

Nonlocal nonstabilizerness from holographic Schwinger pair production

Sebastian Grieninger^{*}

*InQubator for Quantum Simulation (IQUS), Department of Physics, University of Washington,
Seattle, Washington 98195, USA*



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We analyze the emergence of nonlocal quantum correlations in Schwinger pair creation in strong non-Abelian (chromo)electric fields using holography. The produced quark-antiquark pair is entangled into a color singlet, yet accelerates into causally disconnected Rindler wedges. Using the Casini-Huerta-Myers conformal mapping and the probe-brane framework, we compute the refined Rényi entropy and its derivative, which captures the antiflatness of the entanglement spectrum for a spherical bipartition. We find that for boundary spacetime dimension $d > 2$, the entanglement spectrum is nonflat, implying the dynamical generation of nonlocal nonstabilizerness in the pair creation process. Interestingly, the generation of nonlocal nonstabilizerness in the holographic dual is encoded in the free energy of the probe action.

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Introduction. The emergence of nonlocal quantum correlations in nonequilibrium processes is a central question in quantum many body systems [1]. While entanglement entropy quantifies the amount of correlations generated, it does not capture their internal structure.

In this work, we focus on nonlocal magic: the part of nonstabilizerness across a bipartition that cannot be removed by local unitary transformations. We use the term “magic” in this resource-theoretic sense, rather than as circuit complexity. While magic is closely related to the difficulty of classical simulation in stabilizer-based descriptions, it is not identical to algorithmic quantum complexity and should not be treated as a universal measure of computational hardness.

Recent developments in quantum information science have emphasized that entanglement alone does not fully characterize the structure of quantum many-body states [2–32]. While entanglement entropy captures the total amount of quantum correlations, it is insensitive to how these correlations are organized within the Hilbert space. Measures of quantum complexity, such as nonlocal magic [17,33], provide complementary information by probing the microscopic structure underlying these correlations.

In particular, nonlocal magic is sensitive to fluctuations and variance of the entanglement spectrum [15,17], revealing features that remain invisible to the entanglement entropy [26]. This distinction has proven important in a range of physical settings, where states with similar entanglement can exhibit vastly different internal structures and dynamical properties [34,35]. For example, nonlocal magic can detect nontrivial organization in many-body wave functions even when conventional observables or classical order parameters appear simple, and can expose transient regimes of high complexity during dynamical evolution that act as barriers to efficient simulation [2,25].

From this perspective, complexity measures provide a refined diagnostic of quantum correlations, capturing how information is distributed and processed across a system (see also holographic magic-harvesting approaches in Refs. [36,37]). This additional sensitivity makes them particularly well suited for probing far-from-equilibrium dynamics, strongly coupled systems, and emergent phenomena in quantum field theory.

Schwinger pair creation [38] in strong chromoelectric fields offers a natural setting to study this question. In this process, an external field pulls quark-antiquark pairs out of the vacuum, similar to the strong color fields generated in the early stages of high-energy heavy-ion collisions.

At strong coupling, the produced pair is described holographically by an open string whose end points accelerate apart, providing a controlled framework in which both the tunneling process and the subsequent real-time dynamics can be analyzed. The Schwinger pair provides a concrete realization of Einstein-Podolsky-Rosen (EPR) entanglement in gauge theory [39–41]. The produced quark

^{*}Contact author: segrie@uw.edu

and antiquark are entangled into a color singlet yet accelerate into causally disconnected Rindler wedges. Connected correlators between the pair are nonzero but cannot arise from interactions after nucleation, identifying the pair's entanglement as EPR-type. Hence, it is natural to ask what quantum correlations the Schwinger pair carries beyond its total entanglement.

This problem is closely related to string formation and breaking in confining gauge theories. In such systems, (color) charges are connected by flux tubes whose dynamics govern hadronization and particle production. Recent studies of lattice models in one and two spatial dimensions, respectively, using tensor-network and quantum simulation methods have shown that quantum complexity measures, including entanglement and magic, reveal nontrivial structure in the wave function during string formation and fragmentation [26,42–66]. These results suggest that particle production processes are accompanied not only by the generation of entanglement, but also by a reorganization of quantum correlations.

In this work, we investigate the emergence of nonlocal quantum correlations in the pair creation process by analyzing the entanglement structure of a fixed real-space subsystem containing one member of the pair (i.e., a half-space bipartition). We find that the Lorentzian structure reshapes the entanglement spectrum and leads to nonlocal magic for $d > 2$.

Nonlocal magic of the schwinger pair.

String world sheet as an AdS_2 black hole: We consider a quark-antiquark pair produced by the Schwinger mechanism in a large- N_c , large- λ gauge theory with a holographic dual. In this limit string breaking is suppressed, and the pair remains connected by an open string. Fundamental matter is introduced via a probe Dirichlet-brane at radial position $z = z_M$, on which the quark has mass [67–76]

$$M = \frac{\sqrt{\lambda}}{2\pi z_M}. \quad (1)$$

A constant background electric field E on the brane induces pair production. The holographic description of this process is a string world sheet in AdS_{d+1} governed by the Nambu-Goto action [70],

$$S_{\text{NG}} = \frac{\sqrt{\lambda}}{4\pi} \int d\tau \int d\sigma \frac{1}{z^2} (\partial X_\mu \bar{\partial} X_\mu + \partial z \bar{\partial} z) + i \oint A, \quad (2)$$

where $(\partial, \bar{\partial}) = (\partial_\sigma \pm i\partial_\tau)$ and Virasoro constraints are imposed. The boundary term $i \oint A$ captures the Lorentz force on the quark end points.

The Euclidean instanton describing the tunneling event has the quark worldline tracing a circular path of cyclotron

radius R_{cyc} on the Dirichlet-brane. The world sheet solution satisfying the equations of motion and Virasoro constraints was obtained in Ref. [70]; extremizing the on-shell action over R_{cyc} fixes the cyclotron radius through

$$R_{\text{cyc}}^2 + z_M^2 = \frac{1}{a^2}, \quad (3)$$

where $a \equiv E/M$. The instanton exists for $E \leq E_c$, with critical field $E_c = M/z_M$ corresponding to $R_{\text{cyc}} = 0$. Note that the pair appears spatially separated in real space at $t = 0$, i.e. it appears nonlocally in the Lorentzian spacetime.

The induced metric on the world sheet instanton [70] is

$$ds_W^2 = \frac{1}{\sinh^2(2\pi n_{\text{inst}}\sigma)} (d\tau^2 + d\sigma^2), \quad (4)$$

where n_{inst} is the instanton winding number. A single pair corresponds to $n_{\text{inst}} = 1$.

Under analytic continuation to Lorentzian signature, $\tau \rightarrow it$, this becomes

$$ds_W^2 = \frac{1}{\sinh^2(2\pi\sigma)} (-dt^2 + d\sigma^2). \quad (5)$$

Introducing the coordinate $w = \coth(2\pi\sigma)$, the metric takes the manifestly static form

$$ds_W^2 = -(w^2 - 1)dt^2 + \frac{dw^2}{w^2 - 1}, \quad (6)$$

which is an AdS_2 black hole with horizon at $w = 1$.

The Hartle-Hawking construction patches the Euclidean instanton to the Lorentzian solution at $t = 0$ hence producing the maximally extended two-sided geometry, whose Penrose diagram has two asymptotic boundaries (one near each string end point) connected through the horizon [73]. This world sheet wormhole is the holographic geometrization of EPR entanglement between quark and antiquark: whereas entanglement among $\mathcal{O}(N^2)$ degrees of freedom gives rise to spacetime wormholes, the $\mathcal{O}(\sqrt{\lambda})$ entanglement carried by the color-singlet pair is encoded on the string world sheet [73,77]. Note that the Hartle-Hawking construction yields a thermofield-double state of the form

$$|\Psi\rangle \sim \sum_{k,c,\bar{c}} e^{-\pi E_k/a} \delta^{c\bar{c}} |k, c\rangle_q \otimes |k, \bar{c}\rangle_{\bar{q}}, \quad (7)$$

where $|k, c\rangle_q$ are eigenstates of the Rindler Hamiltonian with energy E_k and color index c , and the inverse temperature is $\beta_U = 2\pi/a$. We emphasize, however, that the entanglement we analyze is not the entanglement of the color singlet structure Eq. (7), but rather the entanglement

of a spherical spatial region A containing one member of the pair with its complement $\bar{A}$ which is equivalent to the quark-antiquark (bipartite) entanglement.¹ The reduced density matrix ρ_A therefore encodes the interplay of the pair's color-singlet structure with the full QFT vacuum.

Casini-Huerta-Myers (CHM) map and topological black holes: To compute Rényi entropies for a spherical entangling region of radius R_A we employ the conformal mapping of Casini, Huerta, and Myers [79]. The Euclidean AdS_{d+1} in Poincaré coordinates,

$$ds^2 = \frac{L^2}{r^2} dr^2 + \frac{r^2}{L^2} (dt_E^2 + d\rho^2 + \rho^2 d\Omega_{d-2}^2), \quad (8)$$

can be mapped to [79,80]

$$ds^2 = L^2 \left[\frac{d\zeta^2}{f(\zeta)} + f(\zeta) d\tau^2 + \zeta^2 du^2 + \zeta^2 \sinh^2 u d\Omega_{d-2}^2 \right], \quad (9)$$

with $f(\zeta) = \zeta^2 - 1$ (see Supplemental Material [81] for details).

This coordinate system has two remarkable properties. First, in the field theory, the transformation maps the causal development of the spherical region A onto $S^1_\tau \times \mathbb{H}^{d-1}$, where $\mathbb{H}^{d-1}$ is the hyperbolic plane parametrized by (u, Ω_{d-2}) [79]. The reduced density matrix ρ_A thereby becomes a thermal density matrix on $\mathbb{H}^{d-1}$ at inverse temperature $\beta = 2\pi$. Second, the metric (9) has a Euclidean horizon at $\zeta_h = 1$, which corresponds in Poincaré coordinates to $r^2 + \rho^2 = R_A^2$. This is precisely the Ryu-Takayanagi surface of the spherical region A [82,83]. The entanglement entropy maps to the Bekenstein-Hawking entropy of this horizon. Since the transformation is a coordinate change, the geometry remains globally AdS_{d+1} and the horizon is of Rindler type [79,80].

The n th Rényi entropy is obtained by extending the periodicity to $\tau \sim \tau + 2\pi n$, which in the bulk corresponds to replacing $f(\zeta)$ by the topological black hole [80,84]

$$f_n(\zeta) = \zeta^2 - 1 - \frac{\zeta_h^d - \zeta_h^{d-2}}{\zeta^{d-2}}, \quad (10)$$

$$\zeta_h(n) = \frac{\sqrt{1 + n^2 d(d-2)} + 1}{nd}, \quad (11)$$

where smoothness of the Euclidean geometry with period $2\pi n$ fixes $\zeta_h(n)$. At $n = 1$ one recovers $\zeta_h = 1$ and

¹The position dependence of the entanglement entropy of the accelerated pair was studied in Ref. [78]. The symmetric cut used in this work is the unique bipartition (which coincides with the quark-antiquark bipartition) and the reduced state is thermal on $S^1 \times \mathbb{H}^{d-1}$.

$f_1 = \zeta^2 - 1$. Here d denotes the spacetime dimension of the boundary gauge theory ($d = 4$ for $\mathcal{N} = 4$ SYM dual to $\text{AdS}_5 \times S^5$), so the topological black hole lives in the $(d+1)$ -dimensional bulk. Although the string world sheet is two-dimensional, it is embedded in this $(d+1)$ -dimensional bulk and inherits its n -dependence from the ambient geometry.

The relation between this two-dimensional world sheet and the $(d+1)$ -dimensional bulk is the construction used by [71] for the entanglement entropy of a quark. The accelerated pair, equivalently a static quark at the center of the sphere A after the conformal map (see Supplemental Material [81]) maps under the CHM map to a Polyakov loop at the center $u = 0$ of $\mathbb{H}^{d-1}$, wrapping τ . Its holographic image is the probe string at $u = 0$ running from the brane to the horizon ζ_h , and at $n = 1$ the induced metric of this string (see Supplemental Material [81]) reduces to the AdS_2 black hole of Sec. II A, whose horizon is the RT surface ζ_h .

The same accelerated world sheet underlies both the EPR pair of Sec. II A and the pair produced here by tunneling; the instanton fixes only the separation R_{cyc} (3), which does not enter the n -dependence and hence, as we will see, affects neither S_{EE} nor C_E . The quark end point is fundamental matter on the brane, so the world sheet ends on the brane in the UV and caps off on the horizon ζ_h (the RT surface) in the IR, at fixed $u = 0$.

Rényi entropies from the probe action: We now connect the topological black hole backgrounds (10)–(11) to Rényi entropies, following the approach of Lewkowycz and Maldacena [85] in probe brane language [84,86]. This is motivated by the quantum information-theoretic connection between the refined Rényi entropy and nonlocal magic as follows (see [17] and Supplemental Material [81] for details). The capacity of entanglement, $C_E = -\partial_n \tilde{S}_n|_{n=1}$, equals the variance of the modular Hamiltonian spectrum, $C_E = \text{Var}_\rho(-\log \rho_A)$. By Lemma 1 of [17], C_E vanishes if and only if the entanglement spectrum is flat, which is equivalent to vanishing nonlocal magic (lower bounded by antiflatness)

$$C_E(\rho_A) = 0 \Leftrightarrow \mathcal{M}^{(\text{NL})}(\psi_{AB}) = 0. \quad (12)$$

Thus $C_E > 0$ is both necessary and sufficient for the presence of nonlocal magic. This relation is purely information-theoretic.

In QFT, the n th Rényi entropy $S_A^{(n)} = \frac{1}{1-n} \log \text{Tr} \rho_A^n$ may be computed via the replica trick: $\text{Tr} \rho_A^n$ equals the partition function on a manifold $\mathcal{M}_n$ with a conical deficit $2\pi/n$ at the entangling surface. Defining the generating functional $W[\mathcal{M}_n] = -\log Z[\mathcal{M}_n]$, the entanglement entropy is [87–89]

$$S_A = \lim_{n \rightarrow 1} (\partial_n - 1) W[\mathcal{M}_n]. \quad (13)$$

In holography, the n th Rényi entropy follows from $S_A^{(n)} = \frac{n}{n-1} [\hat{I}(n) - \hat{I}(1)]$, $\hat{I}(n)$ is the on-shell gravitational probe action $\hat{I}(n)$ [90]

$$\tilde{S}_n = n^2 \partial_n \hat{I}(n). \quad (14)$$

The derivative at $n = 1$ yields the capacity of entanglement (see Supplemental Material [81] for more details),

$$C_E = -\partial_n \tilde{S}_n|_{n=1}, \quad (15)$$

which through (12) provides a direct diagnostic for nonlocal magic in the state of the produced pair. We emphasize that the vanishing of boundary terms at all orders in the n -expansion ensures that Eq. (15) is obtained from Eq. (14) without additional contributions.

Nonlocal magic of the Schwinger pair: We now evaluate the refined Rényi entropy $\partial_n \tilde{S}_n|_{n=1}$ for the probe string.²

The on-shell Nambu-Goto action in the n th topological black hole background is (up to a constant; see Supplemental Material [81] for more details)

$$\hat{I}(n) = -\sqrt{\lambda} \zeta_h(n). \quad (16)$$

The refined Rényi entropy (14) is therefore

$$\tilde{S}_n = n^2 \partial_n \hat{I}(n) = -\sqrt{\lambda} n^2 \zeta'_h(n). \quad (17)$$

At $n = 1$ this gives the entanglement entropy (in agreement with Refs. [39,67,71,84,91]),

$$S_{EE} = \tilde{S}_1 = -\sqrt{\lambda} \zeta'_h(1) = \frac{\sqrt{\lambda}}{d-1}. \quad (18)$$

For $d = 4$ (i.e. $\mathcal{N} = 4$ SYM), this yields $S_{EE} = \sqrt{\lambda}/3$.

The capacity of entanglement $C_E = -\partial_n \tilde{S}_n|_{n=1}$ (15) requires the slope of the refined Rényi entropy,

$$\partial_n \tilde{S}_n|_{n=1} = -\sqrt{\lambda} [2\zeta'_h(1) + \zeta''_h(1)] = -\sqrt{\lambda} \frac{d-2}{(d-1)^3}. \quad (19)$$

and is thus

$$C_E = \frac{\sqrt{\lambda}(d-2)}{(d-1)^3}, \quad (20)$$

which for $d = 4$ gives $C_E = \frac{2\sqrt{\lambda}}{27}$.

²The dominant vacuum contribution to the entanglement entropy of the spherical region is of order N_c^2 and is independent of the produced pair. Our calculation isolates the $\mathcal{O}(\sqrt{\lambda})$ probe-string contribution associated with the heavy quark-antiquark pair. All entropies and Rényi derivatives below are therefore excess contributions relative to the vacuum, induced by the probe.

This is strictly positive for $d > 2$. By Lemma 1 of Ref. [17], a nonvanishing capacity of entanglement implies a nonflat entanglement spectrum, which in turn establishes the presence of nonlocal magic $\mathcal{M}^{(NL)}$ in the Schwinger pair production. Note that C_E depends only on the spacetime dimension d and the 't Hooft coupling λ , and hence reflects the intrinsic structure of the pair's quantum state in the gauge theory.

The physical picture is as follows. The world sheet wormhole encodes the entanglement between the quark and antiquark and the total entanglement entropy is determined by the horizon area (position) at $n = 1$. The structure of the entanglement spectrum, however, is encoded in the response of the horizon position to changes in n . This response is governed by the full n -dependent topological black hole family (10). For $d > 2$, the topological black hole acquires a nontrivial “charge” that depends on n and the geometry genuinely deforms under the replica parameter.

In $d = 2$, the topological black hole is three-dimensional (the BTZ black hole), and the n -deformed geometry is locally equivalent to the undeformed one. C_E and hence the antiflatness vanish. This does not exclude subleading $1/N$ or finite- λ corrections which could yield a finite C_E at subleading orders.

A possible field-theory interpretation of the vanishing of C_E in $d = 2$ is the following. The bipartition considered in this work is the standard spherical region of the boundary CFT. CHM maps the reduced density matrix ρ_A to a thermal density matrix on $S_\beta^1 \times H^{d-1}$, with $\beta = 2\pi n$. Accordingly, $\tilde{S}_n$ is the thermal entropy in this hyperbolic space, while C_E is the corresponding heat capacitylike response.

For $d = 2$ the causal diamond of the entangling region is conformally equivalent to a Rindler wedge with no transverse spatial directions. The corresponding modular Hamiltonian is the conformal image of a boost generator, and the CHM thermal problem lives on $S_\beta^1 \times H^1$. The Schwinger-pair contribution then shifts the entanglement entropy, but its leading contribution to the CHM free energy is linear in the modular temperature. Consequently, its second derivative with respect to the modular temperature vanishes: the pair contributes to S_{EE} but not to C_E . This should not be confused with the vacuum itself, whose capacity is nonzero.

For $d > 2$, the same modular problem maps to a thermal theory on $S_\beta^1 \times H^{d-1}$ with transverse hyperbolic directions. Hence, the Schwinger-pair contribution to the free energy is no longer linear in the modular temperature [which is the field-theory counterpart of the nonlinear replica dependence of $\zeta_h(n)$] and $C_E > 0$.

Note that the calculation implicitly contains the backreaction of the probe onto the spacetime which is necessary to compute any Rényi entropy to begin with [92]. However, by employing the CHM map for the spherical entangling surface we can evade constructing the backreacted geometry explicitly since all necessary information is encoded in

the free energy and can be extracted from there (see Supplemental Material [81] for details).

It is interesting to compare the ratio C_E/S_{EE} in the for the probe correction with the vacuum result of [93,94]. For the vacuum of a d -dimensional CFT with holographic dual restricted to Einstein gravity and to a fixed regularization scheme, $C_E/S_{EE} = 1$ for a spherical region. This result is tied to the conformal collider bounds [95,96]. For the Wilson insertion, we find the coupling independent ratio of vacuum subtracted capacity to vacuum subtracted entanglement entropy $\frac{C_E}{S_{EE}} = \frac{d-2}{(d-1)^2}$, which only depends on the spacetime dimension. For $d = 4$, $C_E/S_{EE} = 2/9$.

Discussion. We have shown that nonlocal magic is generated dynamically in Schwinger pair creation at strong coupling. The capacity of entanglement, $C_E = \frac{\sqrt{\lambda}(d-2)}{(d-1)^3}$, is strictly positive for $d > 2$, establishing that the produced quark-antiquark pair carries genuinely nonlocal quantum correlations that go beyond entanglement. The color-singlet pair, despite its structural similarity to a Bell state, does carry nonlocal magic.

Notably, this result is independent of the acceleration $a = E/M$ and hence of the Unruh temperature $T_U = a/2\pi$. The nonlocal magic is therefore not a thermal artifact of the Rindler horizon, but reflects the intrinsic structure of the pair's quantum state as dictated by the gauge interaction for a given spacetime dimension.

The ratio $C_E/S_{EE} = (d-2)/(d-1)^2$, is independent of the 't Hooft coupling and reduces to $2/9$ for $d = 4$. It characterizes how far the entanglement spectrum deviates from flatness relative to its total weight.

The CHM map for a spherical entangling region produces a topological black hole whose equation of state has a curvature correction proportional to $(d-2)$. A probe string in this background has an action linear in $\zeta_h(n)$, so the n -dependence of all Rényi quantities is governed entirely by $\zeta_h(n)$. The curvature correction renders $\zeta_h(n)$ nonlinear in $1/n$, producing $C_E \neq 0$. For a planar black hole $\zeta_h(n)$ is linear in $1/n$ and $C_E = 0$. The capacity of entanglement is therefore controlled by the spatial curvature of the black hole horizon, which in the CHM construction is the curvature of H^{d-1} . This is only nontrivial for $d > 2$.

A possible field theory interpretation is as follows. Under the CHM map, the reduced density matrix ρ_A of a spherical region is mapped to a thermal density matrix on $S_\beta^1 \times H^{d-1}$, with $\beta = 2\pi n$. In this representation, $\tilde{S}_n$ is the thermal entropy and C_E is the corresponding heat-capacitylike response to changing the modular temperature.

For $d = 2$, the CHM thermal problem lives on $S_\beta^1 \times H^1$. Equivalently, the modular Hamiltonian is the conformal image of a boost generator. The Schwinger-pair contribution to the CHM free energy is linear in the modular temperature, and therefore shifts the entanglement entropy without adding a heat-capacitylike response. Thus, at

leading order, the pair contributes to S_{EE} but not to C_E . This does not mean that the full two-dimensional interval spectrum is flat; in fact, the vacuum has nonzero capacity. For $d > 2$, the CHM thermal problem is instead on $S_\beta^1 \times H^{d-1}$, and the pair contribution to the free energy has nonlinear modular temperature dependence, producing a positive capacity of entanglement.

This observation has a broader significance. Recent work [17] has established that gravitational backreaction encodes nonlocal magic in holographic systems. Our result demonstrates that it can be recovered from the free energy in the probe limit. The probe calculation should not be interpreted as neglecting backreaction. The Wilson-line insertion changes the full gauge-theory state, and in the bulk this change is represented by the first-order gravitational response of the geometry. The simplification for spherical regions is that the CHM replica construction maps the calculation to the free energy (which is on shell). Thus constructing the gravitational response is unnecessary, not the backreaction itself. This opens the door to studying nonlocal magic across the wide landscape of probe-brane systems in holography, including, for example defects, impurities, and flavor sectors, where entanglement has been extensively studied but complexity measures remain largely unexplored.

Several extensions merit investigation. In heavy-ion collisions, the electric fields are not constant but pulsed and it would be interesting to study the modification of the nonlocal magic during Sauter pulses [68,97]. It would be intriguing to understand the nonlocal magic during black hole evaporation and in the context of quantum extremal surfaces [98–106]. For example, it would be interesting to answer how nonlocal quantum correlations change across the Page curve in black hole evaporation. Third, one should extend our calculation to all values of the coupling λ via localization [71,107,108] which would require the connected stress-tensor correlator $\langle T_{\tau\tau} T_{\tau\tau} \rangle_W^{\text{conn}}$ on $S^1 \times H^3$ in the presence of the Polyakov loop. Moreover, the authors of [109–111] investigated the entanglement of the flux tube in Yang-Mills theories on the lattice using, among other observables, Rényi entropies and it would be interesting to understand the connection to our results.

Finally, the dynamical transition from zero to nonzero nonlocal magic in pair creation provides a concrete, analytically controlled example of complexity generation in a nonequilibrium process, which may be amenable to study via quantum simulation of lattice gauge theories.

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