

Control of tunneling ionization and subcycle interference dynamics with squeezed coherent light

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The influence of the quantum state of light on strong-field ionization has emerged as a frontier in attosecond physics. Here, we develop a quantum-optical extension of the quantum-trajectory Monte Carlo framework to study the ionization of hydrogen atoms driven by an orthogonally polarized two-color quantum field, composed of an 800 nm coherent field and a 1600 nm squeezed coherent field. We show that the quantum properties of light significantly influence the subcycle dynamics of photoelectron wave packets. We systematically analyze how the quantum statistics of light influence the photoelectron phase, momentum, yield, and number distribution. Our results reveal a direct correspondence between the quantum fluctuations of the driving field and the electron final state, unveiling an avenue for controlling electron wave packet interference via the quantum properties of light.

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I. INTRODUCTION

Ionization is one of the central topics in strong-field physics. The photoelectrons emitted during this process carry rich dynamical information [1–3] and encode structural characteristics of the parent atoms or molecules [4,5]. Consequently, strong-field ionization serves as the cornerstone for numerous significant applications, including molecular orbital imaging [6], photoelectron spin polarization studies [7,8], attosecond metrology [9], and high-order harmonic generation [10,11]. In most current studies within the strong-field community, the driving laser pulse is typically modeled as a classical field (i.e., a coherent state). With the rapid progress in generating ultrashort nonclassical light pulses [12–16], the intersection between quantum optics and strong-field physics has become an emerging frontier [17–21]. Fully quantum-mechanical treatments of the light-matter interaction are attracting increasing attention and have begun to uncover phenomena beyond the scope of classical models. Notably, recent studies [22,23] have explored how the intercycle interference in above-threshold ionization can be modulated by the quantum properties of light.

Temporally multicolor synthesized laser fields have been widely used in strong-field physics for processes such as single ionization [24–26], double ionization [27,28], and high-order harmonic generation [10,11]. Compared with monochromatic fields, multicolor fields introduce additional

degrees of freedom, including the intensity ratio [29], frequency ratio, and time delay [30] between the different color components. More importantly, each component in a multicolor field can possess its own polarization state, which exhibits distinct characteristics, enabling the exploration of various physical phenomena. For example, circularly polarized fields can be used to investigate circular dichroism [26], while orthogonally polarized two-color fields can be employed to study subcycle photoelectron interference [25]. Furthermore, given the current limitations in the intensity of available quantum light sources, employing a hybrid light configuration—comprising a strong coherent pulse and a weaker quantum field—greatly enhances experimental feasibility. The synthesized laser fields with quantum light might provide an avenue for imprinting the quantum statistical features of light onto HHG and photoelectrons.

In this work, we introduce a quantum-optical extension of the quantum-trajectory Monte Carlo (q-QTMC) framework to investigate how quantum light modulates subcycle photoelectron interference. The driving field combines a 1600 nm quantum field with an orthogonally polarized 800 nm coherent field, forming a quantum orthogonally polarized two-color (q-OTC) field. This configuration enables the resolution and control of ultrafast electron dynamics in two-dimensional momentum space. Unlike previous studies that primarily focused on modulating electron wave packets (WP) using classical light, our approach incorporates the quantum properties of light as an additional degree of freedom to control the mapping of electron quantum states onto momentum space. Specifically, we consider two types of quantum light, i.e., phase-squeezed coherent (PSC) and amplitude-squeezed coherent (ASC) states. We analyze the resulting photoelectron momentum distributions and systematically examine how quantum light affects the photoelectron phase, momentum, yield, and number distribution. Our findings reveal that the

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quantum statistical properties of photons can be imprinted onto photoelectrons, leaving distinct signatures in both momentum spectra and interference structures.

II. THEORETICAL MODEL

The evolution of the joint electron-field density matrix follows the Liouville-von Neumann equation

$$i\hbar \frac{\partial \rho_0(t)}{\partial t} = [\hat{H}_A + \hat{H}_F + \hat{\mathbf{d}} \cdot \hat{\mathbf{E}}, \rho_0(t)], \quad (1)$$

where $\hat{H}_A = \frac{\hat{\mathbf{p}}^2}{2m} + V(\hat{\mathbf{r}})$ is the Hamiltonian of a bound electron and $\hat{H}_F = \sum_{\mathbf{k}} \hbar\omega_{\mathbf{k}}(\hat{a}_{\mathbf{k}}^\dagger \hat{a}_{\mathbf{k}} + \frac{1}{2})$ is that of the quantized field. We study a hydrogen atom interacting with an orthogonally polarized two-color (OTC) field, composed of a 1600 nm squeezed field along the x axis and an 800 nm coherent field along the y axis. Denoting the 1600 and 800 nm fields as modes 1 and 2, the field Hamiltonian becomes $\hat{H}_F = \hbar\omega_1(\hat{a}_1^\dagger \hat{a}_1 + \frac{1}{2}) + \hbar\omega_2(\hat{a}_2^\dagger \hat{a}_2 + \frac{1}{2})$. The squeezed coherent state reads $|\alpha_{01}, r_1\rangle = \hat{D}(\alpha_{01})\hat{S}(r_1)|0\rangle$, with $\hat{D}(\alpha_{01})$ and $\hat{S}(r_1)$ the displacement and squeezing operators. The complex amplitude is $\alpha_{01} = \alpha_{01x} + i\alpha_{01y}$, and the squeezing parameter is $r_1 = |r_1|e^{i\theta_1}$. For simplicity, we set $\theta_1 = 0$. When $r_1 < 0$, the phase-space distribution is squeezed along the α_y direction (phase squeezing), and when $r_1 > 0$, the distribution is squeezed along the α_x direction (amplitude squeezing).

Transforming into the interaction picture, the evolution equation becomes

$$i\hbar \frac{\partial \rho_0(t)}{\partial t} = [\hat{H}_0(t), \rho_0(t)], \quad (2)$$

where $\hat{H}_0(t) = \hat{H}_A + \hat{\mathbf{d}} \cdot \hat{\mathbf{E}}(t)$. The initial density matrix is a tensor product of the atomic and field states

$$\begin{aligned} \rho_0(0) &= \rho_A(0) \otimes \rho_{F1}(0) \otimes \rho_{F2}(0), \\ \rho_A(0) &= |g\rangle\langle g|, \\ \rho_{F1}(0) &= \int d^2\alpha_1 d^2\beta_1 P(\alpha_1, \beta_1^*) |\alpha_1\rangle\langle\beta_1|, \\ \rho_{F2}(0) &= |\alpha_2\rangle\langle\alpha_2| = |\alpha_{02}\rangle\langle\alpha_{02}|. \end{aligned} \quad (3)$$

Here, $|g\rangle$ denotes the ground state of the atom, and $P(\alpha_1, \beta_1^*)$ is the Glauber P function, which is the quasiprobability distribution of a quantum state in the coherent state basis, widely used to describe nonclassical light. We substitute it into the Schrödinger equation and apply a unitary displacement transformation using the coherent displacement operator [31]. Note that the enormous photon number of the driving field renders electron backaction negligible, so that the entanglement between light and matter is ignored [32]. Tracing out the field degrees of freedom, we obtain the electronic density matrix:

$$\rho_e(t) = \int d^2\alpha_1 d^2\beta_1 P(\alpha_1, \beta_1^*) |\phi_{\alpha_1}(t)\rangle\langle\phi_{\beta_1}(t)|, \quad (4)$$

where $|\phi(t)\rangle$ denotes the wave function of the electron. The probability distribution of photoelectrons with momentum $\mathbf{p}$ is then

$$W(\mathbf{p}) = \int d^2\alpha_1 d^2\beta_1 P(\alpha_1, \beta_1^*) M_{\alpha_1}(\mathbf{p}) M_{\beta_1}^*(\mathbf{p}), \quad (5)$$

where $M_{\alpha_1/\beta_1}(\mathbf{p}) = \langle \mathbf{p} | \phi_{\alpha_1/\beta_1} \rangle$ represents the transition amplitude for photoelectrons with momentum $\mathbf{p}$. Furthermore, the Glauber P function $P(\alpha_1, \beta_1^*)$ can be defined through the Husimi Q function as

$$P(\alpha_1, \beta_1^*) = \frac{1}{4\pi} \exp\left(-\frac{|\alpha_1 - \beta_1^*|^2}{4}\right) Q\left(\frac{\alpha_1 + \beta_1^*}{2}\right). \quad (6)$$

Introducing $\varepsilon_{\alpha_1} = 2\epsilon_1^{(1)}\alpha_1$ and $\varepsilon_{\beta_1} = 2\epsilon_1^{(1)}\beta_1$, where $\epsilon_i^{(1)} = \sqrt{\frac{\hbar\omega_i}{2V\epsilon_0}}$ denotes the single-photon field amplitudes of mode i . We choose a typical mode volume for strong-field ionization, $V = 10^{-16} \text{ m}^3$, which leads to $\epsilon_1^{(1)} \rightarrow 0$, and employ the approximation

$$\lim_{\epsilon_1^{(1)} \rightarrow 0} \frac{1}{4\pi} \frac{1}{4|\epsilon_1^{(1)}|^2} \exp\left(-\frac{|\varepsilon_{\alpha_1} - \varepsilon_{\beta_1}|^2}{16|\epsilon_1^{(1)}|^2}\right) = \delta(\varepsilon_{\alpha_1} - \varepsilon_{\beta_1}), \quad (7)$$

the expression then simplifies to

$$W(\mathbf{p}) = \int d^2\alpha_1 Q(\alpha_1) |M_{\alpha_1}(\mathbf{p})|^2. \quad (8)$$

It is important to emphasize that this approximation does not imply the disappearance of quantum modulation of the driving field; the quantum modulation arising from the statistical distribution of the driving field's quantum state remains. To calculate the photoelectron momentum distribution driven by q-OTC fields, we developed the q-QTMC model by considering the quantum states of light. The QTMC theory [33] has been widely used to study the photoelectron momentum distribution in strong-field ionization driven by classical coherent light. Briefly, we sampled the initial transverse momentum and tunneling probability with the Ammosov-Delone-Krainov theory [34]. After tunneling, electron motion is computed via Newton's equation $\ddot{\mathbf{r}} = -\mathbf{r}/r^3 - \mathbf{E}(t)$ until the laser turns off. To consider the photoelectron interference, we consider the trajectory phase given by $S = \int_{t_0}^{\infty} [v^2(t)/2 - 1/r + I_p] dt$, t_0 is the ionization time, $v(t)$ is the electron velocity, and I_p is the ionization potential. For each field strength, we sample 2×10^8 trajectories. Including the quantum nature of squeezed coherent light, we then obtain the final photoelectron momentum distribution for q-OTC fields with Eq. (8).

III. RESULTS AND DISCUSSION

Figure 1(a) shows the calculated photoelectron momentum distribution driven by an OTC field composed of coherent state pulses at 800 nm ($E_2 = 0.063 \text{ a.u.}$, $\theta_2 = -\frac{\pi}{2}$) and 1600 nm ($E_1 = 0.020 \text{ a.u.}$, $\theta_1 = 0$, $r_1 = 0$). We refer to this configuration as the coherent state OTC field. In such OTC fields, subcycle photoelectron wave packets emitted at different times interfere coherently, forming distinct momentum-space interference patterns, including stripelike and carpetlike structures [24,35,36].

It is important to study quantum-optical effects on photoelectron interference when one component of the OTC field is replaced by quantum light. In further study, the 1600 nm component adopts either a phase- or amplitude-squeezed coherent state, with $E_1 = 0.020 \text{ a.u.}$ and $\theta_1 = 0$. These configurations

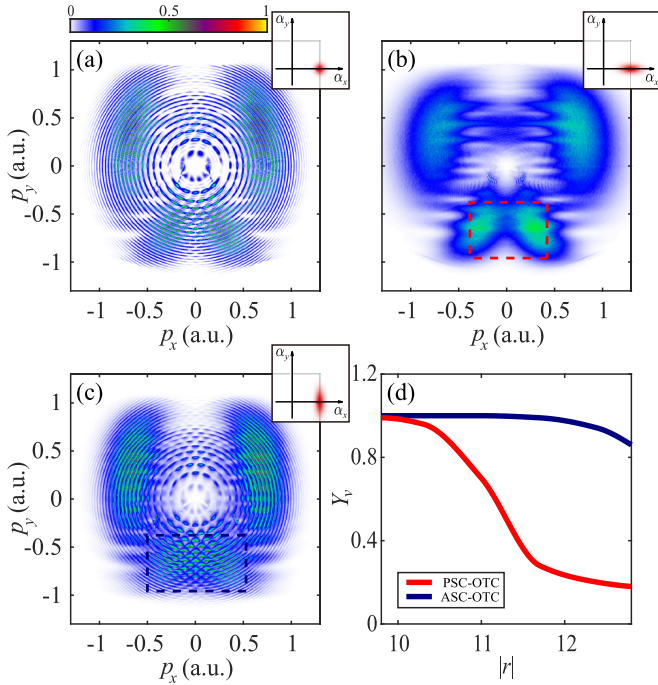


FIG. 1. (a)–(c) Photoelectron momentum distributions under the coherent state, PSC- and ASC-OTC fields. Insets in panels (a)–(c) show schematic Husimi Q functions of the 1600 nm components. Dashed boxes in panels (b) and (c) highlight regions exhibiting the carpetlike interference patterns. (d) Decoherence parameter Y_V vs squeezing degree $|r|$ for photoelectrons in $p_x \in [-0.05, 0.05]$, $p_y \in [-0.65, -0.55]$.

are termed PSC-OTC and ASC-OTC fields. The corresponding squeezing parameters for the 1600 nm quantum light are set to $r = -12.8$ and $r = 12.8$. In Figs. 1(b) and 1(c), we present the photoelectron momentum distributions in the PSC-OTC and ASC-OTC fields, respectively. Obviously, the quantum properties of the driving field strongly modulate the photoelectron momentum distributions. Both the PSC-OTC and ASC-OTC fields exhibit pronounced decoherence in the momentum spectra. This decoherence reflects a transition in the behavior of photoelectrons from wavelike to particlelike, implying that the quantum properties of light offer an effective knob for controlling this wave-particle duality. To examine subcycle photoelectron interference in detail, we focus on the carpetlike pattern as a representative feature. As marked in Figs. 1(b) and 1(c), the ASC-OTC case gives rise to a pronounced and widely distributed carpetlike structure. In contrast, the PSC-OTC case exhibits a significantly reduced overlap region between the left and right branches, and even within this limited area, the carpetlike structure is largely suppressed.

To quantify decoherence in the carpetlike interference, we introduce a parameter based on photoelectron yields:

$$Y_V = \frac{Y_{\max} - Y_{\min}}{Y_{\max} + Y_{\min}}, \quad (9)$$

where $Y_{\max}$ and $Y_{\min}$ denote the maximum and minimum photoelectron yields within $p_x \in [-0.05, 0.05]$ and $p_y \in [-0.65, -0.55]$. The values of Y_V are 0.18 for the PSC-OTC

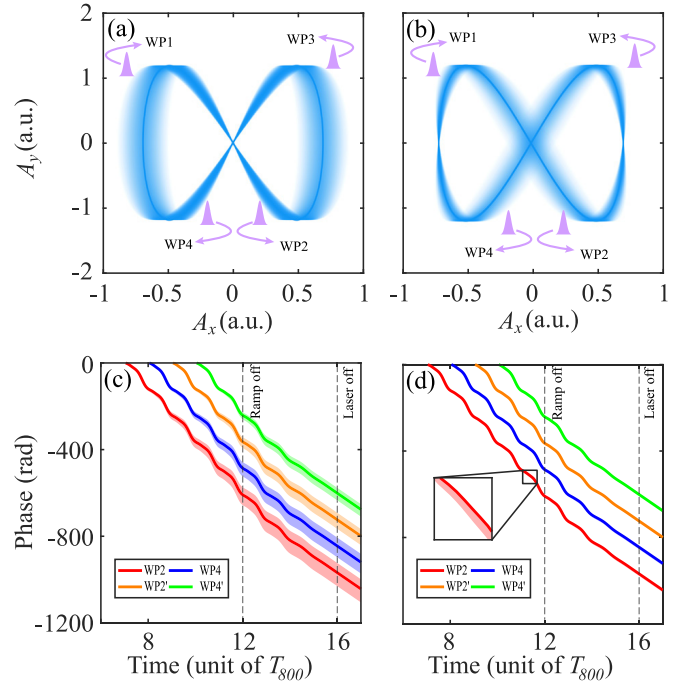


FIG. 2. Negative vector potentials of PSC-OTC [panel (a)] and ASC-OTC [panel (b)] fields, with tunneling photoelectron wave packets (WP1–WP4) labeled. (c), (d) Time evolution of the phases of WP2, WP4, WP2', and WP4' in the PSC-OTC and ASC-OTC fields, respectively.

case and 0.86 for the ASC-OTC case. Then, we compute the dependence of Y_V on the squeezing degree $|r|$, as shown in Fig. 1(d). It is evident from the figure that the photoelectron momentum distribution is highly sensitive to the value of r . The results for positive and negative r exhibit distinct differences, which become increasingly pronounced as $|r|$ grows. This indicates that the quantum properties of the driving light have a strong influence on the electron momentum distribution.

The quantum properties of the driving field influence the phases of the tunneling electron wave packets. Figures 2(a) and 2(b) show schematic negative vector potential distributions for PSC-OTC and ASC-OTC fields. The central blue curves represent the vector potential of the coherent state OTC field, while shaded areas indicate the quantum uncertainties of the squeezed fields. In each cycle, four electron wave packets (WP1–WP4) are emitted, as shown in Figs. 2(a) and 2(b). After tunneling, the electron wave packets evolve under the combined influence of the oscillating laser field and the Coulomb potential, accumulating phase over time. This accumulated phase exhibits an uncertainty under quantum light, defined as $\Delta S = \max(S_{\mathbf{E}_{\alpha_1}}(\mathbf{p}, t)) - \min(S_{\mathbf{E}_{\alpha_1}}(\mathbf{p}, t))$. Here, $\mathbf{E}_{\alpha_1}$ is sampled from the interval $[\mathbf{E}_1 - \mathbf{E}_0, \mathbf{E}_1 + \mathbf{E}_0]$, with $E_0 = 0.006$ a.u., the standard deviation of vacuum fluctuations. As shown in Figs. 2(a) and 2(b), the carpetlike interference predominantly results from the interference between WP2 and WP4. Therefore, we analyze the influence of quantum light on the phase evolution of different electron wave packets by taking WP2, WP4, WP2', and WP4' as representative examples, where WP2' and WP4' refer to

electron wave packets ionized in the subsequent cycle relative to WP2 and WP4.

Figures 2(c) and 2(d) illustrate the time evolution of the photoelectron wave packet phase modulated by the PSC-OTC and ASC-OTC fields, respectively. Solid curves represent wave packet phases in the coherent state OTC field; shaded areas show the corresponding uncertainties. In the coherent state OTC field, each wave packet carries a well-defined phase, while under quantum light, phase uncertainty accumulates from ionization to the end of the pulse. Notably, the electron wave packets emitted earlier accumulate a larger amount of phase uncertainty by the time the pulse terminates. Furthermore, the phase uncertainty is significantly larger in the PSC-OTC case than in the ASC-OTC case. This enhanced phase uncertainty results in stronger decoherence and thus a more pronounced suppression of interference fringes. These findings underscore the critical role of phase uncertainty, driven by the quantum properties of light, in shaping the photoelectron momentum distribution and controlling the emergence or suppression of the carpetlike interference structures.

Apart from the phase, quantum statistics also crucially affect the momentum-space distribution of photoelectron wave packets. By varying the ionization times, the momentum-space distributions of individual wave packets can be isolated and analyzed. Due to the symmetry of the laser field configuration, WP3 and WP4 exhibit similar momentum-space characteristics to WP1 and WP2. Therefore, we focus on WP1 and WP2 for further analysis. Figures 3(a) and 3(b) show the momentum-space distributions of WP1 and WP2 in the PSC-OTC field. The long-range Coulomb potential pulls more electrons back toward the ionic core, suppressing the direct ionization trajectories while enhancing rescattered ones. Similarly, the momentum-space distributions of WP1 and WP2 in the ASC-OTC field are simulated and compared with those in the PSC-OTC field. For the p_x distribution, quantum light induces noticeable modulation on the photoelectron momentum. Figures 3(c) and 3(d) show the p_x distributions of WP1 and WP2, with PSC-OTC curves shifted for clarity. It is evident that the p_x distribution of WP1 ionized by the PSC-OTC field is significantly broader than that in the ASC-OTC field, while the distribution of WP2 is narrower in the PSC-OTC field compared to the ASC-OTC field. These results demonstrate that the quantum properties of the field directly influence the final momentum distribution of the photoelectron wave packets. Specifically, the final momentum is given by $p_x = p_{0x} - A_x(t_0)$, where p_{0x} denotes the initial momentum. The photoelectrons experience strong modulation from $-A_x(t_0)$. In different q-OTC fields, the uncertainty in the negative vector potential $-A_x(t_0)$ varies. In the PSC-OTC field, the negative vector potential at the time of WP1 ionization exhibits greater uncertainty, leading to a broader p_x distribution for WP1. Similarly, in the ASC-OTC field, the negative vector potential at the time of WP2 ionization exhibits greater uncertainty, resulting in a broader p_x distribution for WP2.

For the p_y distribution, the modulation by the quantum light is relatively limited. As an example, we analyze the p_y distribution of WP1 in the q-OTC field, as depicted in Fig. 3(e). The ASC-OTC distribution is slightly broader, though the difference is minimal. The photoelectron wave

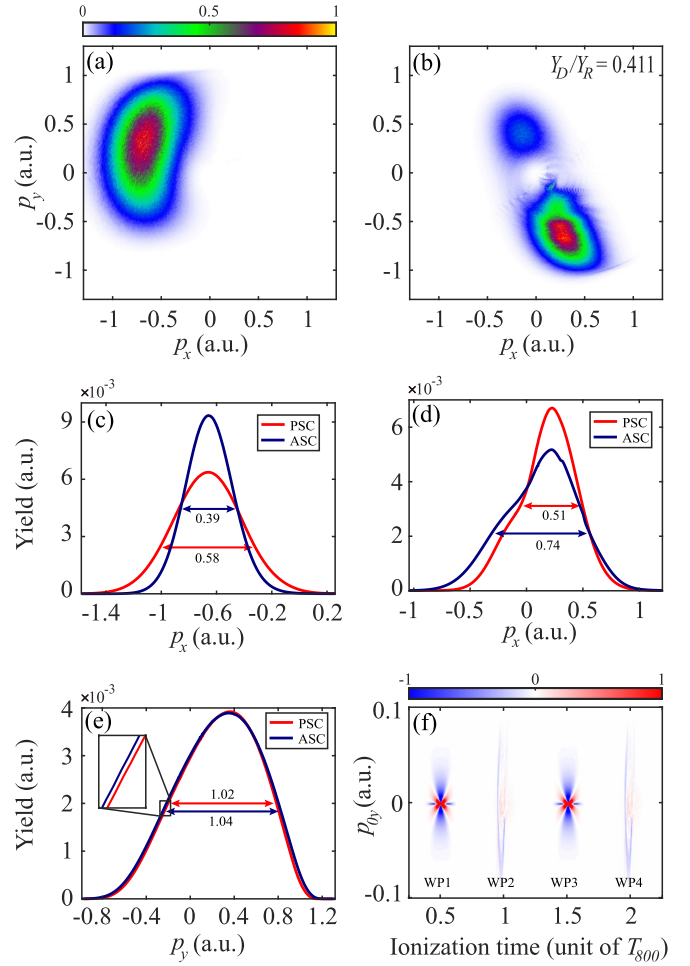


FIG. 3. Momentum distributions of WP1 [panel (a)] and WP2 [panel (b)] in the PSC-OTC field. (c), (d) p_x distributions of WP1 and WP2 in the two types of q-OTC fields. (e) p_y distribution of WP1 in the two types of q-OTC fields. (f) Differential photoelectron yield between PSC-OTC and ASC-OTC fields as a function of t_0 and p_{0y} . The full widths at half maximum of the curves in panels (c)–(e) are indicated in the corresponding panels. PSC and ASC refer to the PSC-OTC and ASC-OTC fields; the same notation is used throughout the remaining figures.

packets satisfy $p_y = p_{0y} - A_y(t_0)$, where p_{0y} denotes the initial momentum. In both q-OTC fields, $-A_y(t)$ is identical, and the ionization time t_0 is predominantly determined by the 800 nm coherent field, which is essentially the same in both cases. Therefore, the quantum field primarily modulates p_{0y} . To examine this effect, we trace back the p_{0y} of electrons ionized at different times and calculate the difference in ionization yields between the PSC-OTC and ASC-OTC fields, as shown in Fig. 3(f). In ASC-OTC, photoelectrons typically acquire larger p_{0y} , resulting in broader final-state p_y distributions.

The carpetlike interference in the q-OTC field mainly arises from WP2–WP4 interference. In the PSC-OTC field, their momentum distributions are significantly narrower than in the ASC-OTC field [Fig. 3(d)]. The broader overlap in ASC-OTC leads to a wider interference region. Besides phase uncertainty, this momentum-space distribution also controls

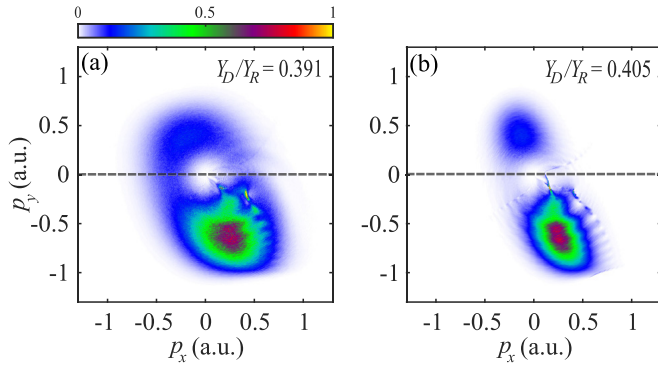


FIG. 4. Momentum distributions of WP2 in the ASC-OTC [panel (a)] and coherent state OTC [panel (b)] fields. The dashed line separates the main momentum regions corresponding to directly ionized and rescattered photoelectrons. Yield ratios Y_D/Y_R are indicated in the upper-right corner of each panel.

the emergence and suppression of the carpetlike structure. Furthermore, we analyze how q-OTC fields modulate the momentum distribution of photoelectron wave packets. The quantum light influences both the initial and final states of the photoelectron. The p_x broadening of WP1 is more sensitive to amplitude uncertainty, while that of WP2 is more sensitive to phase uncertainty. The p_x distributions of photoelectrons ionized at different times may help characterize the driving field's amplitude and phase uncertainties.

In addition to the phase and momentum-space distributions, the photoelectron yield also plays a crucial role in shaping the electron wave packets. The direct and rescattered photoelectrons in WP2 exhibit a clear separation in momentum, allowing us to compare the yields of these two contributions. In PSC-OTC [Fig. 3(b)], the direct-to-rescattered yield ratio is $Y_D/Y_R \sim 0.411$. Similarly, in the ASC-OTC field [Fig. 4(a)], the yield ratio is $Y_D/Y_R \sim 0.391$. For comparison, we also show the WP2 distribution in the coherent state OTC field in Fig. 4(b), which gives $Y_D/Y_R \sim 0.405$. Compared to the coherent state OTC field, the larger ratio in the PSC-OTC field is attributed to amplitude uncertainty. Due to the focusing effect, the dominant contributions to the PSC-OTC field originate from regions with stronger electric fields, thereby enhancing the proportion of directly ionized electrons. The smaller ratio in the ASC-OTC field compared to the coherent state OTC field is attributed to phase uncertainty. When WP2 is ionized in the ASC-OTC field, the negative vector potential $-A_x(t)$ exhibits larger fluctuations near zero. This gives rescattered electrons higher transverse momentum, reducing recapture and boosting rescattering yield. Since the carpetlike interference pattern is primarily formed by the interference of rescattered electrons, the increased rescattering yield in the ASC-OTC field is one of the reasons why the interference structure becomes more pronounced.

We have further calculated the photon and corresponding photoelectron number distributions. The total photon number $N = N_{\text{ph}} + N_{\text{coherent}}$ includes contributions from the squeezed coherent field (N_{ph}) and the coherent field (N_{coherent}), with N_{coherent} constant due to the fixed intensity. We therefore focus on N_{ph} . Given the known E_{α_1} distribution, the

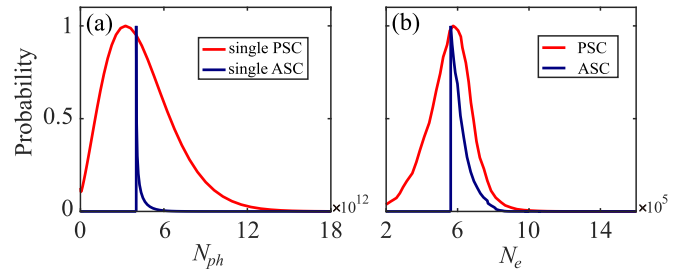


FIG. 5. Normalized photon number distribution of the quantum light [panel (a)] and the corresponding normalized photoelectron number distribution driven by the q-OTC field [panel (b)].

photon number distribution $P_{\text{ph}}(N_{\text{ph}})$ is obtained via $N_{\text{ph}} = KE_{\alpha_1}^2$, where K is determined by laser parameters via $N_{\text{ph}} = I_{\alpha_1} \tau A / (\hbar \omega)$, with τ denoting the pulse duration and A the focal area of the beam. Here, we set $K = 10^{16}$. The resulting photon number distribution is shown in Fig. 5(a). The photoelectron number distribution $P_e(N_e)$ is obtained through $P_e(N_e) = P_{\text{ph}}(N_{\text{ph}}) \frac{dN_{\text{ph}}}{dN_e}$, which follows directly from probability conservation under a change of variables $P_e(N_e) dN_e = P_{\text{ph}}(N_{\text{ph}}) dN_{\text{ph}}$. Here, N_e represents the relative yield of the photoelectrons. Based on $P_{\text{ph}}(N_{\text{ph}})$, we numerically evaluate dN_{ph}/dN_e , and subsequently obtain $P_e(N_e)$, as shown in Fig. 5(b). It clearly reveals that the quantum statistical properties of the field are projected onto the photoelectron number distributions. Moreover, a broader photon number distribution causes decoherence, effectively blurring the interference structures in momentum space.

IV. CONCLUSION

We have studied strong-field ionization by a q-OTC field, focusing particularly on the emergence and suppression of subcycle carpetlike interference structures. The quantum field's statistical properties strongly affect the subcycle electron dynamics. Quantum light imprints distinct signatures on photoelectron phase, momentum, yield, and number distributions. Given the limited intensity of available squeezed coherent light sources, such a hybrid light configuration—comprising a strong coherent pulse and a weaker quantum field—offers a promising route for strong-field ionization. Controlling wave packet interference with quantum light is also crucial for the development of interferometric techniques to probe ultrafast multielectron dynamics and steer chemical reactions [37]. These findings open new pathways for manipulating ultrafast electron dynamics [38] by using structured quantum light.

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DATA AVAILABILITY

No data was created in this article.

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