Vacuum Muon Decay and Interaction with Laser Pulses

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Muons decay in vacuum mainly via the leptonic channel to an electron, a muon neutrino and an electron antineutrino. Previous investigations have concluded that muon decay can be significantly altered only in a strong electromagnetic field when the muonic strong-field parameter is of order unity, which is far beyond the reach of lab-based experiments at current and planned facilities. In this Letter, an alternative mechanism is presented in which a laser pulse affects the vacuum decay rate of a muon *outside* the pulse. Quantum interference between the muon decaying with or without interacting with the pulse generates fringes in the electron momentum spectra and can increase the muon lifetime by up to a factor 2. The required parameters to observe this effect are available in experiments today.

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The highest intensity of electromagnetic fields that can be produced in the lab has been increasing in recent years [1] and is set to increase still further with several multipetawatt lasers operating or in the process of being commissioned or constructed [2]. In anticipation of the extended science reach that these facilities will provide, many suggestions have been made about phenomena that may be studied in strong electromagnetic fields (for reviews, see, e.g., Refs. [3–8]). One aspect of these phenomena is nonperturbativity at small coupling: the fundamental electromagnetic coupling $\alpha \approx 1/137 \ll 1$ is enhanced by the electromagnetic field intensity to be of order unity or larger. In high-power laser labs, an effective coupling to electrons and positrons ξ_e of order unity, corresponding to an all-order interaction of background photons with electrons and positrons, can nowadays be routinely accessed [9].

Recently, it has been demonstrated in experiment how high power lasers can be used to study *electroweak* processes. For example, laser-wakefield accelerated electron beams have been collided with high-Z solid targets to produce muons in setups that are relatively compact compared to traditional muon sources [10–12]. Muon-antimuon pairs are thereby generated in the Coulomb field of the target nuclei via the (two-step) Bethe-Heitler mechanism of an electron emitting a real photon that decays, or via the (one-step) trident process of direct pair production from an electron [13], with some fraction also generated through the decay of charged pions. Charged kaons, along with muons and pions, can even be created by relatively weak lasers, for example in nuclear-induced processes on ultra-dense hydrogen H(0) [14].

Noting that the effective coupling ξ_μ of a muon of mass m_μ to the electromagnetic field of a laser pulse is $\xi_\mu = \delta\xi_e$ (with $\delta = m_e/m_\mu \approx 1/207$) and considering high power lasers coming online can potentially reach $\xi_e \sim O(1000)$, we see that measuring small-coupling nonperturbativity in muon-laser interactions with $\xi_\mu \sim O(1)$ may soon be within experimental reach. Motivated by these developments, we revisit the question of whether electroweak decays involving electromagnetic charges may be modified by intense laser pulses.

Previous work on electroweak decays in electromagnetic backgrounds has focused mainly on infinitely extended fields. In the seminal work by Nikishov and Ritus [15] the leptonic decay of pions was calculated and in [16] also the leptonic decay of muons and neutrino emission of electrons in a constant crossed field were studied in detail. Constant crossed fields are particularly relevant when $\xi \gg 1$ because in that regime, the "formation length" of the process is sufficiently short that it is a good approximation to assume the background is locally constant and crossed [17–20]. However, because this parameter regime is not likely to be accessible in experiment in the near future, it is not a regime of immediate interest. General arguments have also been made [21] for why the total probability of a charged particle decay cannot be modified by an intense laser pulse, but the calculation was performed in the quasiclassical limit and the analysis was again concerned with decay in the laser pulse itself. The situation was clearly formulated by Narozhny and Fedotov [22] that an electromagnetic background can only significantly modify the total probability of an electroweak decay if the effect is (i) classical,

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changing the trajectory and hence the time dilation of a decaying particle or (ii) quantum, such that the muon strong-field parameter, $\chi_\mu = e\hbar\sqrt{-(p\cdot F)^2}/m_\mu^3 c^4$ where m_μ and p are the muon mass and momentum, F is the field tensor, and e is the charge on a positron, must be of order unity. Since $\chi_\mu = \delta^3\chi_e$ and experiment can only currently reach $\chi_e \sim O(1)$, this would imply electroweak decays can only be influenced by an intense laser in current and near future experiments by a very small, likely undetectable amount. These arguments were demonstrated by explicit calculation for the case of a monochromatic wave background with $\xi_\mu \ll 1$ but arbitrary ξ_e by Dicus *et al.* [23,24].

In the current Letter we consider electroweak decay in a laser pulse of *finite* longitudinal extent. This crucial difference allows for the particle to decay before or after interacting with the laser pulse, thus providing two extra routes to decay that are absent when the laser field is infinitely extended, as in previous treatments. We will find these two decay routes interfere with fringes appearing in the emitted particle momentum spectra. This "which-way interference" of histories can occur for standard quantum electrodynamic processes in strong fields [25–30]; here, we will see it for a particle decay process. Focusing on the example of muon decay, we will find the total vacuum decay rate can be suppressed down to 50% of its usual value. The controlling parameter originates from the change in the *classical* position of the muon due to having interacted with the laser pulse, but the effect is clearly quantum in nature, arising from the interference of different decay pathways. This mechanism circumvents the restrictions that previous analyses have placed on manipulating electroweak decays with strong electromagnetic fields.

Outline—The most common decay of a muon, μ^- is $\mu^- \to e^- + \bar{\nu}_e + \nu_\mu$, where e^- is an electron, $\bar{\nu}_e$ an electron antineutrino and ν_μ a muon neutrino. At center of mass energies much lower than the W-boson mass, one can employ Fermi's effective four-fermion interaction (we set $\hbar = c = 1$ in the following). The vacuum term can be written

$$\mathsf{T}_{00} = \int \mathfrak{M}_{00} e^{iQ\cdot x} d^4x; \qquad \mathfrak{M}_{00} = \frac{G}{\sqrt{2}} J_{q,\ell,\sigma} J^{\sigma}_{p,k}, \quad (1)$$

where $Q = q + \ell + k - p$ is the total momentum change (p is the muon momentum, q is the electron momentum, k is the muon neutrino momentum, and ℓ is the electron antineutrino momentum), $G \approx (293\ \mathrm{GeV})^{-2}$ is the Fermi constant [31], and

$$J_{q,\ell,\sigma} = \frac{\bar{u}_q \gamma_\sigma (1-\gamma_5) v_\ell}{\sqrt{(2V)^2 q^0 \ell^0}}; \qquad J^{\sigma}_{p,k} = \frac{\bar{u}_k \gamma^\sigma (1-\gamma_5) u_p}{\sqrt{(2V)^2 p^0 k^0}},$$

represent currents, with subscripts on spinors labeling momentum.

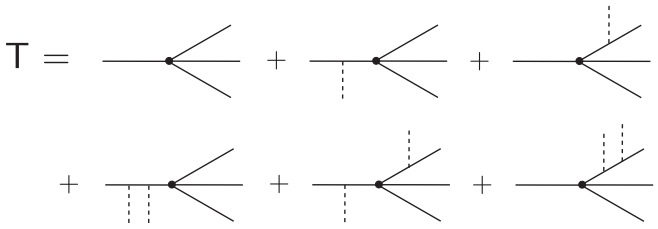


FIG. 1. First two orders of perturbative expansion in charge-field coupling ξ, $\mathsf{T} = \mathsf{T}_{00} + \mathsf{T}_{10} + \mathsf{T}_{01} + \mathsf{T}_{20} + \mathsf{T}_{11} + \mathsf{T}_{02}$, respectively.

Low intensity background (perturbative case)—It is instructive to begin with the case that the field intensity is small enough that $\xi \ll 1$, where ξ represents ξ_e and ξ_μ. Then the transition matrix T can be expanded in the charge-field interaction ξ as $\mathsf{T} = \sum_{ij} \mathsf{T}_{ij}$, where i (j) refer to the number of interactions between the plane wave background and the muon (electron) (see Fig. 1). The laser background is modeled as a plane wave of finite spatiotemporal extent. The scaled vector potential $a = eA$ (with $e > 0$ the charge on a positron) can be written $a(\phi) = m\xi\varepsilon g(\phi)$ where m is a mass, ξ an intensity parameter, ε a polarization vector and ϕ the phase where $\phi = \varkappa \cdot x$ with $\varkappa$ the wave vector. The squared intensity parameter, ξ_s^2 satisfies [32]: $\xi_s^2 = e^2\langle p\cdot T\cdot p\rangle_\phi/[m_s(\varkappa\cdot p)]^2$ where $s \in \{e,\mu\}$ refers to the particle species (electron or muon), p is the particle momentum, T is the energy momentum tensor, and $\langle\cdot\rangle_\phi$ refers to cycle averaging over the phase, ϕ. To represent a finite plane wave pulse, we choose $g(\phi)$ to be nonzero only when $0 < \phi < \Phi$ so that Φ denotes the pulse phase duration. (Zero-frequency components can be included in the pulse description, but as we will see, we are interested in channels that do not change the net particle momentum.) Vacuum decay kinematics are reflected by $\mathsf{T}_{00} \propto \delta^{(4)}(Q)$. When external-field photon interactions are added, we find each T_{ij} contains some contribution with the same kinematics as vacuum decay and some with different kinematics, allowing $Q \neq 0$. At the probability level, contributions with the same kinematics will interfere. The vacuum term scales as $\sim VT$ for typical measurement volume V and time T, whereas laser-only terms scale as $\sim V\tau$ where $\tau = \Phi/\varkappa^0$ is the laser pulse duration. Since $\tau \ll T$, only the vacuum terms are retained. Furthermore, the arguments cited in the introduction from previous work also imply direct modification of muon decay *inside* the pulse should be negligible in the perturbative regime.

Adding external-field photons and writing $\mathsf{T} = \mathsf{T}_{\mathsf{vac}}(\xi)$ where $\mathsf{T}_{\mathsf{vac}}$ indicates contributions with kinematics identical to those of vacuum decay. The vacuum term of the decay amplitude is

$$\mathsf{T}_{\mathsf{vac}}(0) = \mathsf{T}_{00} = (2\pi)^4 \delta^{(4)}(Q)\mathfrak{M}_{00},$$

and we find up to order ξ^2 in the charge-field interaction (see Supplemental Material, Sec. A [33] and Ref. [34]

therein for details):

$$\mathsf{T}_{\mathsf{vac}}(\xi) = \mathsf{T}_{\mathsf{vac}}(0)\left[1 + \frac{i\xi}{2}\mathcal{I}_1 + \frac{\xi^2}{2}\left(-\frac{1}{2}\mathcal{I}_1^2 + i\mathcal{I}_2\right)\right] \quad (2)$$

with the integrals

$$\mathcal{I}_1 = \mathcal{I}_{1q} - \mathcal{I}_{1p}; \quad \mathcal{I}_{1r} = m\frac{r\cdot\varepsilon}{\varkappa\cdot r}\int g(\phi)d\phi \equiv \int \mathcal{I}'_{1r}(\varphi)d\varphi$$
$$\mathcal{I}_2 = \mathcal{I}_{2p} - \mathcal{I}_{2q}; \quad \mathcal{I}_{2r} = \frac{m^2}{2\varkappa\cdot r}\int g^2(\phi)d\phi \equiv \int \mathcal{I}'_{2r}(\varphi)d\varphi, \quad (3)$$

where $r \in \{p, q\}$. Hence, a finite laser pulse can modify the rate of the vacuum decay channel. This occurs when the total momentum absorbed from the laser equals the momentum emitted back into the laser. Therefore, contributions from diagrams with an odd number of interactions with the laser such as T_{10} and T_{01} should be negligible (unless the laser has an exceptionally wide bandwidth or is heavily chirped). Indeed, these channels contribute to the vacuum decay amplitude with $\mathcal{I}_1$ and for a plane wave background with a finite number of cycles and a symmetric pulse envelope, each of the terms in $\mathcal{I}_1$ integrate to zero, i.e., $\mathcal{I}_1 = 0$. Instead, the main modification of the vacuum decay rate is from the ponderomotive term I_2. By considering each diagram in Fig. 1, we can infer where the muon decayed, with the situation illustrated in Fig. 2. The channel T_{20} must have occurred after the muon entered the laser pulse (possibly already having exited it), T_{02} must have occurred before the electron exited the laser pulse, and T_{11} must have occurred when both the muon and electron were in the laser pulse. Since T_{11} is proportional to the cross-term in $\mathcal{I}_1^2$, which is zero we conclude the contribution to vacuum muon decay from processes that can *only* occur *within* the laser pulse, is negligible. Instead, the laser contributes to vacuum decay by interacting with the muon or electron before or after the decay itself. Another way of seeing this is to consider the momentum change $Q(\varphi)$ for a decay in the plane wave pulse, by replacing the muon momentum p and electron momentum q with their plane-wave values. We then find $Q(\varphi) = Q + \Delta Q(\varphi)$ where

$$\Delta Q(\varphi) = \varkappa\left[\frac{2a\cdot q - a\cdot a}{2\varkappa\cdot q} - \frac{2a\cdot p - a\cdot a}{2\varkappa\cdot p}\right]. \quad (4)$$

Therefore $Q(\varphi) = 0$ at a finite and discrete set of points inside the pulse, where $p\cdot a(\varphi) = 0$ and $a(\varphi)\cdot a(\varphi) = 0$ are fulfilled, compared to outside the pulse, where $\Delta Q = 0$ everywhere. Indeed, the result in Eq. (2) is independent of the shape of the laser pulse and depends only on the square of the integral of its potential. Further analysis of where the process takes place, is given in Appendixes A and B.

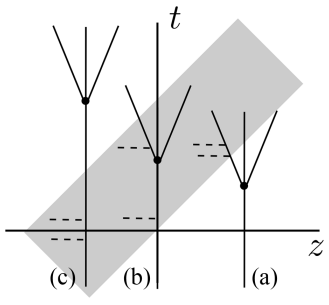


FIG. 2. Spacetime diagram showing various positions of muon decay, (a) before, (b) during, or (c) after laser pulse (gray shaded region), where dashed lines indicate interaction with the pulse.

High-intensity background—In an intense laser pulse, where $\xi \gtrsim 1$, both the muon and electron can become "dressed" in the pulse's electromagnetic field. This charge-laser coupling can be included to all orders of interaction by modeling the pulse as a plane wave and employing Volkov wave functions [3]. The free fermion wave functions then acquire electromagnetic field-dependent additions, for example

$$u_p e^{-ip\cdot x} \to \left[1 + \frac{\not{k}\not{a}(\phi)}{2\varkappa\cdot p}\right]u_p e^{-ip\cdot x + iS_{a,p}(\phi)},$$

i.e., a spinor-valued prefactor and a nonlinear phase given by

$$S_{a,p}(\phi) = \int_{-\infty}^{\phi}\left[-\xi\mathcal{I}'_{1p}(\varphi) + \xi^2\mathcal{I}'_{2p}(\varphi)\right]d\varphi, \quad (5)$$

where $\mathcal{I}'_{1p}(\varphi)$ and $\mathcal{I}'_{2p}(\varphi)$ are defined in Eq. (3). Motivated by the perturbative analysis, we can make a general argument for how a laser pulse of finite extent can modify the rate of vacuum decay of electroweak processes involving electromagnetically charged particles. The transition matrix can be written as

$$\mathsf{T}(\xi) = \int \mathfrak{M}[a]e^{iQ\cdot x + iS_a(\phi)}d^4x, \quad (6)$$

where we have separated out the combined nonlinear phase from the Volkov wave functions with

$$S_a(\phi) = S_{a,p}(\phi) - S_{a,q}(\phi) = \int_{-\infty}^{\phi}\xi\mathcal{I}'_1(\varphi) + \xi^2\mathcal{I}'_2(\varphi)d\varphi.$$

In a plane wave of finite extent $0 < \phi < \Phi$, $S_a(\phi) = 0$ before the pulse, i.e., when $\phi < 0$ and $S_a(\phi) = S_a(\Phi)$ after the pulse, when $\phi > \Phi$. Therefore, after the initial particle has propagated through the pulse the amplitude for decay acquires a constant phase. This contribution interferes with the purely vacuum decay contribution, to modify the total vacuum decay channel. (Although a monochromatic field would also modify the vacuum decay channel, since in that case the field is infinite, there is no interference with the purely vacuum contribution; this interference is the central effect we study here).

Separating the vacuum contributions into those originating before and after the leading edge of the pulse at $\phi = 0$, we find (see Supplemental Material, Sec. B [33]):

$$\mathsf{T}(\xi) = (2\pi)^4\delta^{(4)}(Q)\mathfrak{M}_{00}F(\xi) + 2(2\pi)^3\delta^{\perp,-}(Q)(\cdots), \quad (7)$$

where $\delta^{\perp,-}$ is a delta function in the three conserved momenta in a plane wave background. Using the same arguments as before, we retain only the vacuum channel and assume $\mathsf{T}(\xi) = \mathsf{T}_{\mathsf{vac}}(\xi)$, where $\mathsf{T}_{\mathsf{vac}}(\xi) = \mathsf{T}_{\mathsf{vac}}(0)F(\xi)$ with

$$F(\xi) = \frac{1}{2}\left[1 + e^{iS_a(\Phi)}\right]; \qquad F(0) = 1. \quad (8)$$

Expanding $F(\xi)$ to quadratic order in ξ, we find Eq. (7) tends to the direct, perturbative result from Eq. (2). To acquire the probability requires forming $|F(\xi)|^2$, and noting that $|F(\xi)|^2 = \cos^2[S_a(\Phi)/2] \to 1$ for $\xi \to 0$ but $|F(\xi)|^2 \to 1/2$ when the argument is averaged over, we see already at this stage, the origin of the 50% suppression in the decay rate that can take place. This is similar to the double-slit effect where the observer of the decay products lacks "which-way" information on whether the particle decayed with or without the laser pulse interaction. In this case the which-way information is not the trajectory of which slit is chosen, but rather which history of the particle led to its decay.

We proceed by calculating the decay rate $\mathsf{W}_{\mathsf{vac}}(\xi)$ (probability per unit time):

$$\mathsf{W}_{\mathsf{vac}}(\xi) = \frac{V^3}{T}\int \frac{d^3q d^3\ell d^3k}{(2\pi)^9}|\mathsf{T}_{\mathsf{vac}}(\xi)|^2. \quad (9)$$

The derivation proceeds along standard lines (see, e.g., Ref. [35]) in the muon rest frame. After the neutrino momenta are integrated over, the electron momentum integral in q is cast in spherical polar coordinates with the polar angle θ_q, coinciding with the projection on the laser wave vector, i.e., $\varkappa \cdot q = \varkappa^0(q^0 - |\mathbf{q}|\cos\theta_q)$. Performing the trivial integration over the azimuthal angle and setting the electron mass to zero without an appreciable change in the rate (see Supplemental Material, Sec. B [33] and Refs. [35,36] therein for details) leaves

$$\mathsf{W}_{\mathsf{vac}}(\xi) = \frac{G^2m_\mu^5}{48\pi^3}\int_0^{1/2} dZ \int_{-1}^{1} dX Z^2(3-4Z) \times \left\{1 + \cos\left[\Omega\left(1 - \frac{1}{Z(1-X)}\right)\right]\right\}, \quad (10)$$

where $Z = q^0/m_\mu$, $X = \cos\theta_q$, and $\Omega = \xi_\mu^2\Phi\langle g^2\rangle/2\eta_\mu$ ($\eta_\mu = \varkappa \cdot p/m_\mu^2$ is the muon energy parameter) and $\langle f\rangle = (1/\Phi)\int_0^\Phi f(\phi)d\phi$. The integration in Z and X can be performed analytically but the result is long and not particularly illuminating. In Fig. 3, we plot the electron polar distribution in the muon rest frame. In the absence of the laser, the emission is completely isotropic. As the parameter Ω is increased, fringes build up in the laser wave vector direction and the rate is decreased overall. The higher Ω is made, the more numerous the fringes become. As $\Omega \to \infty$, the emission becomes isotropic again, but with half the rate. In the energy distribution of emitted electrons in the muon rest frame Fig. 4, the suppression of the total rate can also be seen as $\Omega \to \infty$, although the appearance of fringes is much less pronounced, being clearest at the lowest energies. After performing all momentum integrations, we find

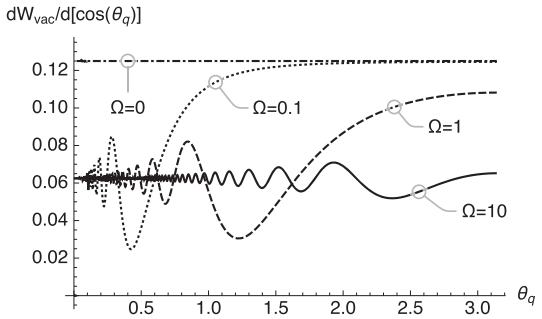


FIG. 3. Electron polar distribution in the rest frame of the muon $d\mathsf{W}_{\mathsf{vac}}/d[\cos(\theta_q)]$ for different values of the input parameter Ω.

$$\mathsf{W}_{\mathsf{vac}}(\xi) = \mathcal{R}[\Omega(\xi)]\mathsf{W}_{\mathsf{vac}}(0), \quad (11)$$

where

$$\mathcal{R}[\Omega] \approx \begin{cases} 1 - \frac{5\pi\Omega}{12} + \frac{\Omega^2}{18}(19 - 15\gamma - 15\log\Omega), & (\Omega \ll 1) \\ \frac{1}{2} - \frac{1}{\Omega^2} + \frac{20}{\Omega^4}, & (\Omega \gg 1) \end{cases}$$

with $\gamma = -\int_0^\infty e^{-x}\ln x dx \approx 0.577$ the Euler-Mascheroni constant. The decay rate in the absence of the laser is then the well-known result [35], $\mathsf{W}_{\mathsf{vac}}(0) = G^2m_\mu^5/192\pi^3$. [See Eq. (27) in Supplemental Material, Sec. B [33] for the explicit expression for $\mathcal{R}[\Omega]$.] We note from Fig. 5 the effect of laser pulse interference on muon decay is a suppression of the rate, up to around half the vacuum value. Although the effect is quantum mechanical in nature,

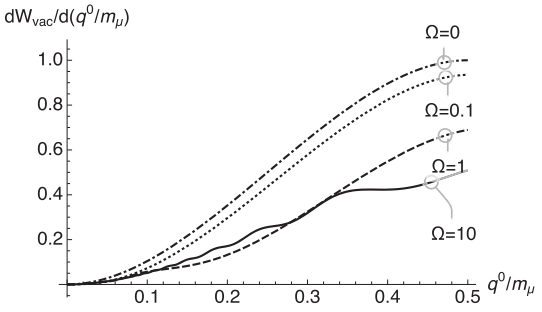


FIG. 4. Electron energy distribution in the rest frame of the muon $d\mathsf{W}_{\mathsf{vac}}/dq^0$ for different values of the input parameter Ω.

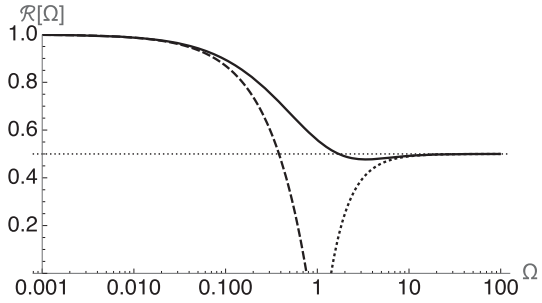


FIG. 5. A plot of the function $\mathcal{R}[\Omega]$ (solid line). The leading-order perturbative and asymptotic limits are indicated with dashed lines and the grid line is at $\mathcal{R}[\Omega] = 1/2$.

the parameter controlling pulse interference Ω is entirely classical. Because the Volkov wave function is semiclassical exact, the nonlinear phase term from the muon and electron, $S_a(\Phi)$, can be understood as originating from the net change in the position of the muon and electron due to acceleration in the laser field. Explicitly, $\Omega = p \cdot [x_f(\xi) - x_f(0)]$, where $x_f(\xi)$ is the muon position after interacting with the laser pulse of intensity parameter ξ. By artificially "turning off" the muon-laser or electron-laser coupling by setting ξ_μ or ξ_e to zero, we find that the dependence of the total rate most closely matches the electron-laser interaction (see Supplemental Material, Sec. B [33]). This suggests that the decay of other heavy particles to electrons or positrons may be affected by the same mechanism. In very intense fields, or fields with a sufficiently long duration, classical radiation reaction may significantly modify a charge's trajectory in a plane wave background. Here, the radiation reaction parameter for the muon or electron is $\nu = (2/3)\alpha\eta\xi^2\Phi$. If the field is very long, then eventually the probability P for the muon to decay, $\mathsf{P} = \mathsf{W}_{\mathsf{vac}}T$, will increase and to maintain unitarity, higher orders in the weak-field coupling G such as loops, must be included. On the other hand, if the field is so intense that $\chi_\mu \sim O(1)$, then we would have to include the nonvacuum part of the probability that changes the decay kinematics, which we have neglected for reasons outlined in the introduction. However, we will see that the laser pulse does not need to be long or intense to affect muon decay, and since $\eta\Phi \ll 1$ this means $\nu \ll 1$ and therefore radiation reaction and channels with modified kinematics should not appreciably modify the result. We also note that the controlling parameter Ω does not depend on the explicit pulse shape, but just as in the perturbative case, depends only on the integral of the square of the potential.

Recent experiments [12] colliding laser wakefield accelerated electrons with solid targets created muons with energies $\approx$300 MeV. If muons from such sources were collided with optical lasers, $\eta_\mu \sim O(10^{-8})$. Noting that $\langle g^2 \rangle \sim O(1)$ and the phase duration $\Phi = 2\pi N$ where $N \gg 1$ is the number of laser cycles, writing $\xi_\mu = \delta\xi_e$, gives $\Omega \sim \pi N\delta^2\xi_e^2/\eta_\mu$. Clearly, one can reach $\Omega \gg 1$ even with weak fields, for which $\xi_e \ll 1$, implying that the suppression of vacuum muon decay could be observable using laser parameters available in today's facilities. If a broadband muon source, such as bremsstrahlung, is to be used in experiment, then it would be important to understand how the bandwidth of the muon wave packet changes the laser's effect on vacuum decay. Using a toy model [36] of a Gaussian wave packet of muons colliding head-on with the laser pulse, we find that the laser's effect is only lightly suppressed, at mainly in the perturbative regime of $\Omega \ll 1$. At moderate Ω, for example $\Omega \gtrsim 2$, there is effectively no change even for a 100% bandwidth muon wave packet (see Supplemental Material, Sec. B [33] for details). It would also be important to understand how focusing effects, in particular the localized nature of the muon-laser interaction point, influence the main result. Any experimental test would need to be able to select for muon decays originating from trajectories that crossed the interaction point whilst the pulse was at the focus. However, since the laser intensity can be much weaker than the those used in all-optical muon sources [10–12], the laser beam can be defocused to provide a larger and more persistent target for the muons.

Throughout, we have assumed that neither the muon nor the electron radiate when interacting with the weak laser pulse. The probability of Compton scattering can be estimated using literature expressions for the perturbative limit $\xi \ll 1$ and $\eta \ll 1$, e.g., in a circularly polarized background plane wave of finite extent [37]. Then $\mathsf{P}_{e\to e+\gamma} \approx 2\alpha\xi_e^2\Phi\langle g^2\rangle/3$ and likewise for the muon with $e \to \mu$. Comparing probabilities, we find

$$\mathsf{P}_{\mathsf{vac}} \approx 1.5 \times 10^{-3}\mathcal{R}(\Omega)L[m]; \qquad \mathsf{P}_{e\to e+\gamma} \approx 10^{-4}\bar{N}\bar{\xi}_e^2,$$

where $L[m]$ is the distance in meters to the detector, $\Omega \sim 0.8\bar{N}\bar{\xi}_e^2/\bar{\eta}_\mu$ and where parameters have been scaled by typical experimental values $\bar{N} = N/10$, $\bar{\xi}_e = \xi_e/0.02$, $\bar{\eta}_\mu = \eta_\mu/10^{-8}$. For a high-power laser with wavelength 800 nm, the central frequency is 1.55 eV, in which case $\eta_\mu = 10^{-8}$ would correspond to muons with a kinetic energy of 40 MeV. An intensity parameter of $\xi_e = 0.02$ would correspond to an intensity of 1.6×10^{15} W cm^{-2} [38] and $N = 10$ laser cycles to a full-width-at-half-maximum pulse duration of 15 fs for a sine-squared pulse envelope. Therefore, if ξ_e is made small enough $\Omega \sim O(1)$ can still be achieved with electron Compton scattering much less probable than muon decay.

Conclusion—We have shown how muon decay can be significantly influenced by interaction with a laser pulse modeled as a plane wave of finite extent. Interactions with the electron and muon that involve zero *net* momentum change (for example absorbing and emitting the same number of laser photons) result in an interference with the standard vacuum decay channel. This interference of histories between the muon decaying with or without interacting with the laser pulse can modify the emitted

electron momentum spectrum and increase the muon lifetime by up to a factor 2. Because it depends on the muon being able to decay outside the laser pulse, this mechanism circumvents well-known limitations for manipulating electroweak processes with strong electromagnetic fields as it requires neither a large strong-field parameter nor a large intensity; indeed the effect can be demonstrated in weak fields. This Letter focused on muon decay but the same arguments clearly apply more generally to electroweak decays involving electromagnetic charges. All-optical experiments are using high-Z solid targets to create charged pions and muons in the lab [10–12]; future work could involve using a weaker secondary laser to investigate the effect on the number of electroweak decays via the mechanism outlined here. This could be supported by improved modeling that includes the localized nature of the interaction point in collisions with focused laser pulses.

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Data availability—No data were created or analyzed in this study.

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