The Physical Weak Current and Quark Mixing in VERSF

Coherent propagation and the complete recorded weak response

Keith Taylor
VERSF Theoretical Physics Programme
22 November 2026

General reader summary

Quarks come in different types, and the weak force can turn one type into another. Experiments measure the pattern of these changes. VERSF has a promising conditional account of part of that pattern, but a physical explanation must also show what happens while the quarks move and their interactions leave records in their surroundings. This paper carries out that combined calculation in a specified finite model.

It finds that motion and recording together can change the mixing pattern even when recording alone would leave its underlying pattern unchanged. The size of this change depends on which combinations of quark types the surroundings record. It also counts the fluctuations around the average effect. These matter: looking only at the average is like judging a fluctuating signal from its average height, without counting the variation that contributes to what a detector receives. The paper calculates both parts and gives the resulting transition probabilities.

A further calculation shows how one common change in the quarks' motion changes their masses and mixing together. This moves the programme from a compatibility check to an explicit dynamical example that can be tested. It does not yet derive the measured pattern completely: the physical rate of recording, which combinations the surroundings record, how strongly they couple to the quarks, and the experimentally accessible fluctuations still have to be supplied. The calculation keeps those choices visible and does not choose them by fitting the answer.

Abstract

We construct and solve a finite quark evolution combining the published native recording instrument with propagation in the inherited mass frames. Recording alone has an exact positive weak vertex B with $I/9 \le B \le I$ at maximal activity and preserves the polar mixing factor at fixed canonical frames. Interleaved propagation evades that restriction through a derived second-order anti-Hermitian correction, $-i\omega\delta_r\tau^2 \mathrm{Off}_R(H_d - H_u)/2$, with $\delta_r = 0.0463277936554$. The recording basis R is a declared input; an up-mass-frame choice strongly suppresses the displayed propagation correction. The complete 36 × 36 generator returns the current, its noise and all flavour-resolved probe probabilities; a connected family graph gives a unique maximally mixed long-time state. For a specified channel followed by a charged probe, we derive closed time laws and distinguish the inclusive total 0.940828969 from the squared mean 0.777320719 at unit exposure. Their difference is fixed by the channel, not an additional coefficient. A separate declared gauge-covariant kinetic action returns correlated mass and CKM changes with an exactly normalised weak current. We reproduce the bare and conditional AX1 branches, retain their provenance, and expose sensitivity to the kinetic coefficient and propagation-to-event-rate ratio. These are constructive dynamical and response results. Physical source admission, the event rate and the energy-resolved bath attachment remain required for a measured CKM prediction.

Contents

1 Question and source inheritance 3
2 One kinetic form and one weak current 3
3 The native recording operation 4
4 The positive weak vertex theorem 5
5 Repeated events and their time dependence 6
6 Executed matrix checks 7
7 The coherent quark construction 8
8 The bath metric and the axle amplitude 9
9 A common kinetic action with changing mass frames 10
10 The complete recorded weak response 12
11 Propagation between recording events 15
12 Consequences for the November ledger 19
13 Conclusion 20
Appendix A Product state coefficient 20
Appendix B Reproduction and source coverage 20
References 21

Notation

Symbols used throughout the paper

Symbol	Meaning
$\mathcal{F}$, $\mathbb{C}$, $\mathcal{W}$	Family, colour and weak spaces, of dimensions 3, 3 and 2
Z_Q, U_u, U_d	Common left-handed kinetic form and canonical left mass frames
R, ρ, z, χ	Recording basis, reference ray, its unit vector and native activity
C_α, $\mathcal{E}^*$	Hermitian recording operators and their Heisenberg map
B, Φ	Weak vertex family matrix and weak-adjoint family map
α, β	Product-state coherence and diagonal response coefficients
h, γ, τ	Event probability, event rate and exposure τ = γt
w, z_E, κ_E	Charged E_2 response, E_1 source and E_2 maintenance coefficient
H, M_0, λ_J	Factor overlap matrix, current metric and common intensity coefficient
G, η	Hermitian kinetic generator and its declared embedding coefficient
Ψ, Φ, b_W, r	Singlet and adjoint family maps; weak and family attenuation
P, W, A(τ)	Probe weights, conditional probabilities and total lowering weight
p, ω, H_q, L	Momentum, propagation/rate ratio, Hamiltonian and evolution generator

Status convention. “Exact” refers to a proved identity under its stated hypotheses. “Executed” refers to the finite calculation reproduced here. “Inherited” denotes a result taken from a named public paper. A well-posed construction is not, by that description alone, the physically admitted VERSF realisation. No uniqueness of realisation is required.

1 Question and source inheritance

The November ledger asks for masses, kinetics and charged-current mixing on a common compatible construction [1, SM-13, SM-17 and SM-22]. The relevant question is therefore not whether a new matrix resembles the measured Cabibbo–Kobayashi–Maskawa matrix. It is whether the same construction returns the quark mass frames and the operator through which the weak interaction sees them. This paper calculates that operator and then solves its interleaving with quark propagation. The constructive result is Proposition 6: the combination generates a coherent correction that recording alone cannot generate. Section 10 returns complete transition probabilities, and Proposition 7 determines the long-time limit of the joint model.

The bath criterion and the recording instrument perform different jobs. A shared conserved bath permits coherent transfer of weight between alternatives under the hypotheses of the bath papers [2,3]. It does not prescribe numerical CKM angles. The native instrument describes the change of state when an event is recorded [9]. Coherent evolution before recording may change the Yukawa tensors; a completed recording operation may instead change the weighting of an already specified weak current. Both effects must be retained when comparing with experiment. The neutral-sector analysis likewise separates a mass construction from its weak-current attachment [10].

The quark series already contains a target-blind regulator return, RRHF, and a conditional curvature extension, RRHF-AX1 [4,5]. The first has deficient third-family mixing and CP violation; the second gives a much better conditional pattern. The full-carrier and primitive-matter papers also supply more than a generic permission for a correction: they give the quadratic axle, a nonzero cubic source jet, and a retained R4-to-charged carrier [5,6]. The later current and time papers provide magnitude-bearing geometry and nontrivial bounds [7,8]. The earlier quark ledger makes phase provenance explicit [13], and the assignment–bath analysis treats the mass hierarchy as a separate response problem [14].

Table 1. Source results retained in this calculation

Source	Result used	Scope
[2,3] Bath papers	Shared weight permits coherent mixing	Does not determine CKM entries
[4,5] Quark papers	RRHF and conditional AX1 matrices; quadratic axle	Source admission remains qualified
[6] Primitive matter	Cubic coefficient 1/2; E_2 maintenance form; R4 retention	Factor coordinate and completion attachment still matter
[7,8] Current and clock papers	Factor matrix H; current-cycle metric; portal bounds	Absolute physical descent is not supplied by H alone
[9] Completed event paper	Full resolved native state operation	This paper applies it to the weak current

The result below is not a replacement for the force programme. In particular, the inherited one-calibration force branch and its 15/14 premise are not removed or recalibrated here [12]. The present task concerns the family structure of the weak vertex. Numerical attenuation of that vertex is not identified with a change in the gauge kinetic coefficient.

2 One kinetic form and one weak current

Work first in a fixed weak-doublet source frame. Let Z_Q be a positive 3×3 family kinetic form common to the up and down components. Gauge covariance of the corresponding covariant derivative gives the

charged-current family factor Z_Q. With right-handed kinetic forms equal to the identity for this finite construction, canonical normalisation gives

$$Q_L^{can} = Z_Q^{1/2} Q_L, \quad J_{can} = Z_Q^{-1/2} Z_Q Z_Q^{-1/2} = I. \tag{1}$$

The right-hand equality is the required current check. The identity removes an explicit kinetic factor from the current. It does not keep the mass frames fixed when the source Yukawa tensors are held fixed; Section 9 calculates that distinct response. If the unnormalised current is a different operator, that operator must be returned and canonicalised; its value cannot be inferred from the mass eigenvalues.

Proposition 1. For any positive Z_Q, positive diagonal singular-value matrices D_u, D_d, and unitary U_u, U_d, the finite construction

$$Y_u^{source} = Z_Q^{1/2} U_u D_u, \quad Y_d^{source} = Z_Q^{1/2} U_d D_d \tag{2}$$

has canonical Yukawa singular values D_u, D_d and weak-current mixing

$$V_0 = U_u^\dagger U_d. \tag{3}$$

Proof. Left multiplication by $Z_Q^{-1/2}$ returns $U_f D_f$. Their left singular vectors are U_f; Eq. (1) then gives Eq. (3). ▫

This constructs a common mathematical setting for a kinetic form, masses and a weak current. It is not a proof that an arbitrary Z_Q or U_f is physically supplied by VERSF. The numerical check uses the inherited charged singular values and a fixed non-diagonal positive Z_Q; canonical-current and singular-value residuals are below 8×10^{-16}. A simultaneous change of source basis transforms both mass tensors and current and leaves the observable vertex unchanged.

All following recording operators are defined in this canonical Hilbert-space norm. Reusing the source kinetic factor as an additional recording weight would count the same normalisation twice.

3 The native recording operation

The resolved quark-doublet space is $\mathcal{F} \otimes \mathbb{C} \otimes \mathcal{W}$. Colour and weak fundamental generators have $\mathrm{Tr}(t_a t_b) = \delta_{ab}/2$. Hypercharge is scalar on this resolved sector, so its commutator below vanishes. For a fixed unit reference ray z, define

$$\rho = |z\rangle\langle z|, \quad v_a = (T_a - \langle z|T_a|z\rangle)z, \quad s = \Sigma_a \|v_a\|^2. \tag{4}$$

$$K_a = i[T_a, \rho], \quad C_a = \sqrt{(\chi/s)}K_a, \quad C_0 = \sqrt{(I - (\chi/s)\Sigma_a K_a^2)}. \tag{5}$$

$$\mathcal{E}^*(X) = \Sigma_\alpha C_\alpha X C_\alpha, \quad 0 \le \chi \le 1, \quad \Sigma_\alpha C_\alpha^2 = I. \tag{6}$$

These are the native operators used in the completed-event paper [9, Section 12.5]. Each C_α is Hermitian. For each fixed reference ray, $\mathcal{E}$ is a completely positive, trace-preserving and unital linear map. Dependence of that map on its reference ray does not license treating a state-dependent iteration as one linear channel. Event occurrence and the conditional activity χ are separate: $\chi = 1$ is the inherited maximal-positivity convention, not a derived probability of birth.

Let $T_a^W = I_F \otimes I_C \otimes \sigma_a/2$. We extract the colour-singlet weak-adjoint response by averaging the channel over SU(3) × SU(2) gauge orientations with normalised Haar measure. Denote the averaged map by $\mathcal{E}^*$. Representation covariance then defines a family map Φ through

$$\mathcal{E}^*(F \otimes I_C \otimes t_a) = \Phi(F) \otimes I_C \otimes t_a, \quad B = \Phi(I_F). \tag{7}$$

The averaging is an explicit restriction. It is the isotropic gauge component appropriate to a background with no retained colour or weak orientation. A fixed oriented background can generate additional tensor

structures that this projection removes. No claim about those structures is inferred from B. The projection is nevertheless exact, basis covariant, and executable without a choice of a weak axis:

$$\Phi(F) = (2/9)\Sigma_{a=1}{}^{3}\mathrm{Tr}_{CW}[T_a{}^{W}\,\mathcal{E}^{*}(F\otimes I_C\otimes t_a)]. \tag{8}$$

The factor 2/9 follows from $\mathrm{Tr}_{CW}[(I_C\otimes t_a)^2] = 3/2$ and the sum over three weak generators. Thus neither a charge Gram nor an empirical mixing value is inserted as an extra normalisation.

4 The positive weak vertex theorem

Theorem 2. For every pure reference ray on the quark-doublet carrier and every native activity $0 \le \chi \le 1$, the family coefficient in Eq. (7) obeys

$$[1 - (4/9)(1 - \sqrt{(1-\chi)} + \chi)]I \le B \le I. \tag{9}$$

In particular $I/9 \le B \le I$ at $\chi = 1$. The same bounds hold for a fixed convex average of reference-ray channels. For fixed canonical mass frames, the dressed weak vertex and its right polar factor therefore satisfy

$$V_{rec} = U_u{}^{\dagger}BU_d, \quad \mathrm{polar}(V_{rec}) = U_u{}^{\dagger}U_d = V_0. \tag{10}$$

Proof. Write $A = \Sigma_a|v_a\rangle\langle v_a|$. Orthogonality $v_a \perp z$ gives

$$\Sigma_a K_a{}^2 = s\rho + A, \quad A\rho = 0, \quad A \ge 0, \quad \mathrm{Tr}\,A = s. \tag{11}$$

The fundamental SU(2) completeness identity is $\Sigma_a t_a X t_a = (\mathrm{Tr}\,X)I/2 - X/4$, with any family and colour indices left intact. Applied to Eq. (8), it gives the exact matrix identity

$$B = (1/9)\Sigma_\alpha \mathrm{Tr}_C[(\mathrm{Tr}_W C_\alpha)^2] - I/3. \tag{12}$$

All partial weak traces are Hermitian, so every squared term contributes a positive family operator after the colour trace. Put $D = I - C_0$. Its eigenvalue on z is $1 - \sqrt{(1-\chi)}$. On the perpendicular support, if a_k are the eigenvalues of A, then $0 \le a_k/s \le 1$ and $\Sigma_k a_k/s = 1$. Since $1 - \sqrt{(1-x)} \le x$,

$$0 \le D \le I, \quad \mathrm{Tr}\,D \le 1 - \sqrt{(1-\chi)} + \chi. \tag{13}$$

Retain only the $\alpha = 0$ term in Eq. (12), expand $\mathrm{Tr}_W C_0 = 2I - \mathrm{Tr}_W D$, and discard the remaining positive square. Since $\mathrm{Tr}_{CW} D \le (\mathrm{Tr}\,D)I$, this yields the lower bound in Eq. (9). The upper bound follows independently from

$$I - B = (1/9)\Sigma_{\alpha,a}\mathrm{Tr}_{CW}([C_\alpha, T_a{}^{W}]^{\dagger}[C_\alpha, T_a{}^{W}]) \ge 0. \tag{14}$$

The smallest lower bound is 1/9, so B is strictly positive. Finally $V_{rec}{}^{\dagger}V_{rec} = U_d{}^{\dagger}B^2U_d$, whose positive square root is $U_d{}^{\dagger}BU_d$. Multiplication by its inverse proves Eq. (10). Convex averaging preserves both operator inequalities. ▫

The physical distinction is exact. A positive matrix can alter individual transition amplitudes, including off-diagonal entries in a mass frame. It can also produce non-unitarity and apparent changes in quantities extracted under a unitary assumption. These effects do not change the polar unitary factor. Calling such a weighting a new CKM prediction would discard part of the computed weak vertex.

This is not a general statement that an environment cannot change flavour. It fixes the Yukawa frames, uses the native Hermitian instrument, and projects onto the isotropic weak-adjoint channel. Coherent feedback on the mass tensors, a different recording operation, additional gauge tensors or a changed kinetic form require a new joint calculation. The polar factor is a useful decomposition of this vertex; experimental fits must still use the full vertex, not silently replace it by that factor.

5 Repeated events and their time dependence

For a product reference $z = f \otimes c \otimes w$, let $P = |f\rangle\langle f|$ and $Q = I - P$. Fundamental-generator completeness gives $s = 3/2$. Define $r = \sqrt{(1 - \chi)}$ and $u = \sqrt{(1 - \chi/3)}$. The entire weak-adjoint family map, including its action on off-diagonal family operators, is

$$\Phi(F) = QFQ + \alpha(PFQ + QFP) + \beta PFP, \tag{15}$$

$$\alpha = (r + 3u + 2)/6, \quad \beta = [ru + 2u + 3/2 - \chi/3]/(9/2). \tag{16}$$

At maximal activity these reduce to

$$\alpha = (2 + \sqrt{6})/6 = 0.741581624, \quad \beta = (7 + 4\sqrt{6})/27 = 0.622146629. \tag{17}$$

The coefficient β is the isotropic weak-adjoint operator response. It differs from the factor $1/3$ multiplying the weak mean of a particular quark-doublet product ray in [9, Table 18]. The latter is a state expectation in an oriented preparation; averaging the full operator channel answers a different question.

Theorem 3. The map in Eq. (15) is completely positive for all $0 \le \chi \le 1$. Every finite sequence of these maps, with arbitrary product family rays and activities, maps I to a strictly positive matrix. Consequently the fixed-mass-frame polar identity in Eq. (10) survives that whole sequence.

Proof. With $A_f = Q + \alpha P$, the map is

$$\Phi(F) = A_f F A_f + (\beta - \alpha^2)PFP. \tag{18}$$

It is enough to show $\beta \ge \alpha^2$. Substituting $\chi = 1 - r^2$, $0 \le r \le 1$, reduces the inequality to $\sqrt{3}(r + 2)\sqrt{(r^2 + 2)} \ge 2r^2 + 6r + 1$. Both sides are nonnegative. The difference of their squares is $(1 - r^2)(r^2 + 12r + 23)$, which is nonnegative. Also $\alpha > 0$, so A_f is invertible and maps every strictly positive F to a strictly positive contribution. Composition proves the assertion. ▫

This proof uses the family superoperator, not an assumed product B_2B_1. Multiplying the one-event positive matrices instead would generally introduce a spurious polar rotation because the matrices need not commute. It is the action on the operator between events that has to be composed.

Table 2. Product-state response at fixed native activity

χ	α	β	Equal-family factor $(2 + \beta)/3$
0.0	1.000000000	1.000000000	1.000000000
0.1	0.983043256	0.970175274	0.990058425
0.5	0.907619928	0.845460537	0.948486846
1.0	0.741581624	0.622146629	0.874048876

For a fixed source ensemble and an independent Poisson event clock of rate γ, the weak-family observable evolves according to

$$dF/dt = \gamma[\Phi(F) - F], \quad F(0) = I, \quad F(t) = \exp[\gamma t(\Phi - \mathrm{Id})](I). \tag{19}$$

The exponential is the positive Poisson sum of completely positive powers of Φ. For a fixed family ray, $F(t) = Q + \exp[-\gamma t(1 - \beta)]P$. For an equal mixture of three orthogonal family rays,

$$F(t) = \exp[-\tau(1 - \beta)/3]I, \quad \tau = \gamma t. \tag{20}$$

At maximal activity and $\tau = 1$, the common amplitude factor is 0.881657938 and its square is 0.777320719. These are the response of the declared event model, not a physical rate prediction. No measured weak strength or quark angle fixes τ. On the homogeneous clock attachment of [8], one may

write $\gamma = (c_*/a_t)[-\log(1 - h)]$ when the event law supplies a per-opportunity probability h. The physical current background and h still have to be supplied for an absolute time prediction. Time conversion cannot create a polar mixing correction that vanishes for every τ.

The all-sequence result is proved for product family–colour–weak reference rays, including fixed convex ensembles of such channels. The one-event theorem allows entangled rays, but its repeated-event extension is not assumed: the reduced weak-adjoint family map need not then be completely positive, even though the full matter channel is.

6 Executed matrix checks

The implementation builds all eight colour and three weak fundamental generators on the 18-state carrier, includes the scalar hypercharge generator, and forms the recording operators directly. Two independent expressions compute B: the generator projection in Eq. (8) and the Fierz formula in Eq. (12). The Haar average is evaluated analytically by the colour partial trace and weak-adjoint contraction in Eq. (8), checked independently with Eq. (12); no sampled gauge-angle integral is used. The 128 random draws below are reference rays, not gauge orientations. These tests check the implementation; the operator bounds themselves follow from the proof.

Table 3. Numerical verification

Check	Result
128 complex reference rays; fixed seed 20261121	No CKM target enters source generation
Largest completeness residual $\|\|\Sigma C^2 - I\|\|_F$	1.25e-14
Largest projection versus Fierz residual	1.62e-15
Smallest sampled eigenvalue of B	0.628110084
Largest polar identity residual	7.12e-15
Product-state family-map formula residual	5.56e-17

A fixed entangled source, generated before comparison with any quark target, gives

$$B = \begin{bmatrix} 0.912813 & 0.035091 - 0.041432i & -0.016946 + 0.017480i \\ 0.035091 + 0.041432i & 0.759427 & -0.000711 - 0.020170i \\ -0.016946 - 0.017480i & -0.000711 + 0.020170i & 0.803825 \end{bmatrix}. \quad (21)$$

Its eigenvalues are 0.736039646, 0.805247360 and 0.934778112. A simultaneous colour and weak basis rotation changes B by less than 2×10^{-15}. On the bare quark frames of Section 7, the raw vertex has $\|V^\dagger V - I\|_F = 0.591201421$, while its polar factor returns the original unitary to 5×10^{-16}.

The raw quartet $|\mathrm{Im}(V_{11}V_{22}V_{12}^*V_{21}^*)|$ is 0.0104025 in that non-unitary example. It is not the CKM Jarlskog invariant. The polar unitary has $|J| = 1.84350627 \times 10^{-5}$, exactly the inherited bare value. This gives a concrete check against interpreting a large record-induced change in a matrix entry or quartet as an explanation of quark CP violation.

The reduced weak-family map for this source has a Choi eigenvalue −0.0243804. Projection onto an adjoint operator sector need not preserve complete positivity, although the full matter channel has that property. Thus Theorem 3 cannot simply be extended to entangled sequences.

For a separate repeated-event test, use product family rays (1,0,0), (1,1,i)/√3 and (1,i,1)/√3 in that fixed cyclic order, each with χ = 1. They are not selected from mass matrices. Table 4 follows the actual map composition.

Table 4. Changing product-family references

Events	Eigenvalues of F	Polar residual
1	0.62214663, 1.00000000, 1.00000000	5.75e-16
3	0.40144696, 0.71834263, 0.85023209	4.07e-16
10	0.06697058, 0.36000131, 0.53944164	7.20e-16
30	0.00757184, 0.03228674, 0.16030128	1.34e-15

7 The coherent quark construction

The conditional coherent axle is retained as a separate mechanism. For the explicitly reconstructible bare branch, define the anti-Hermitian generator

$$\Omega = [0, a, c\exp(i\varphi); -a, 0, b; -c\exp(-i\varphi), -b, 0], \tag{22}$$

$$a = 9/40, \quad b = 81/2000, \quad c = 243/100000, \quad \varphi = 2\pi/3. \tag{23}$$

These are inherited inputs, not parameters estimated in this paper. In particular the previously calibrated quark information is not relabelled as a new prediction. With $X = \zeta E_{23} - \zeta^* E_{32}$, the common role-even extension is

$$U_u = \exp(X - \Omega/2), \quad U_d = \exp(X + \Omega/2), \tag{24}$$

$$V_0(X) = \exp(-X + \Omega/2)\exp(X + \Omega/2). \tag{25}$$

Putting the same X inside both exponentials is not a simultaneous post hoc basis rotation. It changes their relative product because X need not commute with Ω. A common left multiplication of U_u and U_d, by contrast, cancels from Eq. (3). This is the coherent distinction needed to interpret the quadratic axle.

The inherited illustrative AX1 branch uses $|\zeta| = b/\sqrt{3}$ and $\arg\zeta = 5\pi/6$. Its source admission is not established here by choosing that branch: the full-carrier paper keeps the same-carrier readout, phase and completion obligations explicit [5], while PMS supplies stronger retention and nonlinear-source results [6]. We reproduce the conditional branch as published and report it alongside the baseline.

Table 5. Quark returns with their provenance

Construction	$\lvert V_{ub}\rvert$	$\lvert V_{cb}\rvert$	$\lvert V_{td}\rvert$	$\lvert J\rvert \times 10^5$
Bare Ω, reproduced	0.003931186	0.040284873	0.006108170	1.843506
Bare Ω plus AX1, reproduced	0.003556877	0.040537720	0.008689898	2.996768
RRHF level 4, inherited [4]	0.004198	0.040639	0.005895	1.860329
RRHF-AX1 level 4, inherited [4]	0.003662158	0.040891323	0.008479608	3.023692

The full complex regulator-dependent RRHF generators are not reconstructed from rounded moduli. The two full-carrier rows are quoted results with their public provenance, not new re-executions. The bare and regulated branches are therefore kept separate.

For context, the PDG 2025 update quotes direct averages $|V_{cb}| = 0.0411 \pm 0.0012$ and $|V_{ub}| = 0.00382 \pm 0.00020$, a mixing-based $|V_{td}| = 0.0086 \pm 0.0002$, and the global-fit invariant $J = (3.12^{+0.13}_{-0.12}) \times 10^{-5}$ [11]. These different determinations are benchmarks, not independent inputs to a combined fit here. They show why the conditional curvature branch remains worth completing, while the bare and RRHF baselines are inadequate in the third-family and CP entries.

Theorem 2 applies to either set of canonical mass frames. Native recording preserves the polar mixing of the conditional coherent construction just as it preserves the baseline. Thus recording neither supplies the missing coherent axle on the baseline nor destroys its polar effect once the axle is present under the stated assumptions. Its positive vertex factor remains a separate observable obligation.

A consistent finite completion is obtained by inserting either pair U_u, U_d into Proposition 1 and applying the native recording map in the resulting canonical norm. This returns masses, a current and an explicit record response together. It does not derive their common-action feedback. In particular, it would be circular to tune X to offset the positive filter and then describe the adjusted result as a prediction.

8 The bath metric and the axle amplitude

The coherent source cannot be treated as wholly unspecified. PMS gives a nonzero cubic invariant, with z_E the E_1 coordinate and w the E_2 response:

$$\Gamma(w,z_E) = (\kappa_E/2)|w|^2 + g_{ax}\,\mathrm{Re}(w^* z_E^2) + \ldots, \tag{26}$$

$$g_{ax} = 1/2, \quad w_* = -(g_{ax}/\kappa_E)z_E^2. \tag{27}$$

The value 1/2 is the printed source-coordinate jet [6, EX82]. Whether it is a per-factor coefficient or already the normalised common coefficient must be checked when composing carriers; a diagonal average of three independent copies would introduce $1/\sqrt{3}$. This is an input-typing obligation, not a reason to require VERSF to have only one possible construction.

PMS also establishes that the relevant R4 and charged source planes coincide, so their retained norm is not to be multiplied by the unrelated $\sqrt{(5/6)}$ attenuation. The surviving quadratic maintenance coefficient is $\kappa_E = 2 + 4\omega$, with $\omega = e_\Sigma^T M_{short} e_\Sigma$ and $e_\Sigma = (1,1,1)/\sqrt{3}$ [6, EX115–EX123]. Later work makes the direct factor geometry explicit [7,8]:

$$H = (1/49)[36,18,36;\ 18,26,26;\ 36,26,47], \quad e_\Sigma^T H e_\Sigma = 269/147. \tag{28}$$

$$M_0 = K_{cyc} H, \quad K_{cyc} = 0.212610974570295, \tag{29}$$

$$M_{short} = \lambda_J M_0 - q_C q_C^T, \quad 0 \le e_\Sigma^T M_{short} e_\Sigma \le \lambda_J e_\Sigma^T M_0 e_\Sigma. \tag{30}$$

Here q_C encodes the allowed Abelian-closure portal and positivity of the full block is imposed. These formulas already include correlations between the force factors. The terminal-Born factor shape is not a licence to choose a new target metric or to rescale it until the desired axle amplitude is obtained.

The current-time paper proves a conditional path-to-record information bound: if the terminal law is a coarse-graining of the canonical current-cycle history with the same factor quotient, then $\lambda_J K_{cyc} \le 2/31$. Two separate irreversible path-law coefficients are also printed. Contracting those inherited values with Eq. (28) gives Table 6. This contraction is performed here; the underlying path calculations are inherited.

Table 6. Conditional current-history envelopes

History branch	I_{cyc}	λ_J	Upper κ_E	Lower $\lvert w_*\rvert$ for $\lvert z_E\rvert = 1$
canonical cochain	0.064516129	0.303446843	2.472240509	0.202245695

History branch	I_{cyc}	λ_J	Upper κ_E	Lower $\lvert w_* \rvert$ for $\lvert z_E \rvert = 1$
irreversible direct	0.055129628	0.259298129	2.403533875	0.208027024
irreversible reduced	0.050481473	0.237435877	2.369510648	0.211014034

For the canonical branch the exact envelope is $2 \le \kappa_E \le 11266/4557 = 2.472240509$ and, when $g_{ax} = 1/2$ is the common source coefficient, $0.202245695 \le |w_*| \le 1/4$ at unit $|z_E|$. The portal can only lower κ_E and therefore increase this induced response. It cannot raise the upper stiffness to the unit-residual value 26/5. The two metrics belong to different declared constructions and cannot be joined by calling them a change of coordinates.

The canonical information bound is conditional on the path-to-terminal attachment. The strict separable canonical history generator itself is current-cycle blind. The irreversible branch has a nonzero score but is not established as the required positive heat transfer [8, Section 24.31]. Table 6 is therefore a useful restriction on those attachments, not a physical axle determination. Neither w nor its bound is numerically equal to ζ: the physical charged-completion response that maps between them is still required.

This exposes two distinct uses of the bath. Its magnitude-sensitive history cost contributes to the coherent axle through κ_E; its completed recording operation contributes to the weak vertex through B. Adding the same fluctuation contribution to both as if they were independent corrections would need a derivation from a joint action. The present calculation supplies the second operator explicitly and leaves the first on its correctly typed source branch.

9 A common kinetic action with changing mass frames

The fixed-frame theorem leaves open a different operation: changing the kinetic form while holding the source Yukawa tensors fixed. Canonical normalisation then changes the mass frames themselves. The statement in Section 2 that the canonical current is the identity does not remove this effect. Proposition 1 compensates it by changing the source tensors together with Z_Q; the calculation here does not make that compensation.

9.1 A declared action and its complete quadratic return

Consider a local effective construction with the existing three quark doublets, the existing right-handed singlets, and a complex, gauge-singlet bath coordinate w. No fourth quark family is introduced. Its homogeneous potential uses the source form of Section 8. The right-handed kinetic forms are identity, and the left form is positive for every finite w:

$$U(w,z_E) = (\kappa_E/2)|w|^2 + g_{ax}\,\mathrm{Re}(w^*\, z_E^2),$$
$$G(w) = (\eta/\sqrt{3})(wE_{23} + w^*E_{32}), \quad Z_Q(w) = \exp(2G(w)). \tag{31}$$

$$\mathcal{L}_Q = (i/2)\,\bar{Q}_L Z_Q(w)\gamma^\mu \leftrightarrow D_\mu Q_L$$
$$- [\bar{Q}_L Y_u{}^0 H^c u_R + \bar{Q}_L Y_d{}^0 H d_R + \text{h.c.}]. \tag{32}$$

The bath has its ordinary positive scalar kinetic term; H^c is the charge-conjugate Higgs doublet. The symmetrised derivative makes the kinetic term Hermitian when w varies. This is a declared gauge-covariant effective action, not a recovery of the complete primitive VERSF action. In particular, the exponential kinetic embedding, its coefficient η, the 2↔3 readout and identity right-handed residues are construction choices. The source potential is inherited in form; selecting its conditional benchmark below does not establish its physical normalisation. Only the homogeneous, zero-fermion-background tree-level saddle is executed. Scalar loops and feedback from occupied matter are outside this calculation.

Stationarity and canonical normalisation give all the mass and current data on the same branch:

$$w_* = -(g_{ax}/\kappa_E)z_E^2, \quad C = \exp(-G(w_*)),$$
$$Y_f^{can} = CY_f^0, \quad J_W{}^{raw} = Z_Q, \quad J_W{}^{can} = CZ_QC = I. \tag{33}$$

For nonconstant w the canonical action also contains the anti-Hermitian connection $(C^{-1}\partial_\mu C - \partial_\mu C C^{-1})/2$. It vanishes on the saddle used here. Both the kinetic correction and the weak current therefore come from differentiating the displayed action; no independent current matrix is appended. The complete recording channel can subsequently act in this canonical norm, as in Section 10.

Proposition 4. At G = 0, with fixed source tensors Y_f^0, the common-action first variation is

$$\delta Z_Q = \delta J_W{}^{raw} = 2\delta G, \quad \delta Y_f^{can} = -\delta G\, Y_f^0,$$
$$\delta S_f = -\{\delta G, S_f\}, \quad S_f = Y_f^0 Y_f^{0\dagger}. \tag{34}$$

For nondegenerate singular values, write $S_f = U_f D_f^2 U_f^\dagger$ and choose the diagonal eigenvector phase derivative to vanish. With $U_f(\varepsilon) = U_f(I + \varepsilon A_f) + O(\varepsilon^2)$,

$$(A_f)_{ij} = (U_f^\dagger \delta S_f U_f)_{ij}/(d_{f,j}^2 - d_{f,i}^2), \quad i \neq j,$$
$$\delta V = -A_u V + V A_d, \quad \delta \log d_{f,i} = -(U_f^\dagger \delta G\, U_f)_{ii}. \tag{35}$$

Proof. Differentiate exp(−G) at zero and then the two-sided mass square. Differentiation of $U_f^\dagger S_f U_f = D_f^2$ fixes each off-diagonal entry of A_f; differentiation of $U_u^\dagger U_d$ gives δV. The diagonal mass-square derivative gives the last identity. The current identity holds exactly for finite G, not merely to first order. ▫

Thus a positive common kinetic form can change CKM mixing at fixed source Yukawas. It leaves the explicit canonical current matrix equal to I, while changing the two mass eigenframes differently. This is not a contradiction of Theorem 2, whose frames are fixed. Since Tr G = 0 in this construction, it also gives the exact correlated mass constraint

$$\prod_{i=1}^{3} d_{f,i}(G)/d_{f,i}(0) = \det(\exp(-G)) = 1, \quad f = u,d. \tag{36}$$

9.2 Fixed input execution and sensitivity

Use the two bare frames of Section 7 and the inherited singular values in [4], kept fixed in $Y_f^0 = U_f D_f$. The inputs are

$$D_u = \mathrm{diag}(0.000007214053900982, 0.003704884912807, 0.9725186481825),$$
$$D_d = \mathrm{diag}(0.00001563045011879, 0.0003127500251071, 0.01620450741891). \tag{37}$$

For an executable benchmark take $\kappa_E = 11266/4557$, $g_{ax} = 1/2$ in the common-source convention, $|z_E| = 1$ and η = 1. This uses the conditional upper envelope in Table 6, not a selected physical stiffness. The kinetic off-diagonal magnitude is then $|w_*|/\sqrt{3} = 0.1167666065$. The source phases $\arg z_E = m\pi/3$ give the three distinct w phases 60°, 180° and 300°; all are reported, with none selected by its agreement with quark data. The full finite calculation exponentiates Z_Q, canonicalises both tensors and obtains their singular values and left mass frames by singular-value decomposition.

Table 7. Common kinetic action with fixed bare source tensors

Bath phase	$\lvert V_{ub}\rvert$	$\lvert V_{cb}\rvert$	$\lvert V_{td}\rvert$	$\lvert J\rvert \times 10^5$
No kinetic feedback	0.003931186	0.040284873	0.006108170	1.843506
60°	0.003869956	0.039166850	0.006011560	1.770078
180°	0.003875460	0.039286492	0.006052644	1.796030

Bath phase	$\lvert V_{ub}\rvert$	$\lvert V_{cb}\rvert$	$\lvert V_{td}\rvert$	$\lvert J\rvert \times 10^5$
300°	0.003887550	0.039168811	0.006012357	1.811523

These are nonzero coherent changes with a canonical unitary weak current. They do not repair the deficient bare third-family pattern. At 60° the up-sector mass ratios are (0.999796344, 0.984493198, 1.015957957), and the down-sector ratios are (0.999826414, 0.989026666, 1.011270626). A change in mixing therefore comes with a correlated change in masses; those masses have not been refitted.

The calculation is also applied to the inherited AX1 source tensors, with their phase and amplitude conditions retained. Table 8 varies η at fixed 60° source phase. The full archive includes all three phases and both starting branches. The aim is to expose the effect of the kinetic choice, not to optimise it.

Table 8. Sensitivity of the inherited AX1 branch to the kinetic coefficient

η	$\lvert V_{ub}\rvert$	$\lvert V_{cb}\rvert$	$\lvert V_{td}\rvert$	$\lvert J\rvert \times 10^5$
0	0.003556877	0.040537720	0.008689898	2.996768
0.25	0.003551168	0.040458407	0.008677312	2.987961
0.5	0.003539372	0.040242867	0.008649936	2.968720
1	0.003498143	0.039419730	0.008552257	2.900389
2	0.003352786	0.036426532	0.008203332	2.662312

Canonical-current residuals are below 10^{-14}, and the mass-product identity is verified to 10^{-12} by double-precision SciPy singular-value decomposition (LAPACK gesdd): sort the three singular values of each $\exp(-G)Y_f^0$, form their product relative to the inherited values, and compare with $\exp(-\mathrm{Tr}\, G) = 1$. The largest absolute product residual is 1.1×10^{-14}. Central finite differences of $|V_{ij}|^2$ agree with Proposition 4 to below 10^{-9} over steps 10^{-3}, 10^{-4} and 10^{-5}. Replacing the fixed source tensors by the compensating tensors of Proposition 1 restores the original mixing. This separates a physical parameter change within the declared action from a mere field redefinition. It does not establish that η = 1, or this kinetic embedding, is the admitted VERSF completion.

10 The complete recorded weak response

A mean current is not a transition probability. This distinction can be closed exactly for the native channel, without introducing a new noise strength. Specify the operational order: prepare a canonical up-type mass state with upper weak isospin and colour density $I_C/3$; let the unobserved, gauge-averaged recording channel act; then apply a charged weak probe and resolve the final down-type mass state. The event source is a fixed equal mixture of the three coordinate rays e_1, e_2, e_3 in the source frame in which Ω is written in Eq. (22). These are the recording rays, not the quark mass eigenvectors. In all baseline tables the recording-basis matrix is R = I. Changing these projectors while holding the mass frames fixed changes the physical model, not just its coordinates. Although every such orthonormal ensemble has average density I/3, its averaged recording channel need not be the same, because the channel depends nonlinearly on each reference ray. Its occurrence follows the independent Poisson model already declared in Section 5. The result below is the complete finite channel-and-probe return for that preparation. Its identification with a particular decay experiment still requires the physical preparation, event exposure and energy-dependent current matching.

10.1 The second moment and the missing strength

For a unital completely positive Heisenberg map T, define the current noise operator

$$N_T(J) = T(J^\dagger J) - T(J^\dagger)T(J) \geq 0. \tag{38}$$

The inequality is the completely positive Schwarz inequality [15]. A direct dilation proof writes $T(J) = V_{dil}^\dagger(J\otimes I)V_{dil}$, $V_{dil}^\dagger V_{dil} = I$, so that $N_T(J) = \ell_J^\dagger \ell_J$ with $\ell_J = (I - V_{dil}V_{dil}^\dagger)(J\otimes I)V_{dil}$. It separates the current component returning to the encoded subspace from its orthogonal complement. No extra noise coefficient may be added after the channel is specified.

Proposition 5. For $J_a = I_F\otimes I_C\otimes\sigma_a/2$ and the one-event isotropic channel of Section 3,

$$\Sigma_{a=1}^3 T(J_a^2) = (3/4)I,$$
$$\Sigma_{a=1}^3 N_T(J_a) = (3/4)(I - B^2)\otimes I_C\otimes I_W. \tag{39}$$

Proof. The sum of the three squared weak generators is $(3/4)I$ and unitality preserves it. The first moment is $B\otimes\sigma_a/2$, whose summed square is $(3/4)B^2$. Subtracting gives the result. ▫

The loss from the squared mean current is therefore not automatically a loss of total inclusive weak-current strength. The missing part is fixed by the same operation. The identity sums all three weak-current components; a charged probe with a specified initial isospin has a different, calculable total, given below.

10.2 Complete family and weak response

For an orthonormal recording basis with columns r_k of R, define $D_R(F) = \Sigma_k |r_k\rangle\langle r_k|F|r_k\rangle\langle r_k|$ and $O_R(F) = F - D_R(F)$. Baseline Tables 9–12 use $R = I$, so D_R keeps the diagonal in the Ω source frame. The mass frames and recording projectors are transformed together under a coordinate change; varying R alone is a distinct source choice. Besides the weak-adjoint map Φ, the gauge-averaged channel has a weak-singlet family map Ψ. For a product family source, $\Psi(PFP + QFQ) = PFP + QFQ$ and its off-diagonal coefficient is the same α as in Eq. (16). Averaging the three orthogonal family sources gives

$$\Psi(F) = D_R(F) + [(1 + 2\alpha)/3]O_R(F),$$
$$\Phi(F) = [(2 + \beta)/3]D_R(F) + [(1 + 2\alpha)/3]O_R(F). \tag{40}$$

Exponentiating the actual two superoperators under the common Poisson clock yields

$$\Psi_\tau(F) = D_R(F) + r(\tau)O_R(F), \quad \Phi_\tau(F) = b_W(\tau)D_R(F) + r(\tau)O_R(F),$$
$$b_W(\tau) = \exp[-(1 - \beta)\tau/3], \quad r(\tau) = \exp[-2(1 - \alpha)\tau/3]. \tag{41}$$

The scalar $b_W(\tau)$ is the common weak-current attenuation. The auxiliary singlet map Ψ is not obtained by multiplying weak-current matrices B. Its different diagonal evolution is essential to the inclusive result.

Let $P_j^d = |d_j\rangle\langle d_j|$, $J_- = I_F\otimes I_C\otimes|\downarrow\rangle\langle\uparrow|$ and $V = U_u^\dagger U_d$. The final-flavour effect of the charged probe is $P_j^d\otimes I_C\otimes|\uparrow\rangle\langle\uparrow|$. Decomposing the weak projector into singlet and adjoint parts gives the exact transition weight

$$P_{ij}(\tau) = (1/2)\langle u_i|[\Psi_\tau(P_j^d) + \Phi_\tau(P_j^d)]|u_i\rangle$$
$$= r(\tau)|V_{ij}|^2 + [A(\tau) - r(\tau)]D_{ij}, \tag{42}$$

$$D_{ij} = \Sigma_k|(R^\dagger U_u)_{ki}|^2\,|(R^\dagger U_d)_{kj}|^2, \quad A(\tau) = [1 + b_W(\tau)]/2,$$
$$\Sigma_j P_{ij}(\tau) = \Sigma_i P_{ij}(\tau) = A(\tau). \tag{43}$$

The complement $1 - A$ is the weight transferred to the lower isospin before the lowering probe; it is not discarded by the preparation-and-probe accounting. Conditional on an allowed lowering transition, $W_{ij} = P_{ij}/A$ is a doubly stochastic probability matrix. It is a convex combination of $|V|^2$ and D, since $0 \leq r \leq A$. This does not assign a coherent phase to W or promote $\sqrt{W}$ to a CKM amplitude matrix.

For the total charged current, the same dilation identity gives a separate exact budget:

$$\langle\uparrow|T_\tau(J_-^\dagger J_-)|\uparrow\rangle = A(\tau),$$
$$\langle\uparrow|T_\tau(J_-)^\dagger T_\tau(J_-)|\uparrow\rangle = b_W(\tau)^2, \quad N_{tot}(\tau) = A(\tau) - b_W(\tau)^2. \tag{44}$$

The second line is the coherent encoded-subspace contribution. Its difference from the inclusive total is nonnegative and fully determined. This is a total budget; it is not assigned separately to the flavour outcomes unless the flavour readout and encoded-subspace projector are shown to commute.

Table 9. Complete charged-current strength budget at maximal native activity

Exposure τ	Inclusive A	Squared mean b_W^2	Orthogonal contribution
0	1.000000000	1.000000000	0.000000000
0.01	0.999370641	0.997484148	0.001886493
0.1	0.993741937	0.975124402	0.018617535
1	0.940828969	0.777320719	0.163508250
3	0.842665491	0.469678555	0.372986936

At $\tau = 1$, squaring the mean current gives 0.777320719, whereas the complete charged-probe total is 0.940828969. The difference, 0.163508250, is the contribution omitted by using the squared first moment alone. This is an executed response of the specified protocol, not a fitted correction to a force strength.

Table 10. Square roots of conditional probe probabilities for comparison only

Source branch	τ	$\sqrt{W_{ub}}$	$\sqrt{W_{cb}}$	$\sqrt{W_{td}}$
Bare	0	0.003931186	0.040284873	0.006108170
Bare	0.1	0.003920729	0.040175370	0.006086783
Bare	1	0.003828864	0.039213105	0.005898095
Inherited AX1	0	0.003556877	0.040537720	0.008689898
Inherited AX1	0.1	0.003549562	0.040572875	0.008663739
Inherited AX1	1	0.003485491	0.040876421	0.008433528

The entries in Table 10 are not CKM matrix elements away from $\tau = 0$. Recording preserves the polar factor of the mean current but changes the complete transition probabilities. In particular, the conditional AX1 branch remains a useful input; its experimental assessment must use the full probability return, not only the unchanged polar factor.

10.3 Combined execution and physical scope

The two calculations compose without changing their meaning. Use the canonical frames returned by the action in Section 9 inside the protocol above. For the inherited AX1 source, $\eta = 1$, phase 60° and $\tau = 0.1$, the returned charged-probe matrix is

$$P = [0.943111883, 0.050614228, 0.000015827;\\ 0.050553804, 0.941360025, 0.001828108;\\ 0.000076250, 0.001767684, 0.991898003]. \tag{45}$$

Every row and column sums to 0.993741937 at full precision. The same run returns all six mass shifts, a unitary canonical CKM matrix, the mean current and its inclusive fluctuation budget. The maximal matrix-

map residual is 2×10^{-16}; independent 9×9 superoperator exponentials reproduce the time laws. These checks verify the finite construction, not the source admission of its coefficients.

A lepton doublet provides an additional representation check. With no colour factor the same native construction gives $\alpha_L = \sqrt{(1 - \chi)}$ and $\beta_L = 1 - 4\chi/3$. At maximal activity $\beta_L = -1/3$, so the quark positivity theorem must not be transferred to leptons. Its equal-family Poisson response remains well defined, with $b_L(\tau) = \exp(-4\tau/9)$. Any physical extraction normalised to muon decay must calculate the leptonic preparation and complete current correlator as well; dividing two squared mean currents would omit the noise just derived.

The physically measurable contribution depends on which final records and excitation energies the experiment includes. The finite channel fixes the total fluctuation weight, but it has no energy-resolved bath spectrum. A decay-rate prediction therefore requires the corresponding frequency-dependent correlation function and thresholds. The recording basis R, kinetic coefficient η and exposure τ remain declared choices; none is inferred from measured quark mixing. The complete protocol is a constructive return, while its identification with the physical VERSF process remains conditional.

10.4 First-row normalisation constraint

The independently extracted first-row CKM sum is 0.9984 ± 0.0007 in the cited PDG review [11, Section 12.4]. It is an absolute normalisation test; the conditional matrix $W = P/A$, whose rows sum to one by definition, cannot answer it. At $\tau = 0.1$ the quark-probe total is $A = 0.993741937$, a loss of 0.6258%. This loss is much larger than the quoted first-row uncertainty and must be included in any physical comparison.

The comparison also depends on the Fermi constant inferred from muon decay. In a factorised, process-universal approximation write $S_{ext} = C_{norm}A$, with $C_{norm} = L_{SL}/L_\mu$; L_{SL} and L_μ are complete rate factors for the semileptonic lepton part and muon decay, relative to the same underlying weak normalisation. They must include their accepted final states, fluctuations and energy dependence. Neither is a squared mean current. This paper does not calculate those process factors, so the first-row measurement is not yet an unconditional bound on τ.

For the explicit diagnostic $C_{norm} = 1$, Eq. (43) gives $A = [1 + \exp(-a_d\tau)]/2$ with $a_d = 0.125951124$. Then $\tau = 0.1$ lies 6.65 quoted standard deviations below the measured central sum. The central sum corresponds to $\tau = 0.02545$; its two-standard-deviation lower endpoint corresponds to $\tau = 0.04778$. Requiring the additional deficit alone to be no larger than one quoted standard deviation gives $\tau \le 0.01112$. Thus exposures of order 10^{-2}, rather than 0.1 or 1, are the relevant small-loss scale under unchanged relative leptonic normalisation. These are translations of a stated diagnostic assumption, not a rate fitted for use as a prediction. Tables at larger τ illustrate the mechanism and are not phenomenologically accepted points. A changed C_{norm} requires the full muon-decay and semileptonic calculation before any revised bound is claimed.

11 Propagation between recording events

The recording-only result is not a theorem about recording interleaved with propagation. A finite construction can use the same quark frames, inherited singular values and native event map, with the recording projectors specified relative to those frames, and return a nonzero coherent correction. This section solves that construction. It is a spin-suppressed, positive-energy flavour model; a decay calculation must still restore spin factors, phase space and the energy-resolved environment.

11.1 One continuous generator for propagation and recording

At a fixed dimensionless momentum p, let the canonical source masses of Section 7 or 9 define

$$H_f(p) = U_f\,\mathrm{diag}(\sqrt{(p^2 + d_{f,1}^2)}, \sqrt{(p^2 + d_{f,2}^2)}, \sqrt{(p^2 + d_{f,3}^2)})U_f^\dagger,$$
$$H_q = H_u \otimes P_\uparrow + H_d \otimes P_\downarrow. \tag{46}$$

The dispersion is the positive-energy branch of the canonical quadratic action at constant w; the reduced current here omits the spinor kinematic factors. The inherited mass unit is E_*; p is measured in E_*/c. On the reduced family–weak carrier, let E be the gauge-averaged native channel defined by $E(F\otimes I) = \Psi(F)\otimes I$ and $E(F\otimes\sigma_a) = \Phi(F)\otimes\sigma_a$. With $\tau = \gamma t$ and $\omega = E_*/(\hbar\gamma)$, the joint generator is

$$d\rho/d\tau = L(\rho), \quad L(\rho) = -i\omega[H_q,\rho] + E(\rho) - \rho. \tag{47}$$

Both terms act on the same density matrix. The collision term follows by coupling each event to a fresh record through the isometry $|\psi\rangle \mapsto \Sigma_\alpha C_\alpha|\psi\rangle\otimes|\alpha\rangle$ and retaining the record labels. The sum $C_\alpha^\dagger C_\alpha = I$ preserves the total weight. A Poisson collision process interleaved with unitary propagation gives the displayed generator. Its propagator is completely positive and trace preserving by the product limit of those two operations. The event map, source ensemble and mass inputs are fixed before comparison with CKM data. The common rate γ is a declared dynamical input, not an extracted physical clock.

Only an 18-state colour-independent density description is needed before tracing colour, and the reduced family–weak carrier has six states. Its complete density generator is a 36 × 36 matrix. We exponentiate it directly. The same matrix gives the state evolution, the Heisenberg current, its polar factor and the inclusive flavour probabilities.

11.2 An exact leading coherent correction

Write the lowering block of the Heisenberg current $T_\tau(J_-)$ as $B_{du}(\tau)$, with $B_{du}(0) = I$. Let $\mathrm{Off}_R(F) = F - D_R(F)$ remove the diagonal in the chosen recording basis, with R = I for Tables 11–12, and define

$$a_d = (1 - \beta)/3, \quad a_o = 2(1 - \alpha)/3,$$
$$\delta_r = a_o - a_d = 0.0463277936554 > 0. \tag{48}$$

$$dB_{du}/d\tau = i\omega(H_d B_{du} - B_{du}H_u) - a_d B_{du} - \delta_r\,\mathrm{Off}_R(B_{du}). \tag{49}$$

Proposition 6. Remove the common attenuation and the free propagation by defining $\tilde{B} = \exp(a_d\tau)\exp(-i\omega H_d\tau)\,B_{du}\exp(i\omega H_u\tau)$. Then

$$\tilde{B}(\tau) = I - (i\omega\delta_r\tau^2/2)\,\mathrm{Off}_R(H_d - H_u) + O(\tau^3). \tag{50}$$

Proof. The first derivative of B_{du} at zero is $i\omega(H_d - H_u) - a_d I$. Differentiating again, the extra off-diagonal damping contributes $-i\omega\delta_r\mathrm{Off}_R(H_d - H_u)$. The common attenuation and free propagation remove the other zeroth-, first- and second-order terms. The remaining coefficient is anti-Hermitian, so it enters the polar unitary at this order. ▫

The mechanism is the failure of propagation to commute with the family dephasing induced by recording. A fixed positive weighting is insufficient to represent the interleaved evolution. Theorem 2 and the repeated product-state theorem remain valid for their recording-only maps; the joint generator is outside that restriction. The source-basis correction can contain removable phases in special cases, so genuine changes of rephasing-invariant moduli are checked explicitly below. At ω = 0 the correction vanishes exactly.

For the bare branch at ω = 1, the Frobenius norm of the interaction-picture remainder divided by τ^3 is 0.00020805, 0.00020808, 0.00020809 and 0.00020809 at τ = 0.1, 0.03, 0.01 and 0.003. This independently checks the predicted second-order term against the full exponential.

11.3 Executed coherent and inclusive returns

Tables 11–12 hold the inherited AX1 construction fixed, take $\eta = 0$, $p = 0$ and $\tau = 1$, and use the three Ω-source-frame recording rays, R = I. The recording basis R is a declared choice alongside η, ω and the exposure. The tables vary ω over a fixed logarithmic grid. The unitary entries refer to polar($U_u^\dagger B_{du} U_d$); its free propagation phases do not affect the displayed moduli. No grid point is selected as the physical rate.

Table 11. Coherent polar factor with propagation between events

ω	$\lvert V_{ub}^{pol}\rvert$	$\lvert V_{cb}^{pol}\rvert$	$\lvert V_{td}^{pol}\rvert$	$\lvert J_{pol}\rvert \times 10^5$
0	0.003556877	0.040537720	0.008689898	2.996768
0.1	0.003560016	0.040559456	0.008689081	2.999533
1	0.003580560	0.040498541	0.008634622	3.000433
10	0.003526413	0.038732584	0.008319061	2.839840
100	0.003444378	0.038402967	0.008314174	2.767087

The complete probe probability is obtained separately from the evolved state, $P_{ij}(\tau) = \mathrm{Tr}[(P_j^d \otimes P_\uparrow) \exp(\tau L)(P_i^u \otimes P_\uparrow)]$. The result is not the square of Table 11. The total lowering weight remains $A(\tau) = [1 + \exp(-a_d \tau)]/2$, because the Hamiltonian commutes with the upper-isospin projector and the channel fixes its scalar decay law. Every row and column has that sum. The finite coherent and inclusive returns are therefore compared on identical preparations.

Table 12. Complete conditional probe response and relaxation gap

ω	$\sqrt{W_{ub}}$	$\sqrt{W_{cb}}$	$\sqrt{W_{td}}$	Decay gap in τ
0	0.003485491	0.040876421	0.008433528	0.000e+00
0.1	0.003484203	0.040826722	0.008435475	7.142e-07
1	0.003616687	0.040968940	0.008563963	4.305e-06
10	0.004489530	0.044892379	0.009296257	1.340e-04
100	0.004424613	0.045008002	0.009302773	4.696e-04

As in Table 10, $\sqrt{W}$ is a probability comparison and is not assigned a CKM phase. The archived calculation also uses the bare branch, $p = 1$, and the $\eta = 1$ kinetic example, so the dependence on propagation, momentum and the chosen kinetic embedding is exposed. The one-event channel has zero trace-preservation and unital residuals, and its minimum Choi eigenvalue is -9×10^{-16}, consistent with numerical zero. Evolved states and current-noise budgets remain positive to numerical precision; all probability row and column residuals are below 10^{-12}. The rate scan changes the results appreciably and does not establish a simultaneous agreement with all measured quark observables.

11.4 Recording-basis sensitivity

Keep the inherited AX1 mass frames, $\eta = 0$, $p = 0$, $\chi = 1$, $\omega = 1$ and $\tau = 1$ fixed. Change only the recording ensemble: R = I uses the Ω source frame, $R = U_u$ the up-quark mass frame, and $R = U_d$ the down-quark mass frame. Equation (40) is conjugated with this R before exponentiating the full generator. Table 13 reports both the polar current and one complete conditional probe entry; Δ_{pol} is the maximum absolute change of any polar modulus from $|U_u^\dagger U_d|$.

Table 13. Recording-basis dependence at fixed dynamics

Recording basis	$\lvert V_{ub}^{pol}\rvert$	$\lvert V_{cb}^{pol}\rvert$	Δ_{pol}	$\sqrt{W_{ub}}$
Ω source frame	0.003580560	0.040498541	5.528e-05	0.003616687
Up-quark mass frame	0.003556871	0.040537645	7.524e-08	0.003556877
Down-quark mass frame	0.003533379	0.040270383	2.673e-04	0.004162808

The zero-propagation polar baseline is $|V_{ub}| = 0.003556877$ and $|V_{cb}| = 0.040537720$ for all three recording bases. At $\omega = 1$ the up-frame propagation correction is only 7.52×10^{-8} in the largest modulus, compared with 5.53×10^{-5} in the source frame and 2.67×10^{-4} in the down frame. The up-frame effect is strongly suppressed at this benchmark, not identically zero: at $\omega = 10$ its largest modulus shift is 7.51×10^{-6}. The archive includes the full rate grid and exposures 0.01, 0.1 and 1; at $\tau = 0.01$ and $\omega = 1$ the source-frame shift falls to 2.55×10^{-8}.

This dependence has a useful analytic control. With propagation switched off, $R = U_u$ or $R = U_d$ makes $D_{ij} = |V_{ij}|^2$ in Eq. (43), so $W_{ij} = |V_{ij}|^2$ exactly at every exposure. For propagation restored, the leading polar correction in either mass-aligned recording basis has only row and column phase effects on V; higher orders can change the moduli. In the up basis the non-diagonal Hamiltonian is H_d, whose inherited energy splittings are much smaller than those of H_u at $p = 0$. This explains the strong finite-time suppression, without claiming an all-orders cancellation. A simultaneous coordinate rotation leaves the response invariant; rotating the recording projectors alone does not.

The open physical question is therefore explicit: which family projectors does the VERSF recording source select relative to the quark mass frames, and what source dynamics supplies that selection? Equal ensemble weights alone do not answer it. The paper establishes an available basis-dependent mechanism; it does not establish that VERSF forces the source-frame choice or its numerical correction.

11.5 Long time behaviour

Proposition 7. For maximal activity and $\omega > 0$, form the graph with one vertex per family and an edge ij whenever $(R^\dagger H_u R)_{ij}$ or $(R^\dagger H_d R)_{ij}$ is nonzero. If this graph is connected, the joint semigroup has the unique stationary state $I_6/6$ and converges to it. Consequently,

$$P_{ij}(\tau) \to 1/6, \quad A(\tau) \to 1/2, \quad W_{ij}(\tau) \to 1/3 \quad \text{as } \tau \to \infty. \tag{51}$$

Proof. The recording channel is self-adjoint in the Hilbert–Schmidt inner product. Its fixed space is the set $D \otimes I$ with D diagonal in the recording basis; every other isotype has eigenvalue strictly below one. The real part of the quadratic form of L is therefore nonpositive and vanishes only on this fixed space. A stationary operator must also commute with H_q. In recording coordinates, for a diagonal D, commutation gives $(D_{ii} - D_{jj})(R^\dagger H_f R)_{ij} = 0$, so graph connectivity makes D scalar. The same argument excludes nonzero imaginary-axis eigenvalues: the commutator of H_q with $D \otimes I$ has zero diagonal in recording coordinates and cannot equal a nonzero multiple of $D \otimes I$. Complete positivity excludes growing or Jordan modes on the imaginary axis. In finite dimension, convergence follows. ▫

Both displayed quark branches satisfy the graph condition. The smallest nonzero decay rate is given in Table 12; it can be much smaller than the elementary event rate. This supplies a controlled dynamical statement: a finite amount of interleaved evolution can change the mixing response, while indefinitely repeated unconditioned evolution in this model erases the family distinction. Physical preparation and renewal must therefore be included before a stationary output is identified with the quark mixing observed in experiments.

12 Consequences for the November ledger

The paper returns a solved finite quark evolution rather than only a fixed-frame compatibility test. It calculates a common-action kinetic response, the complete channel-and-probe probabilities, and the coherent correction created by propagation between native recording events. The physical distinction is exact: recording alone preserves the quark polar factor in the proved class, while its interleaving with propagation need not. The full response also contains a positive fluctuation contribution that is absent from the squared mean current. These close mathematical response and finite-execution tasks; they do not by themselves close the physical source obligations of SM-17 or SM-22.

The unresolved task is now to attach the executed dynamics to the physical source: return the kinetic embedding and completion response, the recording-source family projectors, the record-conditioned family preparation, the common event rate and the energy-resolved bath correlator. A physical rate comparison must also pass the muon-normalised first-row constraint of Section 10.4. The master equation can then be evaluated at those inputs rather than scanned. The favourable AX1 source remains conditional; neither a source phase, η nor ω may be selected by reproducing the same quark observables subsequently presented as predictions.

Table 14. Required physical return

Returned object	Use in the calculation	Acceptance test
Same-action δY_u, δY_d	Changes the two canonical mass frames	Zero feedback recovers the selected RRHF branch
δZ_Q and the weak current	Canonicalises the same mass and current data	Common covariant kinetic correction obeys Eq. (1)
E_2 current score and Abelian portal	Determines κ_E in one source norm	Returns the H shape and passes the appropriate path bound
Completion response from w to X	Gives a coherent axle without a target fit	Reproduces the mixed jet and fixes all conversion gains
Recording-source family projectors	Fixes R and D_R, hence the relative recording and mass frames	R = I recovers Tables 10–12; mass-frame choices recover Table 13
Joint state, rate and bath spectrum	Selects ω, τ and the included fluctuation strength	Recovers the native channel and passes the muon-normalised first-row constraint

Explicit failure tests follow from the results. A claimed native one-event rotation of the polar mixing factor at fixed mass frames contradicts Theorem 2 unless an assumption has changed. A non-unitary raw vertex presented as a CKM matrix fails the current test. A physical current metric inferred solely from the factor Gram fails the history attachment. A terminal coarse-graining that exceeds its own complete path Fisher information fails the score pushforward. A change in quark angles obtained by fitting the completion gain or source phase to those same angles is an empirical fit and must be labelled as such.

Conversely, one compatible common-action construction that returns all the objects in Table 14, meets the source and current checks, and predicts the unused flavour observables would materially advance the derivation. It need not be uniquely forced by the abstract principles. Its alternatives should be tested by varying declared construction choices while keeping calibration and prediction separate.

13 Conclusion

The native resolved recording operation has an exact, positive weak-vertex response on the quark-doublet carrier. After isotropic gauge projection it preserves the unitary mixing factor for every pure reference ray in one event, and for arbitrary finite sequences in the solved product-state class. This exact recording-only result is a control case for the propagation and kinetic calculations, not an exclusion of coherent corrections in the combined dynamics.

The joint evolution supplies a constructive extension. Interleaving the native channel with the inherited quark propagation generates a coherent polar correction at second order in the exposure, and the complete 36 × 36 evolution returns both that current and all flavour-resolved probe probabilities. A separate gauge-covariant kinetic action returns correlated mass and mixing changes. The recording basis, physical coefficients and preparation are displayed explicitly; their physical source selection remains open. The result is a solved conditional dynamical construction and a route to testing its physical attachment, not a completed numerical derivation of the observed CKM matrix.

Appendix A Product state coefficient

Choose product reference $c \otimes w = |0,0\rangle$. In the six-dimensional colour–weak space, A has eigenvalue 1/2 on three orthogonal excitations and zero on the remaining directions. Hence C_0 has diagonal entries (r,u,u,1,u,1), with r and u as in Section 5. The off-diagonal family coefficient is the normalised gauge trace of C_0, giving α in Eq. (16). The weak-adjoint trace of C_0 contributes

$$N_0 = ru + 2u + (r^2 + 3u^2 + 2)/4. \qquad (52)$$

The sum of the active native Kraus contributions is $\chi/6$. Dividing $N_0 + \chi/6$ by 9/2 gives β. At $\chi = 1$, $\beta - \alpha^2 = (2\sqrt{6} - 1)/54 = 0.072203324$, which is the extra positive Kraus weight in Eq. (18).

For comparison, on a family–single-fundamental pure ray, the completed-event paper gives $\rho_{F,out} = [(d + t)\rho_F - 2\rho_F^2]/(d - t)$, where ρ_F is the family marginal and $t = \mathrm{Tr}\,\rho_F^2$. It commutes with ρ_F, so that particular marginal update changes eigenvalues but not eigenvectors. For d = 3 and $\rho_F = \mathrm{diag}(0.6,0.3,0.1)$, the output is diag(0.533858268,0.337795276,0.128346457). This control is consistent with the distinction between reweighting and frame rotation, but is not the proof for the full quark doublet; Theorem 2 supplies that proof.

Appendix B Reproduction and source coverage

The companion archive Weak_Current_Quark_Mixing_Calculations.zip uses Python, NumPy and SciPy with double-precision matrices. It contains the generator construction, native Kraus operators, two independent weak-current contractions, polar decomposition, product-state family map, fixed repeated-event sequence, coherent bare and AX1 exponentials, the conditional current-history contractions, common kinetic action, complete second moments, the 36 × 36 propagation-and-recording generator, recording-basis sensitivity and the conditional first-row constraint. A public seed specifies the entangled numerical witnesses; no empirical quantity is used by the source generator. Results and the complete numerical source vector are stored alongside the script. All central formulas and analytic fixed inputs are printed here; the numerical witness arrays are in the companion archive.

The source review covered the public November ledger and the quark, bath, primitive-matter, current-rigidity, clock and completed-event papers cited below. It distinguishes the later magnitude-bearing current and state-level results from the older normalised-ray obstruction. Public PDF references are provided so that the scientific dependencies do not require access to a private conversation archive. The paper does not claim to have re-executed archived RRHF source arrays, the primitive path-law Monte Carlo or the Standard Model force-running codes.

The one-event theorem can be independently checked using only the displayed Kraus law, the SU(2) completeness identity and the operator inequalities in its proof. The empirical benchmarks do not enter that proof. The repeated-event theorem requires product reference rays and composes Φ as a superoperator. Numerical sampled entangled positivity is a verification of the implementation, not a replacement for the exact one-event bound.

References

[1] Keith Taylor. VERSF Standard Model Derivation Ledger November 2026. November 2026. Entries SM-13, SM-17, SM-22 and the shared-construction milestones. Public paper

[2] Keith Taylor. Bath or Ledger. Website edition May 2026. Public paper

[3] Keith Taylor. The Bath Criterion. Website edition June 2026. Public paper

[4] Keith Taylor. The Conditional Derivation of Charged Flavour and the Unique Lowest-Order Quadratic CP Axle in VERSF. Website edition August 2026. Sections 5.2 and 7; RRHF and RRHF-AX1 returns. Public paper

[5] Keith Taylor. The Full-Carrier Quark-Mixing Curvature Closure in VERSF. 20 August 2026. Public paper

[6] Keith Taylor. Primitive Matter Selection in VERSF: Cumulative Executed Theorems on Fact Commitment, Record Retention, Renewal and R4-to-Charged-Flavour Descent. 25 August 2026. Cumulative through PMS-EX126; EX82 and EX115–EX126. Public paper

[7] Keith Taylor. The Twelve-Sector Current-Rigidity, Source-Generated Rank-Six Current-Stiffness and Current-Link Schur Reciprocity Theorem in VERSF. September 2026. Public paper

[8] Keith Taylor. The Sequential Opportunity-to-Lorentzian Commitment-Time Attaching Map in VERSF. Website edition September 2026. Sections 24.14–24.31 and 25. Public paper

[9] Keith Taylor. Completed Event Weighting and Force Response in VERSF. 3 November 2026. Section 12.5, Proposition 13 and Table 18. Public paper

[10] Keith Taylor. Physical Bath Selection and Neutrino Mixing in VERSF. 25 September 2026. Companion example of the common mass–kinetic–current requirement. Public paper

[11] A. Ceccucci, Z. Ligeti and Y. Sakai. CKM Quark-Mixing Matrix. Review of Particle Physics, 2025 update of S. Navas et al., Physical Review D 110, 030001 (2024). Equations 12.1–12.2, 12.11–12.13 and Section 12.4. Public paper

[12] Keith Taylor. Mixed Matter Matching and the One-Calibration Force Test in VERSF. November 2026. Public paper

[13] Keith Taylor. The Quark-Sector Closure Ledger in VERSF. Website edition July 2026. Public paper

[14] Keith Taylor. Assignment, Bath Participation, and Refinement: A Response-Operator Account of the Quark Mass Hierarchy in VERSF. Website edition June 2026. Public paper

[15] Man-Duen Choi. A Schwarz inequality for positive linear maps on C*-algebras. Illinois Journal of Mathematics 18 (1974), 565–574. Public paper