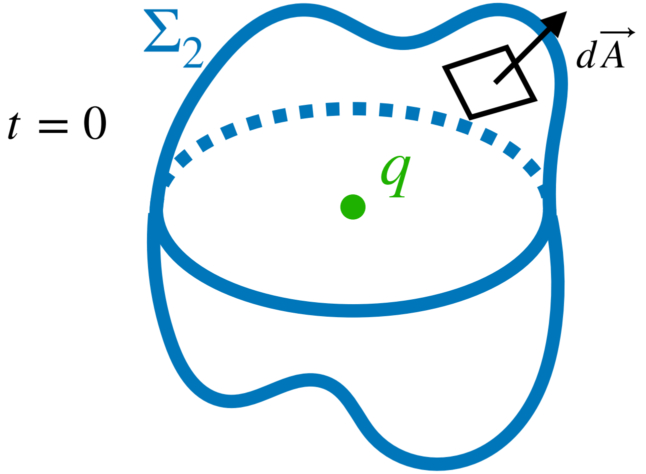
Break only with the appearance of new dynamical degrees of freedom!

Generalized Global Symmetry of Electromagnetism

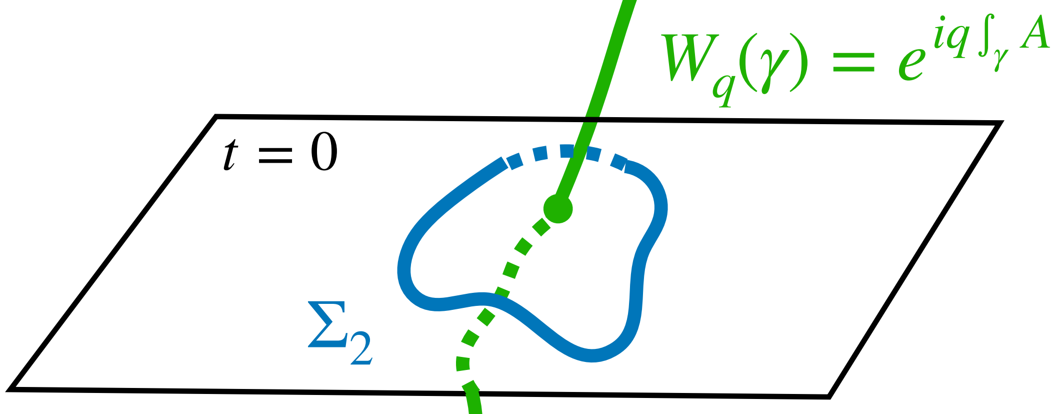
Recall Gauss' law: The Gaussian surface is topological and so computes an invariant charge.

In pure electromagnetism, the photon field strength is conserved $J_E^{\mu\nu} \sim \frac{1}{e^2} F^{\mu\nu}, \quad \partial_\mu J_E^{\mu\nu} = 0$

Gauss' law computes a Noether charge for an electric 1-form symmetry!



$$q = \int_{\Sigma_2} \vec{E} \cdot d\vec{A}$$

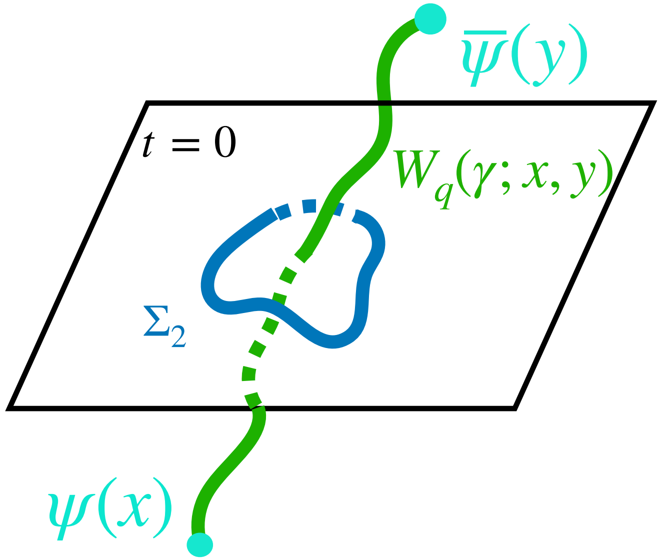


$$q = \int_{\Sigma_2} \varepsilon_{\mu\nu\rho\sigma} F^{\mu\nu} dS^{\rho\sigma}$$

Emergent 1-form symmetry

The 1-form symmetry is emergent in the low-energy, long-distance theory $E \ll m_e$.

Once we see the dynamical electron, then Wilson lines can 'end'.



That is, Gauss' law really breaks for $E > m_e$ because the Gaussian surface is no longer topological.

Or does a discrete 1-form symmetry remain? Test at LHC! SK & A. Martin 2406.17850

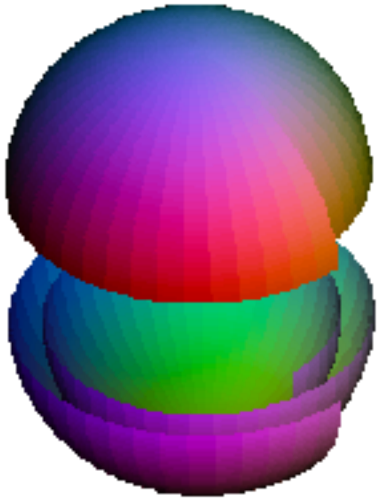
Mutatis mutandis a magnetic one-form symmetry for a theory H with 't Hooft lines classified by $\pi_1(H)$

Instantons

Yang-Mills can have topological quantum numbers $\pi_3(SU(N)) = \mathbb{Z}$

There are other vacuum configurations called ‘instantons’

Semiclassically we must sum over *all* saddle points



$$\int \mathcal{D}\mathscr{A} e^{-S} \simeq \sum_{\mathscr{A}=A^{(n)}+A} e^{-S_{inst}} \int \mathcal{D}A e^{-S}$$

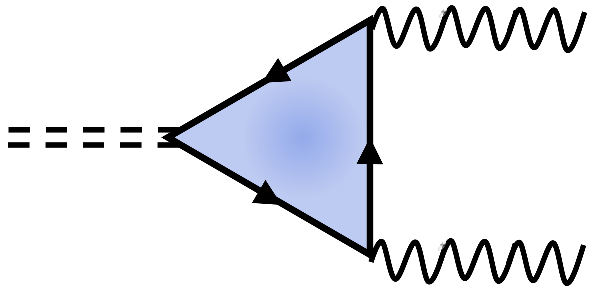
Relative suppression is by instanton action $S_{inst} \geq \frac{1}{4}\int_{\mathbb{R}^4} F^{\mu\nu}\tilde{F}_{\mu\nu} = \frac{8\pi^2}{g^2}n$

Naturally exponentially suppressed effects!

Instantons and Anomalies

Sometimes a global symmetry, say $U(1)_X$, can be good classically but quantum-mechanically be anomalous

$$\partial_\mu J^\mu_X = 0 \quad \longrightarrow \quad \partial_\mu J^\mu_X = \frac{\mathscr{A}}{8\pi^2} F^{\mu\nu}\tilde{F}_{\mu\nu}$$



What sorts of configurations can 'saturate' this anomaly and actually violate $U(1)_X$?

It's the instantons which carry this quantum number

$\int_{\mathbb{R}^4} F^{\mu\nu}\tilde{F}_{\mu\nu} \propto \int_{\partial\mathbb{R}^4 \simeq S^3} \hat{n}_\mu J^\mu_{\mathrm{CS}}$ = number of times A_μ 'winds' around infinity

How do instantons violate symmetries?

Atiyah-Singer '63
Fujikawa '80

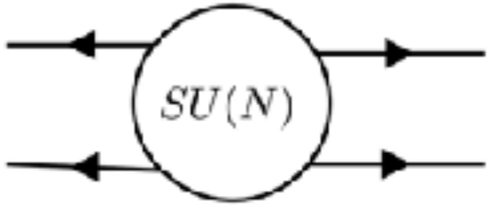
For deep topological reasons charged fermions in G-instanton backgrounds have zero-modes $i\gamma_\mu D^\mu_{\text{inst}}\psi^0 = 0$

Striking, since no action cost to excite these $S_{Dirac} = \int_{\mathcal{M}} i\bar{\psi}\gamma_\mu D^\mu \psi$

't Hooft '76

This leads to effective operators known as 't Hooft vertices which violate G-anomalous symmetries

Multiplicity of ψ_i legs given by Dirac index I_{ψ_i}



Violates anomalous $U(1)_X$ by $\mathscr{A}_X = \sum_{\psi_i} q_{\psi_i} I_{\psi_i}$

Unsaturated Anomalies - Missing Instantons

We said instantons are the field configurations which can saturate the anomaly

$$\partial_\mu J^\mu_X = \frac{\mathscr{A}}{8\pi^2} F^{\mu\nu}\tilde{F}_{\mu\nu}$$

But what about when they don't?

E.g. famously $\pi_3\big(U(1)\big) = 1$ and *there are no Abelian instantons in* $\mathbb{R}^4$, so $\int_{\mathbb{R}^4} F\tilde{F} = 0$

Old lesson: X is anomalous but S-matrix preserves X anyway

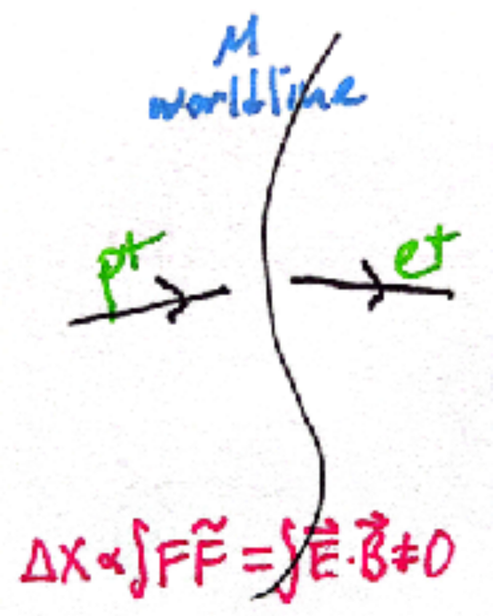
EFT philosophy: If there is ever a zero, there should be a symmetry!

Somehow despite X being anomalous there must remain a subtle sort of symmetry that demands the S-matrix preserves X

A hint: X *can* be violated around magnetic monopoles

c.f. Callan-Rubakov

Dirac ’31
Callan, Rubakov ‘80s
Ongoing…



A confused effective field theorist

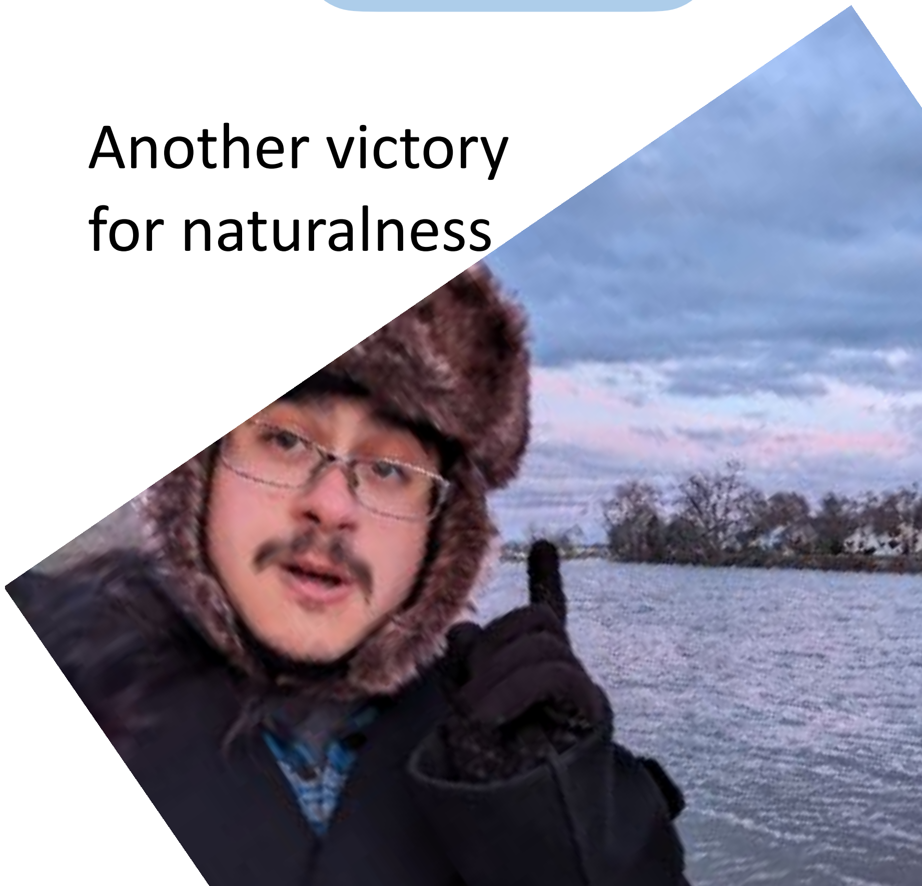
There's a subtler notion of symmetry!

X not fully broken, but converted to a non-invertible symmetry! This must act both on local fields and on 't Hooft lines.

Choi, Lam, Shao
2205.05086
Córdova, Ohmori
2205.06243

$$\psi(x) \to \psi(x)e^{i\alpha} \quad e^{i\oint_\gamma A_m} \to e^{i\oint_\gamma A_m + i\alpha\oint_\gamma A}$$

Another victory for naturalness



Symmetry which acts on local operators and breaks when one-form symmetry is broken! Operators protected by this symmetry must be generated when there are dynamical monopoles!

Dirac (1938) Naturalness → ’t Hooft (1980) Technical Naturalness →

Clay, Kantaro, Seth, Sungwoo (2022-2024): Non-invertible Naturalness

A spurion for a non-invertible symmetry can be *generated by nonperturbative gauge theory effects* in a UV theory.

Infrared symmetry analysis points you to a Dirac natural model! More powerful than learning a UV model need not be destabilized toward the IR

Find a Dirac natural origin for a technically natural parameter:

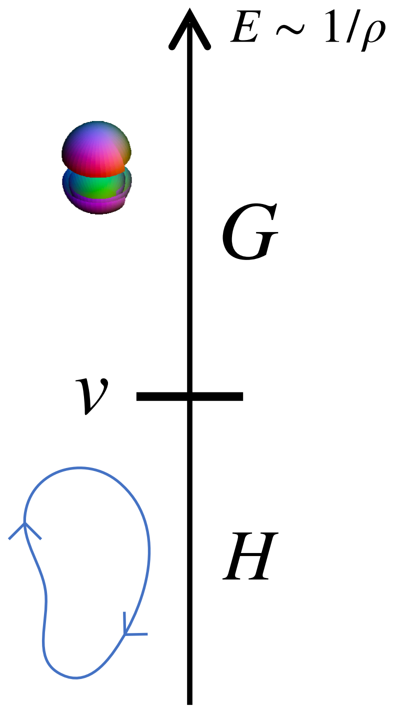
$$y_\nu \sim y_\tau \exp\left(-S_{\text{inst}}\right)$$

Find a Dirac natural origin for an unnatural parameter:

$$y_b \sim y_t \exp\left(-S_{\text{inst}}\right)$$

$$\Rightarrow \bar{\theta} \simeq 0$$

Model-building strategy



A classical global symmetry X protects some operator $\mathscr{O}$ and has an H anomaly

$$\partial_\mu J^\mu_X = \frac{\mathscr{A}}{8\pi^2} H^{\mu\nu}\tilde{H}_{\mu\nu}$$

But some values of $\int_{\mathscr{M}} H\tilde{H}$ not realized for $\mathscr{M} = \mathbb{R}^4$

Non-invertible X symmetry tells us $\mathscr{O}$ could be generated by instantons in the theory $G \supset H$ which has G/H-monopoles

The Standard Model

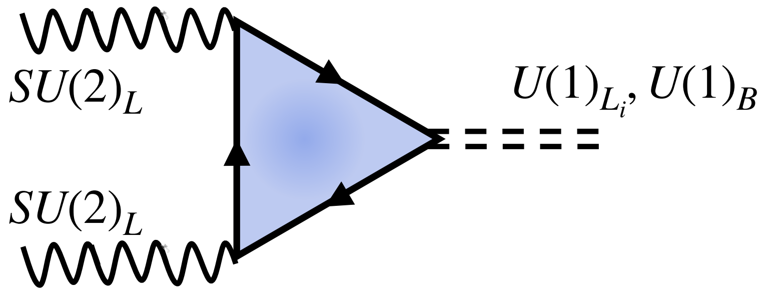
	Q_i	$\overline{u}_i$	$\overline{d}_i$	L_i	$\overline{e}_i$	H
$SU(3)_C$	$\mathbf{3}$	$\overline{\mathbf{3}}$	$\overline{\mathbf{3}}$	–	–	–
$SU(2)_L$	$\mathbf{2}$	–	–	$\mathbf{2}$	–	$\mathbf{2}$
$U(1)_Y$	$+1$	-4	$+2$	-3	$+6$	-3

- Beautiful, yet incomplete.
- Simple way to go beyond is to gauge some of the approximate symmetries of the SM
- E.g. the SU(5) approximate symmetry is broken by $g_1 \neq g_2 \neq g_3$ (and $y_u \neq y_d \neq y_e$)
- Here we're going to go horizontal and play with the $U(3)^5$ approximate symmetries of the SM fermions

SM flavor symmetries actually in 2-group 2212.13193 Córdova & SK

With the Standard Model yukawas

$$\mathscr{L} \supset y^u_{ij}\tilde{H}Q_i\bar{u}_j + y^d_{ij}HQ_i\bar{d}_j + y^e_{ij}HL_i\bar{e}_j$$



Classical Global Symmetry

$$G^c_{\text{SM}} = U(1)_{L_e} \times U(1)_{L_\mu} \times U(1)_{L_\tau} \times U(1)_B$$

Quantum Symmetry after mixed anomaly with $SU(2)_L$

$$G_{\text{SM}} = U(1)_{L_e - L_\mu} \times U(1)_{L_\mu - L_\tau} \times U(1)_{B - N_c L}$$

$G_{\text{SM}} \supset \mathbb{Z}^L_{N_g}$

See my note 2204.01741 on proton stability

Natural to think about gauging some of these symmetries.

(Of course from $\Delta m^2_{ij}, \theta^2_{ij}$ we know these symmetries are broken in the far IR, this just means they must be Higgsed. As with SU(5).)

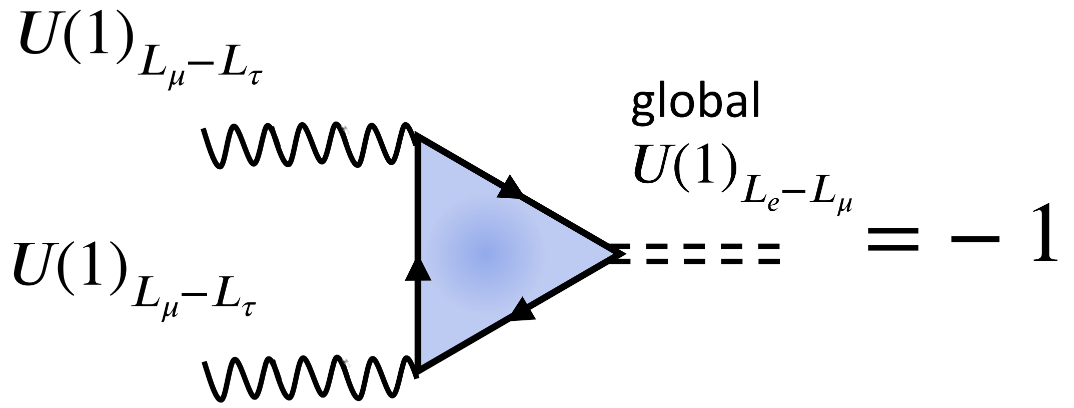
Nonperturbative Quantum Lepton Flavodynamics

Neutrino Masses from Generalized Symmetry Breaking

arXiv:2211.07639, Clay Córdova, Sungwoo Hong, SK, Kantaro Ohmori

Now let's go beyond and gauge $U(1)_{L_\mu - L_\tau}$

There's a new ABJ anomaly diagram to consider



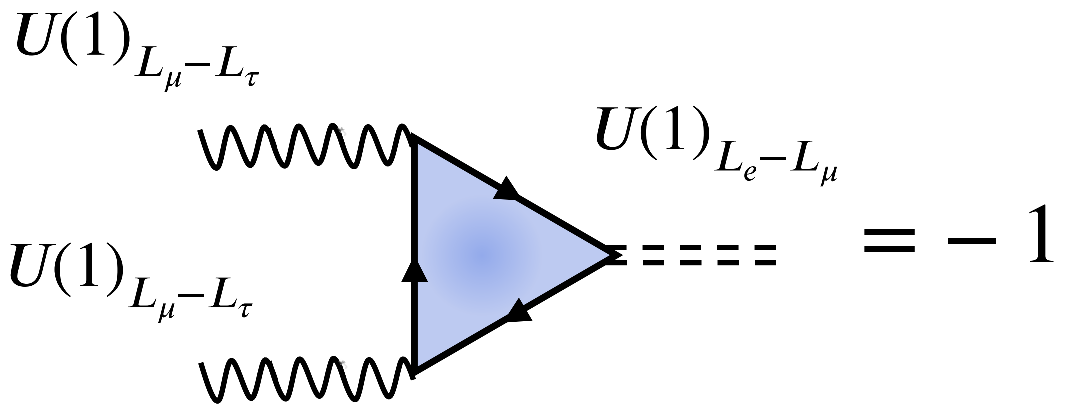
So the global $U(1)_{L_e - L_\mu}$ current is no longer conserved

$$\partial_\mu J^\mu_{L_e - L_\mu} = \frac{-1}{8\pi^2} F_{L_\mu - L_\tau} \tilde{F}_{L_\mu - L_\tau}$$

Non-invertible symmetry!

But the $U(1)_{L_\mu - L_\tau}$ gauge theory cannot saturate this anomaly

Beyond with $Z'_{L_\mu - L_\tau}$



Non-invertible symmetry protects neutrino masses, focus on $\mathbb{Z}_3^L \subset U(1)_{L_e - L_\mu}$

	L_i	$\bar{e}_i$
$\mathbb{Z}_3^L$	$+1$	-1

Disallows $\left(\tilde{H}L\right)^2$

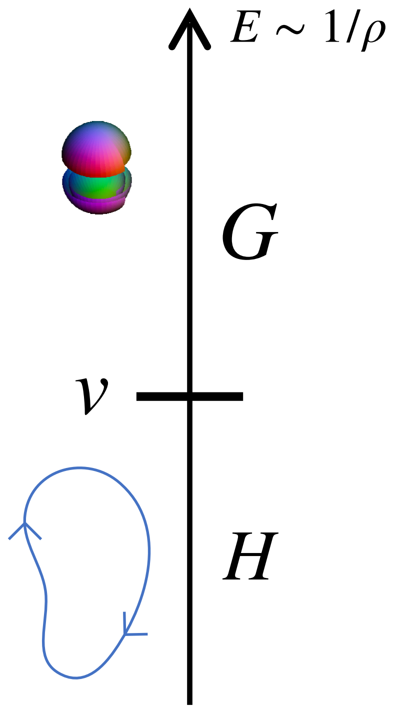
$$L = (L_e - L_\mu) - (L_\mu - L_\tau) \quad (\text{mod } 3)$$

Model-building logic

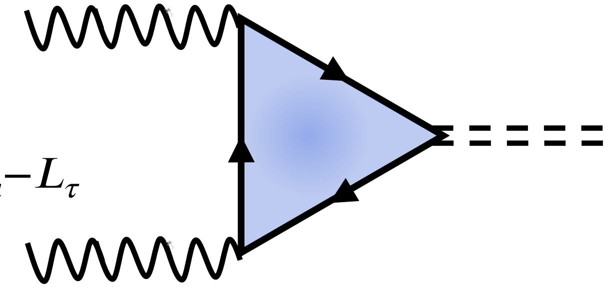
A classical global symmetry $X = \mathbb{Z}_3^L$ protects the operators $\mathcal{O}_{ij} = (\tilde{H}L_i)(\tilde{H}L_j)$ and has an $H = U(1)_{L_\mu - L_\tau}$ anomaly

But while $\int_{\mathcal{M}} H\tilde{H} \in \mathbb{Z}$ generally, $\int_{\mathbb{R}^4} H\tilde{H} = 0$

X is a non-invertible symmetry! In a theory $G \supset H$ with lepton flavor monopoles, $\mathcal{O}_{ij}$ could be classically absent and generated only by G-instantons.



Beyond with $Z'_{L_\mu - L_\tau}$ and N!



Non-invertible symmetry protects neutrino masses
either with or without right-handed neutrinos

	L_i	$\bar{e}_i$
$\mathbb{Z}_3^L$	$+1$	-1

Disallows $\left(\tilde{H}L\right)^2$

	L_i	$\bar{e}_i$	N_i
$\mathbb{Z}_3^{\tilde{L}+N}$	$+1$	-1	$+1$

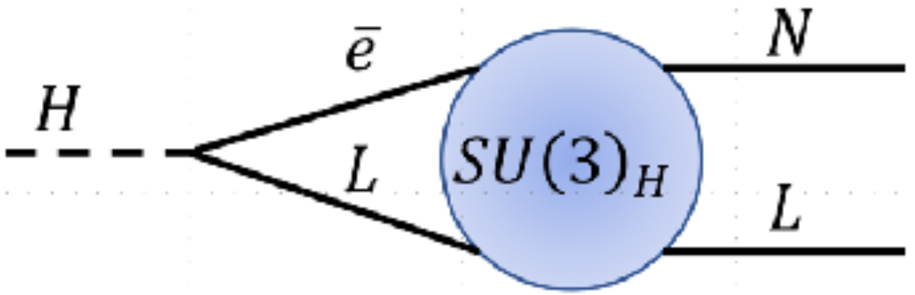
Disallows HLN

Dirac masses:

Write down charged lepton mass $\mathscr{L} \sim y_\tau H \mathbf{L}\bar{\mathbf{e}}$

	$SU(3)_H$	$U(1)_{\mu-\tau}$	$U(1)_L$	$U(1)_N$
$\mathbf{L}$	3	$\begin{pmatrix} L_e \\ L_\mu \\ L_\tau \end{pmatrix} = \begin{pmatrix} 0 \\ +1 \\ -1 \end{pmatrix}$	$+1$	0
$\bar{\mathbf{e}}$	$\bar{3}$	$\begin{pmatrix} \bar{e} \\ \bar{\mu} \\ \bar{\tau} \end{pmatrix} = \begin{pmatrix} 0 \\ -1 \\ +1 \end{pmatrix}$	-1	0
$\mathbf{N}$	$\bar{3}$	$\begin{pmatrix} N_e \\ N_\mu \\ N_\tau \end{pmatrix} = \begin{pmatrix} 0 \\ -1 \\ +1 \end{pmatrix}$	-1	$+1$

Classical $U(1)_N$ symmetry protects the Dirac neutrino mass $\tilde{H}\mathbf{LN}$



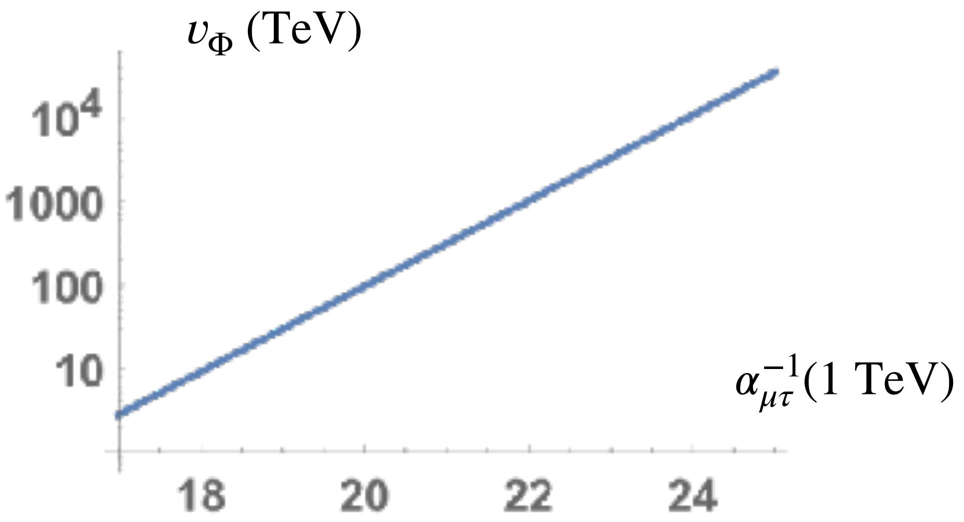
$$\mathscr{L} \sim y_\tau^\star e^{-\frac{8\pi^2}{g_H^2}} \tilde{H}\mathbf{LN}$$

Economical and predictive

Given the discovery of such a Z' for $U(1)_{L_\mu - L_\tau}$, learn the scale

at which $SU(3)_H \to U(1)_{L_\mu - L_\tau}$

$$v_\Phi^2 \sim M_{Z'}^2 \left(\frac{m_\nu}{m_\tau}\right)^{3/2} \exp \frac{3\pi}{4\alpha_{\mu\tau}(M_{Z'}^2)}$$



Texture from Higgses implementing $SU(3)_H \to U(1)_{L_\mu - L_\tau} \to \emptyset$

c.f. [Alonso-Álvarez, Cline '21]

Nonperturbative Quantum Quark Flavodynamics

Non-Invertible Peccei-Quinn Symmetry and the Massless Quark Solution to the Strong CP Problem

arXiv:2402.12453, Clay Córdova, Sungwoo Hong, SK

What about the quark sector?

- If we've found something interesting purely with the leptons, shouldn't there be something interesting to find with the quarks?
- But quarks have this pesky extra $SU(3)_C$ quantum number which means you'll get thrice as many legs in a 't Hooft vertex, and generating a 12-quark interaction is not so interesting.

Require a more subtle notion (and more subtle usage) of non-invertible symmetry

Quark Weak CP and Strong CP Violation

The 'strong CP angle' $\bar{\theta} = \arg e^{-i\theta} \det \left(y_u y_d \right)$ is constrained to $\bar{\theta} \lesssim 10^{-10}$!

Even worse, we also have the 'weak CP angle' $\tilde{J} = \mathrm{Im} \det \left(\left[y_u^\dagger y_u, y_d^\dagger y_d \right] \right)$

oft parameterized by m_i, θ_{ij}, and the phase $\delta_{\rm CKM} \sim 1.14$

A small value of $\bar{\theta}$ is not technically natural $\Rightarrow$ the strong CP problem.

Upon RG evolution, $\delta\bar{\theta} \propto c\delta_{\rm CKM}$

Ellis & Gaillard '79

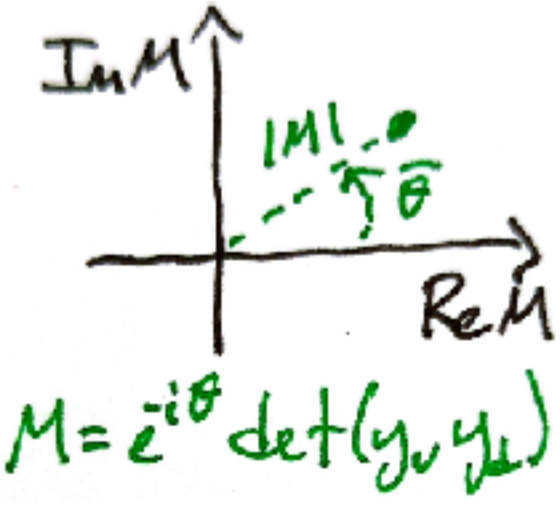
Peccei-Quinn for Strong CP

Now consider a Peccei-Quinn symmetry protecting the up quark mass

$$U(1)_{\mathrm{PQ}}: \quad \bar{u} \to \bar{u} e^{i\alpha} \quad \Rightarrow \quad \tilde{H} Q \bar{u} \text{ charged so } y_u = 0$$

If the PQ symmetry is good, $y_u \to 0$, and so $\det y_u \to 0$ and there's no strong CP violation

Easier to parameterize in 'Cartesian coordinates' for complex parameter $M \in \mathbb{C}$

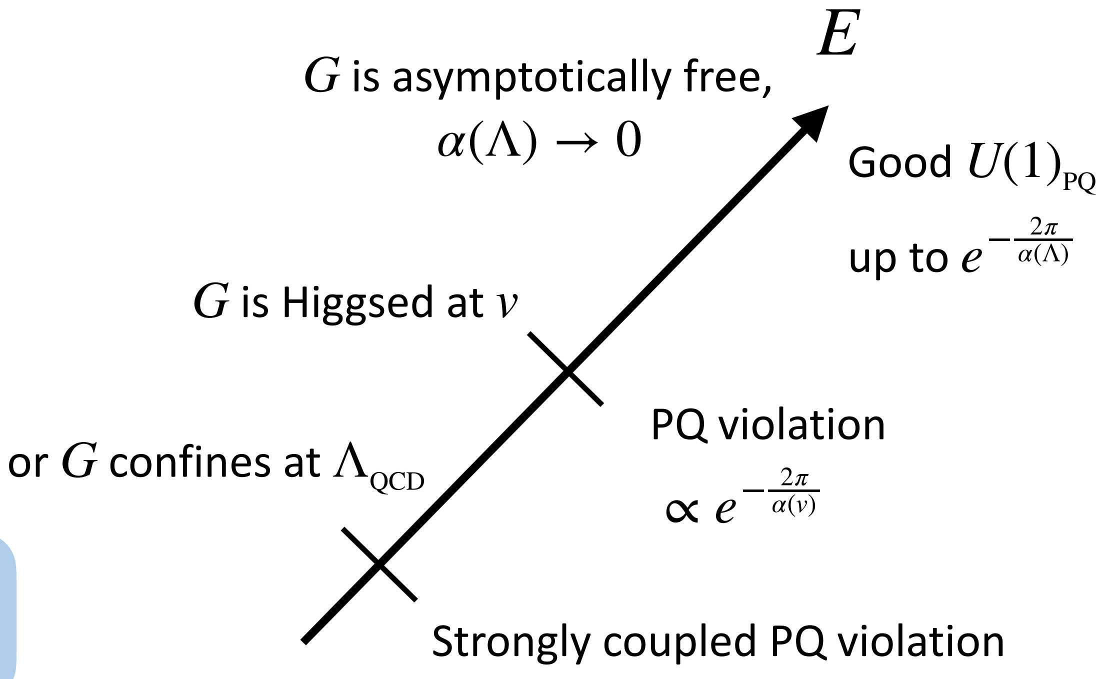


Def $M = e^{-i\theta} \det\left(y_u y_d\right)$, so $\bar{\theta} = \arg M$

Transforms as CP : $\mathrm{Im}(M) \to -\mathrm{Im}(M)$

Peccei-Quinn Violation

Massless up quark?! Not in the IR.
A PQ symmetry which begins good is violated by instantons at low energies



UV $y_u = 0$ is then violated by QCD instantons to generate mass, automatically $M \in \mathbb{R}_+$.

Georgi-McArthur '81
Kaplan-Manohar '86
Choi, Kim, Sze '88

Heroic efforts by lattice physicists tell us the SM does not bear out the massless up quark solution

Flavour Lattice Averaging Group 2019

Could there be any UV model where instantons revive this solution?

Quark Flavor Z'

In lepton sector we had anomaly-free $U(1)_{L_i - L_j}$

Likewise here we can think about gauging e.g. $U(1)_{B_i - B_j}$

Structurally parallel, just broken more by the larger Yukawas

But in fact in the quark sector we can gauge flavored baryon number in a slightly more subtle way because $N_c = 3 = N_g$!

Quark Flavor Z'

One such combination is $U(1)_{B_1+B_2-2B_3}$ which is anomaly-free

Gauge group can be $\left(SU(3)_C \times U(1)_{B_1+B_2-2B_3}\right)/\mathbb{Z}_3$!

At intermediate scales you can realize $\left(SU(3)_C \times SU(3)_H\right)/\mathbb{Z}_3$ along similar lines

	$SU(3)_c$	$SU(3)_H$
$\mathbf{Q}$	3	3
$\bar{\mathbf{u}}$	$\bar{3}$	$\bar{3}$
$\bar{\mathbf{d}}$	$\bar{3}$	$\bar{3}$

With quotient, allowed configurations with fractional magnetic fluxes labeled by 'second Stiefel-Whitney class'

$$w_2(A_C) = w_2(A_H) \in H^2(M, \mathbb{Z}_3)$$

Non-invertible symmetry

When the global structure is non-trivial, there are color and flavor instantons with fractional instanton numbers

E.g. fractional part of color instanton $\mathcal{N}_C = \frac{1}{8\pi^2}\int_M \mathrm{Tr}\left(F_C \wedge F_C\right) = \frac{1}{3}\int_M \omega \wedge \omega \quad \mod 1$

The diagonal quotient locks the fractional parts together $\mathcal{N}_C = \mathcal{N}_H \quad \mod 1$

Classify global zero-form symmetry $X^{(0)}$ with mixed anomaly from

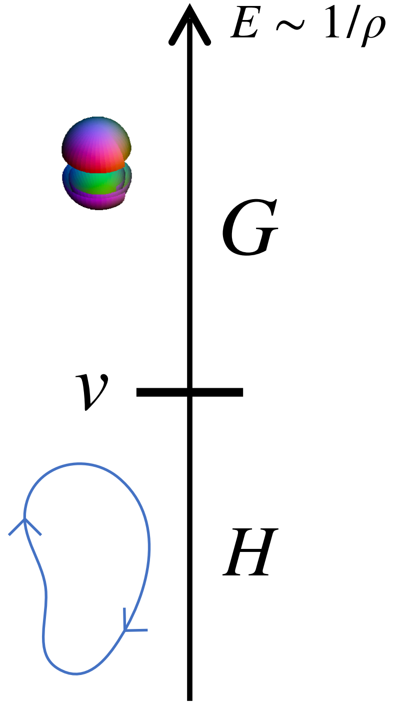
- integer instantons → $X^{(0)}$ broken
- fractional instantons → $X^{(0)}$ noninvertible

Noninvertible $\mathbb{Z}_3^{(0)}$

	Q_i	$\bar{u}_i$	$\bar{d}_i$
$\mathbb{Z}_3^{\tilde{B}}$	$+1$	-1	$+1$

y_d becomes a spurion for this non-invertible symmetry!

Model-building logic



A classical global symmetry $X = \mathbb{Z}_3^{\tilde{B}+d}$ protects the operators $\mathcal{O}_{ij} = HQ_i\bar{d}_j$ and has an $H = \left(SU(3)_C \times SU(3)_H\right)/\mathbb{Z}_3$ anomaly

But while $\int_{\mathcal{M}} H\tilde{H} \in \mathbb{Z}/3$ generally, $\int_{\mathbb{R}^4} H\tilde{H} \in \mathbb{Z}$

X is a non-invertible symmetry! In a theory $G \supset H$ with quark color-flavor monopoles, $\mathcal{O}_{ij}$ could be classically absent and generated only by G-instantons

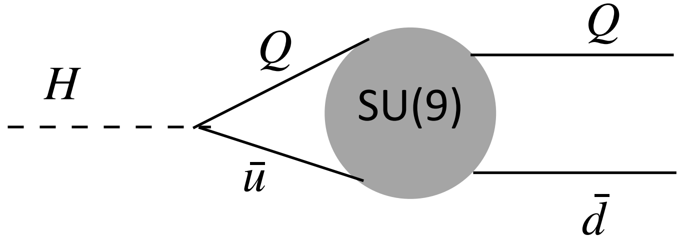
Color-flavor unification!

This all points to a beautiful $SU(9)$ unified theory in which the colors and flavors of the quarks are placed together into the fundamental

	$SU(9)$
$\mathbf{Q}$	9
$\bar{\mathbf{u}}$	$\bar{9}$
$\bar{\mathbf{d}}$	$\bar{9}$

$$\mathscr{L}_0 = y_t \tilde{H} \mathbf{Q}\bar{\mathbf{u}} + \text{h.c.} + \frac{i\theta_9}{32\pi^2} F\tilde{F}$$

Again start with good $U(1)_{\rm PQ}$ and no strong CP violation, then



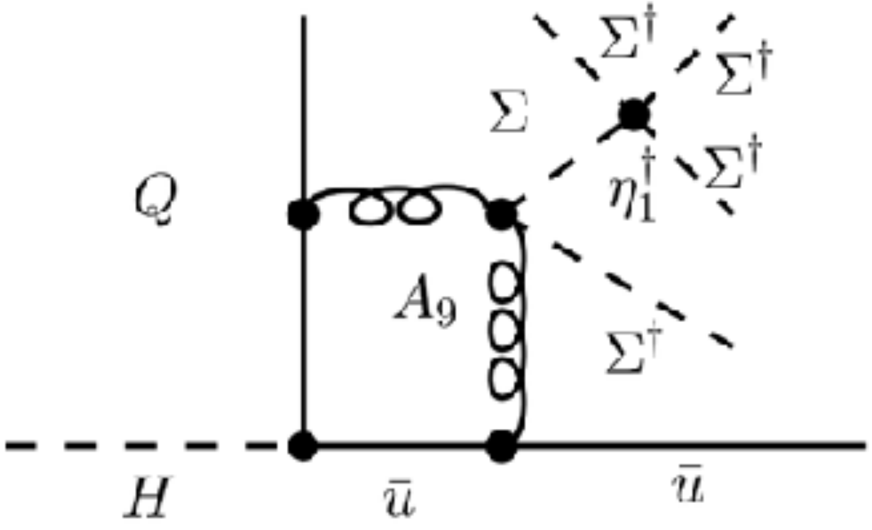
$$\mathscr{L}(\Lambda) \sim y_t H \mathbf{Q}\bar{\mathbf{u}} + y_t^{\star} e^{i\theta_9} e^{-\frac{2\pi}{\alpha_9(\Lambda)}} H \mathbf{Q}\bar{\mathbf{d}} + \text{ h.c.} + \frac{i\theta_9}{32\pi^2} F\tilde{F}$$

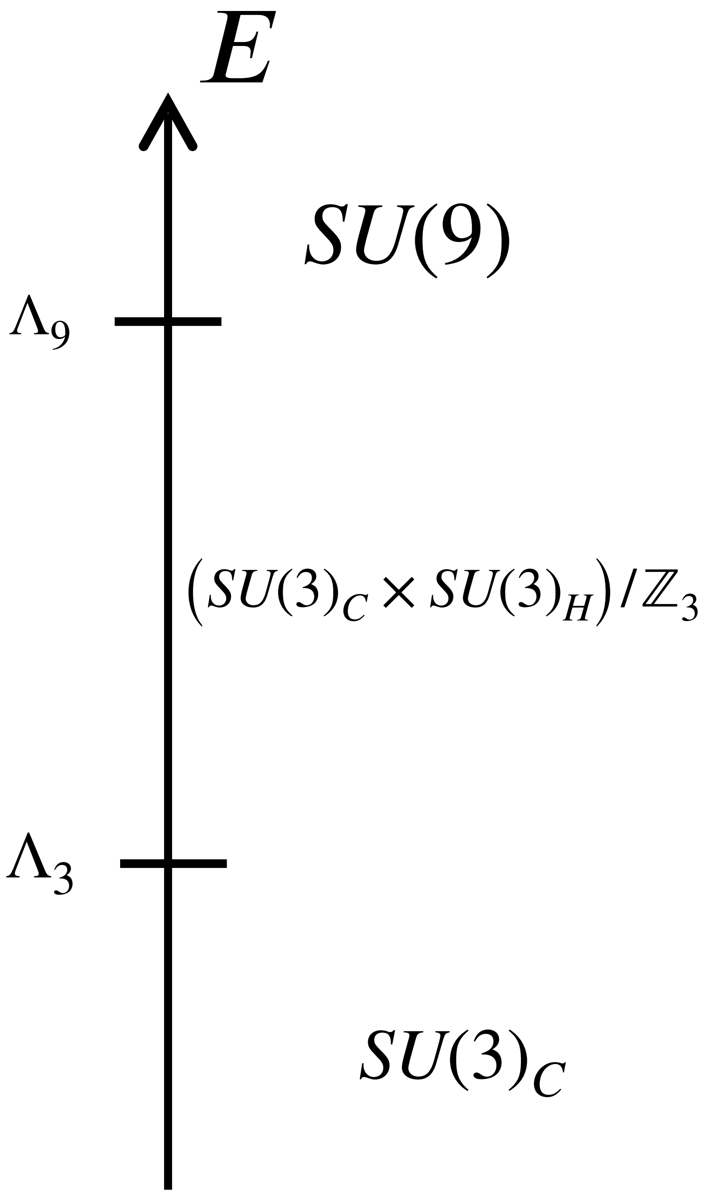
Generating CKM by Bosonic Mediation

General protection

Often strong CP solutions work best only at tree-level and at loop-level protect $\bar{\theta} = 0$ from δ_{CKM} by making use of the small yukawa entries

We instead have a more-general mechanism which would protect $\bar{\theta} = 0$ for general quark yukawas



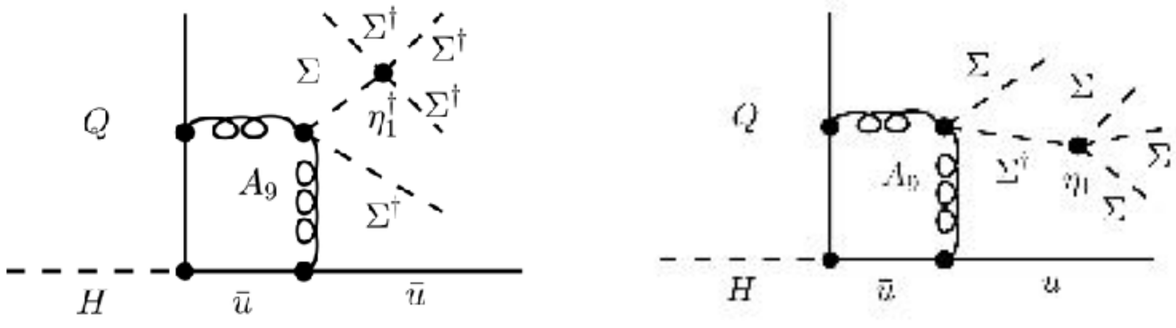


Generating CKM

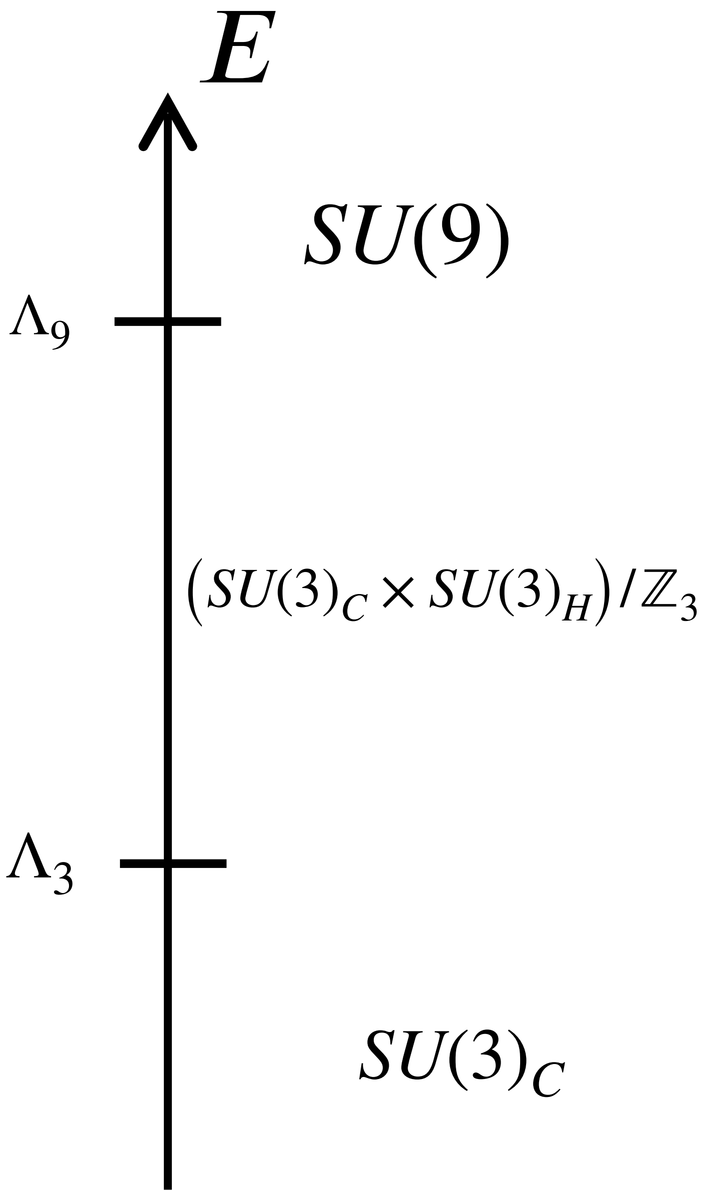
Idea: Communicating flavor-breaking $\langle \Sigma^a{}_b \rangle$ through gauged flavor symmetry lets you generate *hermitian* yukawas

$M = \det(y_u y_d)$ automatically real

$$V_{\mathbb{Z}_4}(\Sigma) = \eta_1 \mathrm{Tr}\left(\Sigma^4\right) + \eta_2 \mathrm{Tr}\left(\Sigma^2\right)^2 + \text{ h.c.}$$



$$(y_u)^a{}_b \sim y_t \left(\mathbb{1}^a{}_b + \frac{\alpha_9}{(4\pi)} \frac{\eta_1^\dagger (\Sigma^{\dagger 4})^a{}_b + \eta_2^\dagger \mathrm{Tr}(\Sigma^{\dagger 2})(\Sigma^{\dagger 2})^a{}_b}{\Lambda_9^4} + \frac{\alpha_9}{(4\pi)} \frac{\eta_1 (\Sigma^4)^a{}_b + \eta_2 \mathrm{Tr}(\Sigma^2)(\Sigma^2)^a{}_b}{\Lambda_9^4} + \dots \right)$$

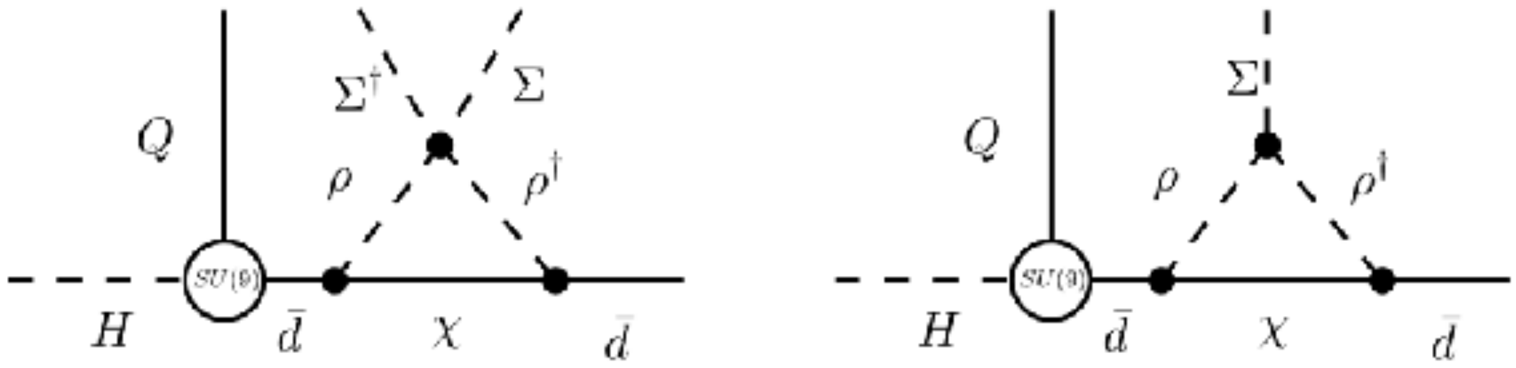


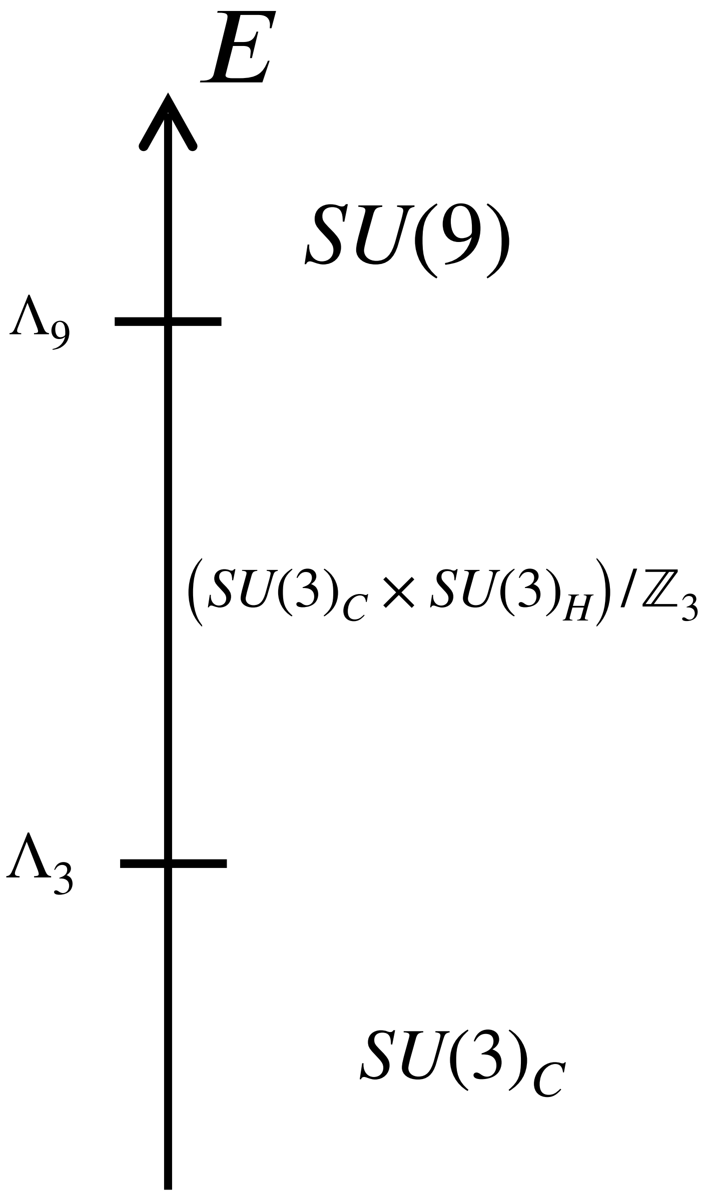
Generating CKM

Yukawas stay hermitian yet $V(\Sigma)$ breaks CP explicitly and/or spontaneously so can generate

$$\delta_{CKM} \propto \arg\det\left(\left[y_u^\dagger y_u, y_d^\dagger y_d\right]\right) \neq 0$$

Another wrinkle: Must treat $\bar{u}, \bar{d}$ differently so they don't commute in flavor space.





Generating CKM

Have shown in principle can generate arbitrary flavor and CP-violation without upsetting Im$M = 0$

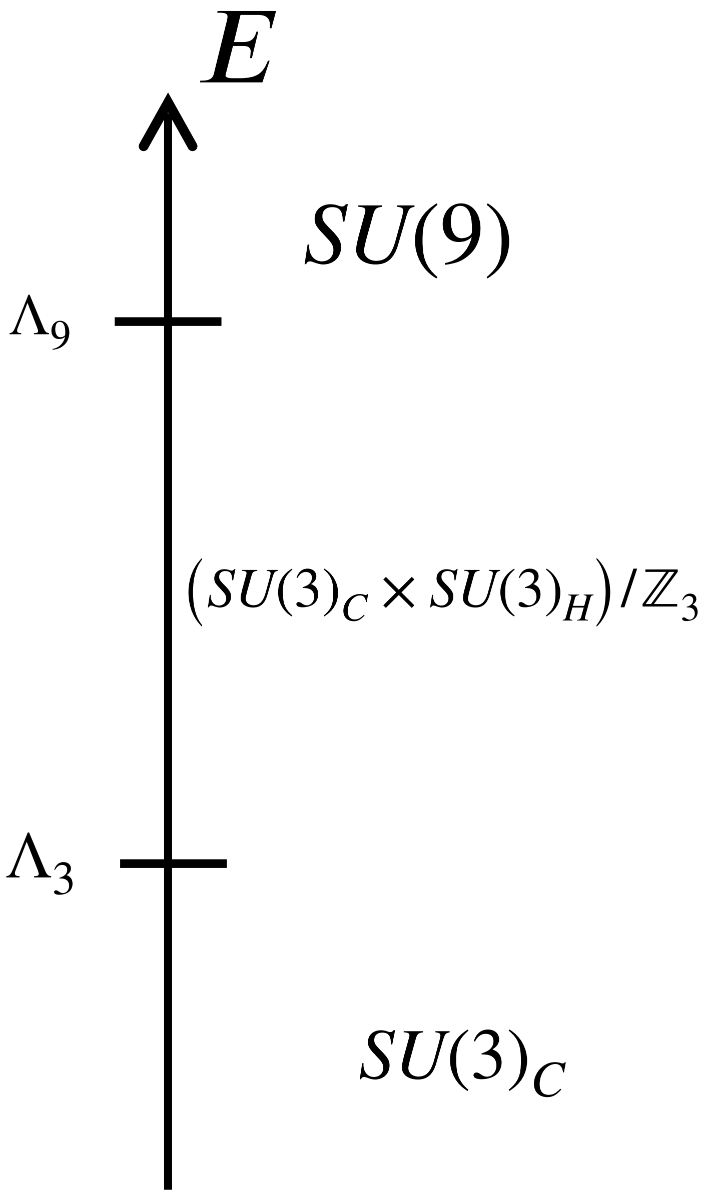
But have not evinced a beautiful way to land specifically on Standard Model flavor.

In particular with H we have $\mathscr{L} \supset y_t \tilde{H} Q \bar{u}$ so get

$y_u = y_t 1^a{}_b + \ldots$ breaking effects

and the quarks don't really want this structure!

Probably need to upgrade $H \to H^a{}_b$ to get flavor correct.

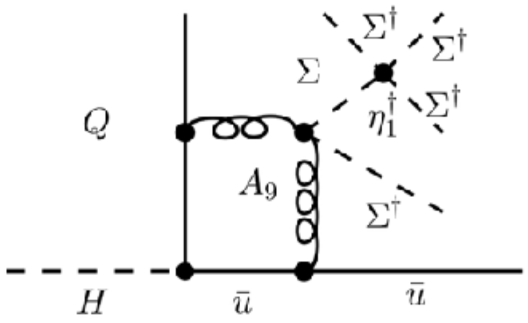


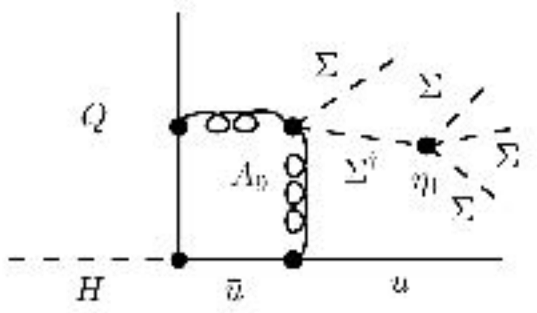
CKM Brief Version

Idea: Communicating flavor-breaking $\langle \Sigma^a{}_b \rangle$ through gauged flavor symmetry lets you generate *hermitian* yukawas

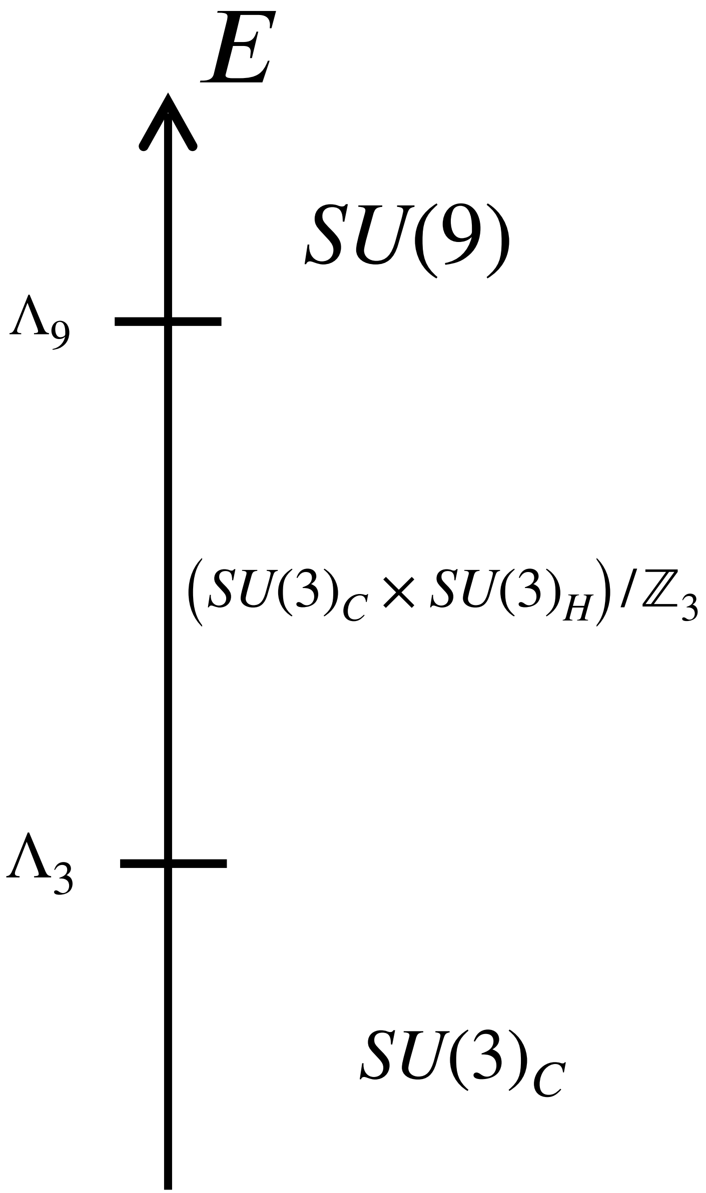
$M = \det(y_u y_d)$ automatically real

No matter what size yukawas or how much CP-violation





$$(y_u)^a{}_b \sim y_t \left(\mathbb{1}^a{}_b + \eta_1^\dagger (\Sigma^{\dagger 4})^a{}_b + \eta_1 (\Sigma^4)^a{}_b + \dots \right)$$



Seth's conclusions

At least any place nonperturbative effects might be phenomenologically relevant, I expect paradigm of generalized global symmetries will offer better understanding.

> Already we have located new unified theories of the SM fermions with instanton effects which can solve SM naturalness issues! Both technically natural, and not.

As particle physicists we are not yet done learning about the role of symmetries!

A primate pleased they newly uncovered some simple, reductionist BSM models