New Regularization-Renormalization Method in Quantum Electrodynamics and Qualitative Calculation on Lamb Shift

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abstract
A simple but effective method for regularization-renormalization (R-R) is proposed for handling the Feynman diagram integral (FDI) at one loop level in quantum electrodynamics (QED). The divergence is substituted by some constants to be fixed via experiments. So no counter term, no bare parameter and no arbitrary running mass scale is involved. Then the Lamb Shift in Hydrogen atom can be calculated qualitatively and simply as \( \Delta E(2S_{1/2}) - \Delta E(2P_{1/2}) = 996.7 \text{MHz} \) versus the experimental value 1057.85 MHz.

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In quantum field theory (QFT), when calculating the Feynman diagram integral (FDI) beyond the tree level, one soon encountered the divergence. After handling it by some regularization method, one further introduced counter terms for cancelling the divergence and established the relation between bare parameter and physical parameter. Many physicists are not satisfied with the above recipe. Can the divergence be avoided? What happens when one pushes the cutoff \( \Lambda \) in momentum integral into infinity? Is this implying the point-like feature in QFT model? Will we need a drastic change in the foundation of our theory? (see the final discussion in the excellent book by Sakurai \[1\].)

After learning the previous methods for regularization-renormalization (R-R)\[2\], beginning from the work by Yang and Ni \[3\] and further by Ni et al. \[4,5\], we proposed a new R-R method as follows. When encountering a superficially divergent FDI, we first differentiate it with respect to external momentum or mass parameter enough times until it becomes convergent. After performing integration with respect to internal momentum, we reintegrate it with respect to the parameter the same times to return to original FDI. Then instead of divergence, some arbitrary constants \( C_i (i = 1, 2, \cdots) \) appear in FDI, showing the lack of knowledge about the model at QFT level under consideration. So they should be fixed by experiment via suitable renormalization procedure. For illustration, let us consider the various FDI at one loop level in QED, using Bjorken-Drell metric \[6\].

1. Self-energy of electron in QED.

The FDI for self-energy of electron reads (see \[1,6,7,8\]) \((e < 0)\)

\[
-i\Sigma(p) = (-ie)^2 \int \frac{d^4k}{(2\pi)^4} \frac{g_{\mu\nu} \gamma^\mu}{ik^2} \frac{i}{k^\nu - m} \gamma^\nu
\]

\[
= -e^2 \int \frac{d^4k}{(2\pi)^4} \frac{N}{D}
\]

with

\[
1 \noalign{\hfill (1) \noalign{\vspace{8pt}}}
\]

\[
\frac{1}{D} = \frac{1}{k^2[(p-k)^2 - m^2]} = \int_0^1 \frac{dx}{[k^2 - 2p \cdot kx + (p^2 - m^2)x]^2}
\]

\[
N = g_{\mu\nu} \gamma^\mu (p^\nu - k^\nu) = -2(p^\nu - k^\nu) + 4m
\]

We first perform a shift in momentum integration: \( k \rightarrow K = k - xp, \) so that

\[
-i\Sigma(p) = -e^2 \int_0^1 dx [-2(1 - x) \cdot (p + 4m)] I
\]

and concentrate on the logarithmically divergent integral:

\[
I = \int \frac{d^4K}{(2\pi)^4} \frac{1}{[K^2 - M^2]^2}
\]

with

\[
M^2 = p^2 x^2 + (m^2 - p^2)x
\]

A differentiation with respect to \( M^2 \) is enough to get

\[
\frac{\partial I}{\partial M^2} = \frac{-i}{(4\pi)^2} \frac{1}{M^2}
\]

(see formula (C.2.5) in Ref.[9]). Thus

\[
I = \frac{-i}{(4\pi)^2} [\ln M^2 + C_1] = \frac{-i}{(4\pi)^2} \ln \frac{M^2}{\mu^2}
\]
carries an arbitrary constant \( C_1 = -\ln \mu_2^2 \). After integration with respect to the Feynman parameter \( x \), one obtains \( (\alpha \equiv e^2/4\pi) \)

\[
\Sigma(p) = A + B \not{p} \\
A = \frac{\alpha}{\pi} m^2 [2 - \ln \frac{m^2}{\mu_2^2} + \left( \frac{m^2 - p^2}{p^2} \right) \ln \left( \frac{m^2 - p^2}{m^2} \right)] \\
B = \frac{\alpha}{4\pi} \left( \ln \frac{m^2}{\mu_2^2} - 3 - \left( \frac{m^2 - p^2}{p^2} \right) [1 + \frac{m^2 + p^2}{p^2} \ln \left( \frac{m^2 - p^2}{m^2} \right)] \right)
\]

(7)

Using the chain approximation, one can derive the modification of electron propagator as:

\[
\frac{i}{\not{p} - m} \rightarrow \frac{i}{\not{p} - m} \frac{1}{1 - \frac{\Sigma(p)}{p^2 - m}} = \frac{iZ_2}{\not{p} - m_R}
\]

(8)

\[
Z_2 = (1 - B)^{-1} \simeq 1 + B \\
m_R = \frac{m + A}{1 - B} \simeq (m + A)(1 + B) \simeq m + \delta m
\]

(9)

(10)

\[
\delta m \simeq A + mB
\]

For a free electron, the mass shell condition \( p^2 = m^2 \) leads to

\[
\delta m = \frac{\alpha m}{4\pi} \left( 5 - 3 \ln \frac{m^2}{\mu_2^2} \right)
\]

We want the parameter \( m \) in the Lagrangian still being explained as the observed mass, i.e., \( m_R = m_{\text{obs}} = m \). So \( \delta m = 0 \) leads to \( \ln \frac{m_R^2}{\mu_2^2} = \frac{5}{3} \), which in turn fixes the renormalization factor for wave function

\[
Z_2 = 1 - \frac{\alpha}{3\pi}
\]

(11)

2. Photon self-energy—vacuum polarization.

\[
\Pi_{\mu\nu}(q) = -(-ie)^2 \text{Tr} \int \frac{d^4k}{(2\pi)^4} \gamma_\mu \frac{i}{k - \not{m}} \gamma_\nu \frac{i}{k - \not{q} - \not{m}}
\]

(12)

Introducing the Feynman parameter \( x \) as before and performing a shift in momentum integration: \( k \rightarrow K = k - xq \), we get

\[
\Pi_{\mu\nu}(q) = -4e^2 \int_0^1 dx (I_1 + I_2)
\]

(13)

where

\[
I_1 = \int \frac{d^4K}{(2\pi)^4} \frac{2K_\mu K_\nu - g_{\mu\nu}K^2}{(K^2 - M^2)^2}
\]

(14)

with

\[
M^2 = m^2 + q^2(x^2 - x)
\]

(15)

is quadratically divergent while

\[
I_2 = \int \frac{d^4K}{(2\pi)^4} \frac{(x^2 - x)(2g_{\mu\nu}q_\mu q_\nu - g_{\mu\nu}q^2) + m^2g_{\mu\nu}}{(K^2 - M^2)^2}
\]

(16)

is only logarithmically divergent like that in Eqs.(3-6). An elegant way for handling \( I_1 \) is modifying \( M^2 \) into

\[
M^2(\sigma) = m^2 + q^2(x^2 - x) + \sigma
\]

(17)
and differentiating $I_1$ with respect to $\sigma$ two times. After integration with respect to $K$, we reinitialize it with respect to $\sigma$ two times, arriving at the limit $\sigma \to 0$:

$$I_1 = \frac{i g_{\mu\nu}}{(4\pi)^2} \left\{ [m^2 + q^2(x^2 - x)] \ln \frac{m^2 + q^2(x^2 - x)}{\mu_3^2} + C_2 \right\}$$

(18)

with two arbitrary constants: $C_1 = -\ln \mu_3^2$ and $C_2$. Combining $I_1$ and $I_2$ together, we find

$$\Pi_{\mu\nu}(q) = \frac{8i e^2}{(4\pi)^2} (g_{\mu\nu} - g_{\mu\nu} q^2) \int_0^1 dx (x^2 - x) \ln \frac{m^2 + q^2(x^2 - x)}{\mu_3^2} - \frac{i 4e^2}{(4\pi)^2} g_{\mu\nu} C_2$$

(19)

The continuity equation of current induced in the vacuum polarization [1]

$$q^\mu \Pi_{\mu\nu}(q) = 0$$

(20)

is ensured by the factor $(g_{\mu\nu} - g_{\mu\nu} q^2)$. So we set $C_2 = 0$. Consider the scattering between two electrons via the exchange of a photon with momentum transfer $q \to 0$ [1]. Adding the contribution of $\Pi_{\mu\nu}(q)$ to tree diagram amounts to modify the charge square:

$$e^2 \to e_R^2 = Z_3 e^2$$

$$Z_3 = 1 + \frac{\alpha}{3\pi} (\ln \frac{m^2}{\mu_3^2} - \frac{q^2}{5m^2} + \cdots)$$

(21)

The choice of $\mu_3$ will be discussed later. The next term in expansion when $q \neq 0$ contributes a modification on Coulomb potential due to vacuum polarization (Uehling potential).

3. Vertex function in QED.

$$\Lambda_\mu(p', p) = (-ie)^2 \int \frac{d^4k}{(2\pi)^4} \frac{i}{k^2} \gamma^\nu \frac{i}{\not{p} - \not{k} - m} \gamma^\mu \frac{i}{\not{p} - \not{k} - m} \gamma^\nu$$

(22)

For simplicity, we consider electron being on the mass shell: $p^2 = p'^2 = m^2$, $p' - p = q$, $p \cdot q = -\frac{q^2}{2}$. Introducing the Feynman parameter $u = x + y$ and $v = x - y$, we perform a shift in momentum integration:

$$k \to K = k - (p + \frac{q}{2})u - \frac{q}{2}v$$

Thus

$$\Lambda_\mu = -ie^2 [I_3 \gamma_\mu + I_4]$$

(23)

$$I_3 = \int_0^1 du \int_{-u}^u dv \int \frac{d^4K}{(2\pi)^4} \frac{K^2}{(K^2 - M^2)^3}$$

(24)

$$M^2 = (m^2 - \frac{q^2}{4})u^2 + \frac{q^2}{4}v^2$$

(25)

$$I_4 = \int_0^1 du \int_{-u}^u dv \int \frac{d^4K}{(2\pi)^4} \frac{A_\mu}{(K^2 - M^2)^3}$$

$$A_\mu = (4 - 4u - 2u^2)m^2 \gamma_\mu + 2i(u^2 - u)mq^\nu \sigma_{\mu\nu}$$

$$- 2 - 2u + \frac{u^2}{2} - \frac{v^2}{2})q^2 \gamma_\mu - (2 + 2u)vmq_\mu$$

(27)

Set $K^2 = K^2 - M^2 + M^2$, then $I_3 = I_3' - \frac{i}{32\pi^4}$. $I_3'$ is only logarithmically divergent and can be treated as before to be

$$I_3 = \frac{-i}{(4\pi)^2} \int_0^1 du \int_{-u}^u dv \ln \frac{(m^2 - \frac{q^2}{4})u^2 + \frac{q^2}{4}v^2}{\mu_1^2}$$

(28)
with $\mu_1^2$ an arbitrary constant. Now $q^2 = -Q^2 < 0$ ($Q^2 > 0$)

$$I_3 = -i\frac{m^2}{(4\pi)^2} \{ \ln \frac{m^2}{\mu_1^2} - \frac{5}{2} + \frac{1}{w} F(w) \}$$  \hspace{1cm} (29)

$$F(w) = \ln \frac{1+w}{1-w}$$

$$w = \frac{1}{\sqrt{4m^2/Q^2} + 1}$$

On the other hand, though there is no ultra-violet divergence in $I_4$, it does have infrared divergence at $u \to 0$. For handling it, we introduce a lower cutoff $\eta$ in the integration with respect to $u$

$$I_4 = \frac{i}{2(4\pi)^2} \{ [4 \ln \eta + 5] \frac{4w}{Q^2} F(w)m^2 \gamma_\mu + \frac{i4w}{Q^2} F(w)mq^\nu \sigma_{\mu\nu}$$

$$+ 4(2 \ln \eta + \frac{7}{4}w F(w)\gamma_\mu + \frac{1}{w} F(w) - 2\gamma_\mu) \}$$  \hspace{1cm} (30)

Combining Eqs.(29) and (30) into Eq.(23), one arrives at

$$\Lambda_\mu(p', p) = -\frac{\alpha}{4\pi} \{ [\ln \frac{m^2}{\mu_1^2} - \frac{3}{2} + \frac{1}{2w} F(w)] \gamma_\mu - (4 \ln \eta + 5) \frac{2w}{Q^2} F(w)m^2 \gamma_\mu$$

$$- i\frac{2w}{Q^2} F(w)mq^\nu \sigma_{\mu\nu} - 2(2 \ln \eta + \frac{7}{4})w F(w)\gamma_\mu \}$$  \hspace{1cm} (31)

When ($Q^2 << m^2$), we get

$$\Lambda_\mu(p', p) = \frac{\alpha}{4\pi} \left( \frac{11}{2} - \ln \frac{m^2}{\mu_1^2} + 4 \ln \eta \right) \gamma_\mu + i\frac{\alpha}{4\pi} \frac{q^\nu}{m} \sigma_{\mu\nu} - \frac{\alpha}{4\pi} \left( \frac{1}{6} + \frac{4}{3} \ln \eta \right) \frac{q^2}{m^2} \gamma_\mu$$

It means the interaction of the electron with the external potential is modified

$$-e\gamma_\mu \to -e[\gamma_\mu + \Lambda_\mu(p', p)]$$  \hspace{1cm} (32)

Besides the important term $i\frac{\alpha}{4\pi} \frac{q^\nu}{m} \sigma_{\mu\nu}$ which emerges as the anomalous magnetic moment of electron, the charge modification here is expressed by a renormalization factor $Z_1$:

$$Z_1^{-1} = 1 + \frac{\alpha}{4\pi} \left( [2 - \ln \frac{m^2}{\mu_1^2} - \frac{1}{2w} F(w)] + (4 \ln \eta + 5) \frac{2wm^2}{Q^2} F(w) + (2 \ln \eta + \frac{7}{4})2w F(w) \right)$$  \hspace{1cm} (33)

The infrared term ($\sim \ln \eta$) is ascribed to the bremsstrahlung of soft photons [6,8] and can be taken care by KLN theorem [10]. We will fix $\mu_1$ and $\eta$ below.

4. Beta function at one loop level.

Adding all four FDI’s at one loop level to the tree diagram, we define the renormalized charge as usual[6-9]:

$$e_R = \frac{Z_2}{Z_1} Z_3^{1/2} e$$  \hspace{1cm} (34)

But the Ward-Takahashi Identity (WTI) implies that [6-8]

$$Z_1 = Z_2$$  \hspace{1cm} (35)

Therefore

$$\alpha_R \equiv \frac{e_R^2}{4\pi} = Z_3 \alpha$$  \hspace{1cm} (36)
Then set $p^2 = m^2$ in $Z_2$ and $Q^2 = 0$ in $Z_1$ with $\mu_1 = \mu_2$, yielding

\[
\ln \eta = -\frac{5}{8} \quad (37)
\]

For any value of $Q$, the renormalized charge reads from Eqs. (19-21):

\[
e_R(Q) = e\left\{1 + \frac{\alpha}{\pi} \int_0^1 dx [(x - x^2) \ln \frac{Q^2(x-x^2) + m^2}{\mu_3^2}]\right\} \quad (38)
\]

\[
e_R(Q) \sim e\{1 + \frac{\alpha}{2\pi} \left[ \frac{1}{3} \ln \frac{m^2}{\mu_3^2} + \frac{1}{15} \frac{Q^2}{m^2}\right]\}, \quad (Q^2 << m^2) \quad (39)
\]

The observed charge is defined at $Q^2 \to 0$ (Thomson scattering) limit:

\[
e_{\text{obs}} = e_R|_{Q=0} = e \quad (40)
\]

which dictates that

\[
\mu_3 = m \quad (41)
\]

We see that $e_R^2(Q)$ increases with $Q^2$. For discussing the running of $\alpha_R$ with $Q^2$, we define the Beta function:

\[
\beta(\alpha, Q) \equiv Q \frac{\partial}{\partial Q} \alpha_R(Q) \quad (42)
\]

From Eq.(38), one finds:

\[
\beta(\alpha, Q) = \frac{2\alpha^2}{3\pi} - \frac{4\alpha^2 m^2}{\pi Q^2} \left\{1 + \frac{2m^2}{\sqrt{Q^4 + 4Q^2m^2}} \ln \frac{\sqrt{Q^4 + 4Q^2m^2} - Q^2}{\sqrt{Q^4 + 4Q^2m^2} + Q^2}\right\} \quad (43)
\]

\[
\beta(\alpha, Q) \approx \frac{2\alpha^2}{3\pi} - \frac{4\alpha^2 m^2}{\pi Q^2}, \quad \left(\frac{Q^2}{4m^2} << 1\right) \quad (44)
\]

\[
\beta(\alpha, Q) \approx \frac{4m^2}{3\pi} - \frac{4\alpha^2 m^2}{\pi Q^2}, \quad \left(\frac{4m^2}{Q^2} << 1\right) \quad (45)
\]

which leads to the well known result $\beta(\alpha) = \frac{2\alpha^2}{3\pi}$ at one loop level at $Q^2 \to \infty$.

5. The renormalization group equation

The renormalization group equation (RGE) in QED is obtained by setting $Q \to \infty$ and $\alpha \to \alpha_R(Q)$ in the right hand side of Eq.(42),

\[
Q \frac{\partial}{\partial Q} \alpha_R = \frac{2\alpha_R^2}{3\pi} \quad (46)
\]

Then after integration, one yields

\[
\alpha_R(Q) = \frac{\alpha}{1 - \frac{2\alpha}{3\pi} \ln \frac{Q}{m}} \quad (47)
\]

where the renormalization is made as

\[
\alpha_R|_{Q=m} = \alpha \quad (48)
\]

Note that, however, the running physical parameter is momentum transfer $Q$, not $\mu$. All $\mu_i (i = 1, 2, 3)$ had been fixed.

A Landau pole emerges at high $Q$,

\[
Q_{\text{Landau}} \equiv \mu_{\text{Landau}} = m \exp \left(\frac{3\pi}{2\alpha}\right) \quad (49)
\]
6. The calculation of Lamb Shift.

The famous Lamb Shift in Hydrogen atom, i.e., the energy difference between \(2S_{1/2}\) and \(2P_{1/2}\) states\[11],

\[
[\Delta E(2S_{1/2}) - \Delta E(2P_{1/2})]_{exp} = 1057.845(9) MHz
\]  

(50)

can be calculated in our formalism qualitatively in an elegant manner. In Eqs. (7) – (10), we set

\[
\delta m = A + mB = -\frac{\alpha^2 m}{2n^2}
\]

(51)

and a deviation parameter \(\eta\) from the mass shell of free motion,

\[
p^2 = m^2(1 - \eta)
\]

(52)

This parameter is precisely the same \(\eta\) used in Eqs. (30) – (33) for avoiding the infrared divergence in vertex function. Then

\[
-\frac{\alpha^2 m}{2n^2} = \frac{\alpha m (-\eta + 2\eta \ln \eta)}{4\pi} \frac{1}{1 + \frac{\alpha}{3\pi}}
\]

(53)

For \(n = 2\) case,

\[
\eta = \eta_2 = 7.44489 \times 10^{-4}
\]

\[
\ln \eta_2 = -7.20281
\]

(54)

Now the \(q^2\) term in vertex function implies a modification on Coulomb potential,

\[
-\frac{e^2}{4\pi r} \rightarrow -\frac{e^2}{4\pi r} - \frac{\alpha^2}{m^2} \frac{1}{6} + \frac{4}{3} \ln \eta)\delta(\vec{r})
\]

(55)

Adding further the contribution from anomalous magnetic momentum in vertex function and the Uehling potential in Eq. (21), one finds the energy shift of \(nS_{1/2}\) state as,

\[
\Delta E(nS_{1/2}) = \frac{\alpha^3}{3\pi n^3} R_y \left( \frac{1}{20} + \frac{1}{\eta} \right)
\]

(56)

where \(R_y \equiv \frac{1}{2}\alpha^2 m\).

Then

\[
\Delta E(2S_{1/2}) = 979.73 MHz
\]

(57)

\[
[\Delta E(2S_{1/2}) - \Delta E(2P_{1/2})]_{Theor} = 996.69 MHz
\]

(58)

where \(2P_{1/2}\) state gets an extra downward shift due to the anomalous magnetic moment. In recent years, the 'absolute' Lamb Shift of \(1S_{1/2}\) state was measured as

\[
[\Delta E(1S_{1/2})]_{exp} = 8172.86(5) MHz
\]

(59)

In the above calculation, we get

\[
[\Delta E(1S_{1/2})]_{Theor} = 6076.79 MHz
\]

(60)

with

\[
\eta = \eta_1 = 3.77274 \times 10^{-3}
\]

(61)

\[
\ln \eta_1 = -5.57995
\]

(62)
But we should add extra contributions from the vacuum polarization with $\mu_3 = m$, but $m \to m - R_y/n^2$

$$[\Delta E^{VP}(1S_{1/2})] = \frac{2Z^4\alpha^3}{3\pi n^4} R_y|_{Z=1,n=1} = 271\, MHz$$  \hspace{1cm} (63)

and the finite radius of proton $r_p = 0.862 \times 10^{-13}\, cm$,

$$[\Delta E^p(1S_{1/2})] = \frac{4}{5} R_y \left(\frac{\alpha}{\alpha_0}\right)^2 = 0.7\, MHz$$  \hspace{1cm} (64)

Altogether, we obtain

$$[\Delta E(1S_{1/2})]_{Theor} \simeq 6349\, MHz$$  \hspace{1cm} (65)

In summary, some remarks are in order.

(a) Our R-R method is really very simple: (i) When encountering a superficially divergent FDI, we use Feynman parameter trick to combine the denominator $D$ into one factor. (ii) Perform a momentum shift from $k \to K$, so that $D \sim (K^2 - M^2)^n$; (iii) Take derivative of FDI with respect to $M^2$ to raise the value of $n$ such that the integral becomes convergent. (iv) After momentum integration, reintegrate it with respect to $M^2$. (v) Then some arbitrary constants ($\mu_i, C_2$) with $\eta$ emerge. They can only be fixed by experiments (the observed mass $m$ and charge $e$) or by some deep reason in theoretical consideration (like the continuity condition of current or WTI). (vi) Since all constants ($\mu_i, C_2, \eta$) are fixed at one loop level ($L=1$) with the meaning of $m$ and $e$ reconfirmed as that at tree level, all previous steps can be repeated at next loop expansion ($L=2$).

(b) In QED, like any renormalizable model in QFT, there will be no trouble in high $L$ calculations because the number of arbitrary constants corresponds to that of so-called primitive divergent integrals, whose number remains finite. On the other hand, in a nonrenormalizable model, there would be more and more arbitrary constants $C_i$'s in high $L$ calculation, showing that such kind of model is not well-defined at QFT level.

(c) The procedure of shifting momentum $k \to K$ is legal in our treatment because, (i) For a logarithmically divergent FDI, the momentum shift brings no change in value; (ii) For a linearly or quadratically divergent FDI, the shift does bring a change (surface integral term) in the value, but eventually it is absorbed into arbitrary constants.

(d) The reason why these constants appear in substitution of divergence was explained in detail in Ref.[4]. It reflects the fact that the world is infinite whereas our knowledge remains finite. In particular, the value of mass $m$ or charge $e$ is beyond the reach of perturbative QED. The calculation on self-energy $\Sigma$ as in this paper has nothing to do with the mass generation of electron except a finite and fixed renormalization on wave function: $Z_2 = 1 - \frac{\alpha}{\pi}$. A model how a fermion can acquire a mass together with the phase transition of vacuum (environment) which provides a second mass scale (the standard weight) was discussed in Ref.[12]. Besides, mass of electron can be modified via radiative correction on self-energy when it is moving inside a cavity (see Ref.[13]). The appearence of Landau pole is also inevitable, showing the boundary of QED theory(see [5]).

(e) No counter term and/or bare parameter is needed any more.

(f) There is also not any arbitrary running mass scale after renormalization. All arbitrary constants $\mu_i$ and $\eta$ should be fixed in renormalization, leaving only the physical running mass scale like external momentum $p$ or momentum transfer $Q$.

(g) It is interesting to see the qualitative calculation of Lamb Shift being so simple. The discrepancy between theoretical and experimental values reflects the fact that the covariant form of QFT is not suitable for dealing with the binding state problem.

(h) This R-R method was used to derive the Higgs mass, 138 GeV, in standard model [4-5]. It works also quite well in QCD, which will be discussed elsewhere.
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