Theory of laser-assisted nuclear fusion

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Date: 2025-06-02T18:57:47+00:00

Abstract

The process of nuclear fusion in the presence of a laser field is analyzed theoretically. The analysis is applicable to most fusion reactions and different kinds of currently available intense lasers, from X-ray free electron lasers to solid-state near-infrared lasers. Laser fields are shown to enhance the fusion yields and the mechanism of the enhancement is explained. Low-frequency lasers are shown to be more efficient in enhancing fusion than high-frequency lasers. Calculation results show enhancements of fusion yields by orders of magnitude with currently available intense low-frequency laser fields. The temperature requirement of controlled nuclear fusion may be reduced with the aid of intense laser fields.

Full Text

Preamble

Theory of laser-assisted nuclear fusion∗ Jintao Qi,1 Zhaoyan Zhou,2, † and Xu Wang3, 4, ‡ 1Shenzhen Key Laboratory of Ultraintense Laser and Advanced Material Technology, Center for Intense Laser Application Technology, and College of Engineering Physics, Shenzhen Technology University, Shenzhen 518118, China 2Department of Physics, College of Science, National University of Defense Technology, Changsha 410073, China 3Graduate School, China Academy of Engineering Physics, Beijing 100193, China 4Southern Center for Nuclear-Science Theory, Institute of Modern Physics, Chinese Academy of Sciences, Huizhou, Guangdong 516000, China

available intense low-frequency laser fields. The temperature requirement of controlled nuclear fusion may be reduced with the aid of intense laser fields.

Keywords: Nuclear fusion, Intense lasers, Enhancement of fusion

Introduction

Controlled nuclear fusion is an active research field with an ultimate goal of supplying sustainable and clean energy solutions to mankind [1–6]. Yet it is difficult to achieve ignition in laboratory environments, essentially due to the fact that nuclear fusion cross sections are very small. To increase the cross section, the nuclear fuel needs to be heated to very high temperatures, typically on the order of 10^7 K. Achieving and maintaining such high temperatures are very challenging in practice. It is therefore sensible and meaningful to consider methods that may increase the fusion cross section and reduce the temperature requirement.

The possibility of using advanced light sources to influence and enhance nuclear-fusion yields has attracted attention recently [7–19]. These are interesting and timely attempts, observing rapid progress in light-source technologies, especially those with extremities in intensity or frequency (photon energy). Light sources with extreme intensities include, for example, the extreme light infrastructure (ELI) of Europe [20–22] and the superintense ultrafast laser facility (SULF) of Shanghai [23–25]. These lasers are expected to reach peak intensities on the order of 10^{23} W/cm^2 in the coming years. The frequency is in the near-infrared regime with very small single-photon energies (around 1.5 eV). Light sources with extreme frequencies include synchrotron radiations and X-ray free-electron lasers [26, 27]. These sources generate light with photon energies on the order of 1 to 10 keV.

The possibility of using these light sources to influence nuclear processes, such as α decay [28–33], nuclear fission [34], quantum optical effects in nuclei [35–38], nonlinear optical effects in nuclei [39, 40], muon production [41], measurements of cross sections and astrophysical S-factor in plasma [42, 43], and isomeric excitation [40, 42, 45–60], has been considered or experimentally demonstrated recently. It is not unreasonable to expect that these light sources can influence the nuclear fusion process. The relevant energy scale of controlled nuclear fusion is on the order of 1 keV (= 1.16×10^7 K). For high-frequency light sources, absorption of a single photon will lift the energy of the fusion system by an order of 1 keV. For low-frequency-high-intensity light sources, simultaneous absorption of 1,000 photons will lift the energy of the fusion system by a similar amount, and it will be shown later that this is actually not difficult with intensities that are readily available nowadays.

Existing works on this topic focus either on light sources with high frequencies [7, 8, 17], or those with low frequencies [9, 18]. This is mainly due to the theoretical techniques used to deal with the problem. For example, Queisser et al. employ a Floquet scattering method [7, 17], and Lv et al. adopt the Kramers-Henneberger

approximation [8], both of which are feasible only for high frequencies. Several studies have comparatively discussed different approaches for laser-assisted nuclear fusion to a certain extent [13–16], with particular emphasis on laser frequencies in the X-ray regime.

A comprehensive theoretical analysis applicable to both high and low laser frequencies is still lacking. Without an analysis that puts different laser frequencies on the same footing, conclusions cannot be drawn on what kind of laser would be most efficient in enhancing fusion. Is an X-ray free-electron laser more efficient in enhancing fusion than a near-infrared laser? One might think the answer would be yes because absorbing a single photon from an X-ray laser is equivalent to absorbing 1,000 photons from a near-infrared laser. However, it will be shown that the answer is actually no.

Extending the theoretical foundations laid by previous studies on low-frequency laser fields [9], this article provides a unified theoretical framework that covers both high-frequency lasers, such as X-ray free-electron lasers, and low-frequency lasers, such as near-infrared solid-state lasers. Different lasers are treated on the same footing. Conclusions are drawn on preferable laser parameters to enhance fusion yields. The analysis is physically oriented with the aim of providing physical understanding using the least possible numerical calculations. Fundamentally, the process of laser-assisted nuclear fusion is a complex many-body problem and an ab initio calculation starting from quantum chromodynamics remains impossible. A feasible theoretical treatment inevitably involves approximations of different levels. A highly precise and numerically intense theoretical technique is not desirable at the current stage.

This article is organized as follows. In Section II the effects of laser fields on each stage (region) of a nuclear fusion process are analyzed. Calculation results are presented in Section III. Discussions on various aspects of our analyses are given in Section IV. A summary and outlook is given in Section V to conclude the article.

II. Analyses of the Laser-Assisted Nuclear Fusion Process

A. Nuclear fusion without laser fields

We start from nuclear fusion without the presence of laser fields. A nuclear fusion process is usually divided into three regions according to the relative distance between the two nuclei, as illustrated in Fig. 1(a). Viewing from the rest frame of one nucleus (noted nucleus 1 for convenience), the other nucleus (nucleus 2) is initially in region III with an asymptotic energy E, which is usually between 1 and 10 keV depending on the temperature of the fusion environment. As it approaches, nucleus 2 will reach a classical turning point, where the Coulomb repulsive energy between the two nuclei equals the energy E. Via a quantum tunneling effect, nucleus 2 nevertheless enters and passes through region II with a small probability. In region I, the two nuclei are very close together and fusion reactions happen. The spatial range of region I is on the order of 1 fm (10^{-15}

m), and the spatial range of region II is on the order of 100 fm for typical energies of controlled fusion research.

Indeed, a fusion cross section is usually written in the following form corresponding to the three-region division [61]:

$\sigma(E) = S(E) * (1/E) * \exp(-\sqrt{(BG/E)})$

The first factor on the right-hand side, S(E), is the result of the many-body nuclear physics happening in region I. The exponential factor is the result of the tunneling process happening in region II [62]. The 1/E factor, called the geometrical factor, is related to region III. The constant BG is called the Gamow constant. For the deuteron-triton (DT) fusion reaction (D + T → 4He + n + 17.6 MeV), BG = 34.38 keV.

An ab initio calculation of the function S(E) is extremely demanding, if possible at all. For most purposes, including that of the current article, it suffices to adopt an empirical form:

S(E) = (A1 + E(A2 + E(A3 + EA4))) / (1 + E(B1 + E(B2 + E(B3 + EB4))))

where the parameters Ai's and Bi's are determined by fitting to experimental data. For the DT fusion, the values of these parameters are given as follows [63]:

A1 = 6.927×10^4, A2 = 7.454×10^8, A3 = 2.050×10^6, A4 = 5.200×10^4, B1 = 6.380×10^1, B2 = -9.95×10^{-1}, B3 = 6.981×10^{-5}, B4 = 1.728×10^{-4}.

Note that these values are accompanied with the value of E in keV when Eq. (2) is used. The DT fusion cross section $\sigma(E)$ and the function S(E) are shown in Fig. 1(b) and (c), respectively. S(E) is a slow-varying function, but $\sigma(E)$ is an exponential function due to the tunneling process. $\sigma(E)$ depends on E very sensitively, especially for relatively low E values.

B. Two nuclei in a laser field, the center-of-mass reference frame

Now let us consider effects of an external laser field on the fusion process. Consider two nuclei with charge and mass {q1, m1} and {q2, m2} placed in a plane-wave laser field. There are two commonly used gauges in light-matter interactions. One is the so-called velocity gauge, with which the laser field is characterized by a vector potential A(t). The total Hamiltonian is written as:

H = [p1 − q1A(t)]^2/2m1 + [p2 − q2A(t)]^2/2m2 + V(|r1 − r2|)

where p1 and p2 are the momenta of the two nuclei, and r = |r1 − r2| is the distance between the two nuclei. Atomic units have been used with which $4\pi_0$ = 1, hence the form of the Coulomb potential. We have assumed the validity of a long-wavelength approximation by omitting the spatial dependency of the vector potential. This is justified by the fact that the spatial range relevant to

nuclear fusion is much smaller than the wavelengths of available intense lasers. Further discussions on this point will be given in Section IV.

For a fusion process it is the relative motion between the two nuclei that is most relevant. It is therefore more convenient to work in the center-of-mass (CM) reference frame. Define:

$R = (m_1 r_1 + m_2 r_2)/(m_1 + m_2)$, $P = p_1 + p_2$, $r = r_1 - r_2$, $p = (m_2 p_1 - m_1 p_2)/(m_1 + m_2)$

Then, after some straightforward algebra, the Hamiltonian in Eq. (3) can be rewritten as:

$H = [P - QA(t)]^2/2M + [p - qA(t)]^2/2\mu + V(r)$

where $M = m_1 + m_2$, $\mu = m_1 m_2/(m_1 + m_2)$, $Q = q_1 + q_2$, and $q = (q_1 m_2 - q_2 m_1)/(m_1 + m_2)$. One sees from Eq. (8) that the Hamiltonian can be separated into a motion of the CM with charge and mass $\{Q, M\}$, and a motion of a virtual particle with charge and mass $\{q, \mu\}$. The mutual Coulomb potential is not affected. The motion of the CM is not of concern here because it is not relevant to the fusion process.

Alternatively one may use the so-called length gauge and the Hamiltonian is given as:

$H = p_1^2/2m_1 + p_2^2/2m_2 + V(|r_1 - r_2|) - (q_1 r_1 + q_2 r_2) \cdot E(t)$

where $E(t)$ is the laser electric field. Again, the long-wavelength approximation has been assumed by neglecting the spatial dependency of the laser electric field. The above Hamiltonian can also be written in the CM reference frame as:

$H = P^2/2M + p^2/2\mu + V(r) - QR \cdot E(t) - qr \cdot E(t)$

C. Assumption on region I

The size of region I is on the order of 1 fm. In this region the two nuclei fuse and new particles are generated. For the DT fusion reaction, an α particle and a neutron is generated as a result. This is a complex many-body nuclear process and an ab initio treatment of this process is usually impossible. Adding influences from an external laser field makes the process even more complicated to deal with.

It is, however, reasonable to expect that the effect of a laser field on this region is negligibly small. From the uncertainty principle, the spatial confinement of Δx 1 fm = 1.9×10^{-5} a.u. indicates an uncertainty in momentum of $\Delta p \sim 5.3 \times 10^4$ a.u., or an uncertainty in energy of $(\Delta p)^2/2\mu$ 6.4×10^5 a.u. 17 MeV. Here $\mu = 2{,}204$ a.u. has been taken as the reduced mass of deuteron and triton. In comparison, the effect of an intense laser field on the fusion process is on the order of 1 keV in energy, which is about 4 orders of magnitude smaller. It seems safe to neglect effects of external laser fields on the nuclear processes happening in region I. This is the basic assumption that our analyses start with.

D. Effects of laser fields on region II

The size of region II is on the order of 100 fm. We shall show in this section that an intense laser field has a finite but small effect on the tunneling process happening in this region. The probability of tunneling, or the "penetrability," through a laser-modified Coulomb potential barrier can be calculated using the Wentzel-Kramers-Brillouin (WKB) method as:

$$P(\theta, t) = \exp\left(-\int_{r1}^{r2} \sqrt{2\mu[VC(r) - E + VI(r, \theta, t)]}\, dr/\hbar\right)$$

where $VC(r) = q_1 q_2/r$ is the Coulomb potential and $VI(r, \theta, t) = -qr \cdot E(t) = -qrE(t)\cos\theta$ is the laser potential from Eq. (14). Here θ is the angle between the laser polarization direction and the DT collision (i.e., relative-motion) direction. The integration is performed between two classical turning points r1 and r2, which are also the two ends of region II.

What is implicit in writing Eq. (15) is a quasi-static approximation. That is, the laser potential can be viewed as static at each time. This approximation is valid when the period of the laser field is much longer than the time scale of the tunneling process. More discussions on this point will be given later in Section IV.

It can be estimated that the magnitude of VI is much smaller than that of VC, or that of $V0 \equiv VC - E$. The Coulomb potential VC in region II can be estimated to be on the order of 100 keV, and the relative collision energy E is on the order of 1 keV in typical nuclear fusion experiments. With an intensity of 10^{20} W/cm^2, the magnitude of VI can be estimated to be on the order of 10 eV. With an intensity of 10^{22} W/cm^2, the magnitude of VI is on the order of 100 eV. Therefore we may expand Eq. (15) as:

$$P(\theta, t) = \exp\left(-\int_{r1}^{r2} \sqrt{2\mu[V0(r) + VI(r, \theta, t)]}\, dr/\hbar\right) \approx \exp\left(-\int_{r1}^{r2} \sqrt{2\mu V0(r)}[1 + VI(r, \theta, t)/(2V0(r))]\, dr/\hbar\right) = \exp\left(-\int_{r1}^{r2} \sqrt{2\mu V0(r)}\, dr/\hbar\right) \times \exp\left(-\int_{r1}^{r2} \mu VI(r, \theta, t)/\sqrt{2\mu V0(r)}\, dr/\hbar\right) \equiv P(E = 0) \times \exp(\gamma^{(1)})$$

where $P(E = 0)$ is the penetrability without external laser fields and $\gamma^{(1)}$ denotes the first-order correction induced by the laser field. Substituting the expression of VI we get an explicit formula:

$$\gamma^{(1)}(\theta, t) = (\mu/\hbar^2)\int_{r1}^{r2} (VI(r, \theta, t)/\sqrt{2\mu V0(r)})\, dr = (qE(t)\cos\theta/\hbar)\int_{r1}^{r2} (r/\sqrt{2\mu V0(r)})\, dr$$

The magnitude of $\gamma^{(1)}$ is maximum at the times when E(t) reaches peaks and when $\cos\theta = 1$. The upper limit of integration r2 is determined by equating the Coulomb potential VC(r) to the collision energy E. For E = 5 keV, we get r2 = 288.3 fm. The lower limit of integration r1 is determined by a DT touching condition: $r1 = rD + rT = 1.13(A_D^{1/3} + A_T^{1/3})$ fm = 3.05 fm. We find that for intensity 10^{20} W/cm^2, $\gamma^{(1)}$ takes a maximum value of 0.18%. For intensity 10^{22} W/cm^2, the maximum value is 1.8%. For E = 10 keV, r2 = 144.2 fm. The corresponding maximum value of $\gamma^{(1)}$ is 0.03% for 10^{20} W/cm^2, and 0.3% for 10^{22} W/cm^2.

Therefore one can see that for currently available state-of-the-art laser intensities, the effects of laser fields on region II are finite, but quite small.

E. Effects of laser fields on region III

Major effects of a laser field on the fusion process come from region III. The laser field can substantially influence the collision energy E. Without external laser fields, the incoming relative-motion virtual particle is usually described asymptotically as a plane wave state:

$(r, t) = \exp(ip \cdot r - iEt/\ )$

where the momentum has magnitude $p = \sqrt{(2\mu E)}$. This plane-wave state has a well-defined energy E.

In the presence of a laser field, the asymptotic plane-wave state becomes a Volkov state [64]:

$\_V(r, t) = \exp(ip \cdot r - iEt/\ ) \times \exp(-i \ ^t HI(t') \ dt'/\ )$

where HI is the interaction Hamiltonian with the external laser field. It is convenient to use the velocity gauge here, and from Eq. (8) $HI(t) = -(q/\mu)p \cdot A(t) + (q^{2/2\mu)A}2(t)$.

Let us assume that the laser field is linearly polarized along the z axis, and the vector potential $A(t) = \hat{z}\ A0\ \sin(\omega t)$. The case of elliptical or circular polarization will be discussed later in Section IV. Note that because the laser field can be very intense, the A^2 term cannot be simply ignored as in low-intensity situations.

The Volkov state can be expanded in terms of photon number states:

$\_V(r, t) = e^{ip \cdot r}\ \_n\ e^{iF\_n(u,v)}\ e^{-i(E+U\_p+n\omega)t/\ }$

where the coefficient $F\_n(u, v)$ is given by the following integral:

$F\_n(u, v) = (1/2\pi)\ \_{-\pi}^{\pi}\ \exp(-iu \cos\ + iv \sin2\ + in\ )\ d$

For convenience we have defined $U\_p = q^{\{\{\{2A0\}\}\{2/4\mu\}\}}$ (the ponderomotive energy), $u = u(\ ) = qpA0 \cos /\mu\omega$, and $v = q^{\{\{\{2A0\}\}\{2/8\mu\omega.\}\}}$ Here is the angle between p and the +z axis, and enters into the formalism through u. In a thermal environment the direction between p and the laser polarization axis is random.

One sees from Eq. (21) that in the presence of a laser field, the collision energy is no longer a well-defined single value. Instead, the energy becomes a distribution, which is centered at $E + U\_p$ (ponderomotive shift) and separated by the photon energy. The energy of the particle can be higher or lower than $E + U\_p$, corresponding to situations of absorption or emission of photons. The probability of finding the system with energy $E\_n = E + U\_p + n\omega$ is $P\_n(u, v) = |F\_n(u, v)|^2$. The total probability summing over the photon number n is equal to unity: $\_n\ P\_n = 1$.

III. Numerical Results

A. Energy distribution with different laser parameters

The energy distribution P_n depends sensitively on laser parameters, especially on the frequency (photon energy). With the same laser intensity, the energy distribution can be very different for lasers with different photon energies. This is illustrated in Fig. 2, which shows the energy distributions for six different photon energies under the same intensity. The bare collision energy without the laser fields is assumed to be 5 keV (corresponding to a temperature of 5.8×10^7 K).

One can see that for a high-frequency laser with photon energy 1 keV (1,000 eV), almost all the population is still with the original collision energy 5 keV. The probability of absorbing (emitting) a photon and changing the energy to 6 keV (4 keV) is very small, with a value of 2.6×10^{-6}. This probability is not visually distinguishable in linear scale as shown in Fig. 2(a). The probability of absorbing (emitting) two photons and changing the energy to 7 keV (3 keV) is on the order of 10^{-12}. It is very difficult to absorb (emit) energy from (to) a high-frequency laser field, even if the intensity is high. The probability of absorbing or emitting more photons shows a perturbative feature. That is, the probability drops substantially as the number of photons increases.

As the photon energy decreases to 100 eV [Fig. 2(b)], the probability of absorbing (emitting) a photon increases to about 0.025. This means that 2.5% of the population has an energy of 5.1 keV and another 2.5% of the population has an energy of 4.9 keV. The probability of absorbing (emitting) two photons is on the order of 10^{-4}. As the photon energy decreases to 50 eV [Fig. 2(c)], the probability of absorbing (emitting) one photon is about 27%, and that of absorbing (emitting) two photons is about 3.3%. The probability of remaining with the original collision energy drops to 39%.

As the photon energy decreases to 30 eV [Fig. 2(d)], the energy distribution shows clear nonperturbative features. For example, the probability of absorbing (emitting) two photons is higher than that of absorbing (emitting) one photon. The number of photons absorbed or emitted is about 5, and the energy range populated is roughly from 4.85 to 5.15 keV.

As the photon energy decreases to 10 eV [Fig. 2(e)], the number of photons absorbed or emitted is about 40. The populated energy range is between 4.6 and 5.4 keV. Besides, the distribution shows an overall structure with peaks near two ends and a valley in the middle. This means that the collision has substantial probabilities with energies away from the original bare energy. As the photon energy decreases to 1.55 eV [Fig. 2(f)], which corresponds to a wavelength of 800 nm from Ti:sapphire intense lasers, the number of photons absorbed or emitted is over 1,300, and energy range between 3.1 and 7.3 keV is substantially populated. In this case the ponderomotive shift U_p = 109 eV, and notice that the distribution is not exactly symmetric.

The main message from the above results is that it is easier for low-frequency lasers to deliver energy to the fusion system. Although the energy of a single photon is small, the number of participated photons can be very large so that the populated energy range is wide.

Note that the energy distribution P_n depends on the parameter u, which depends on the angle between the laser polarization direction and the collision (relative-motion) direction. We have set = 0° for all the cases shown in Fig. 2. As increases from 0° to 90°, the energy distribution becomes narrower, as shown in Fig. 3 for the case of 1.55 eV (800 nm). P_n for > 90° is the same as that for (180° −). It is easier for the laser to deliver energy to the fusion system if the axis of collision and the axis of laser polarization aligns.

B. Enhancement of fusion

We see that in the presence of a laser field, the collision energy E changes from a single value to a distribution, the character of which depends on the laser parameters. The fusion system can either absorb energy from the laser field leading to collision energies higher than E, or lose energy to the laser field leading to energies lower than E (the center of the energy distribution is ponderomotively shifted to E + U_p though). Energies higher than E lead to higher fusion yields and energies lower than E lead to lower fusion yields.

The net effect, however, is an enhancement of the fusion yield. This is because the cross section function of Eq. (1) depends on the collision energy exponentially (concaving upward) as shown in Fig. 4 in linear scale. The fusion yields gained with higher energies are greater than the fusion yields lost with lower energies. This is the mechanism of fusion-yield enhancement in the presence of a laser field.

C. Effective fusion cross section

It seems sensible to define an effective fusion cross section in the presence of a laser field. We denote this laser-assisted cross section as $\sigma\_L(E)$, which can be calculated by averaging over all angles between the collision direction and the laser polarization direction:

$\sigma\_L(E) = \_0\hat{}\pi \sigma\_L(E,) \sin d / \_0\hat{}\pi \sin d$

where $\sigma\_L(E,)$ is defined as:

$\sigma\_L(E,) = \_n P\_n[u( ), v] \sigma(E + U\_p + n\omega)$

Fig. 4 illustrates the enhancement mechanism. Without laser fields, the bare collision energy is assumed to be 5 keV, corresponding to a fusion cross section of $0.538\times$10^{-3} barn (dashed horizontal line). In the presence of a laser field, the collision energy becomes a distribution (an example is shown as the red curve, left axis), with energies higher and lower than 5 keV. Higher energies lead to higher fusion yields and lower energies lead to lower fusion yields. However,

the fusion yields gained with higher energies are greater than the fusion yields lost with lower energies because the cross section function is an exponential function, concaving upward (black solid curve, right axis).

The effective cross section $\sigma\_L(E)$, in comparison to the laser-free cross section $\sigma(E)$, gives a quantitative measure of the effect of a laser field on the nuclear fusion process. It is to be emphasized that the laser field mainly affects the fusion process via modifying the collision energy in region III before the tunneling process. It has very small effects on the processes happening in region I and region II, as explained above.

Fig. 5 shows the laser-assisted effective DT fusion cross section for three bare energies (E = 1, 5, and 10 keV) and different laser intensities and photon energies. The three energies represent typical temperatures in thermonuclear fusion experiments. One sees from Fig. 5 that to have noticeable effects on nuclear fusion, the laser intensity needs to be higher than 10^{18} or 10^{19} W/cm^2.

One can also see that σ_L is larger for lower laser frequencies under the same intensity. For all three cases and with the intensity range shown, the fusion enhancement for the higher frequency cases (100 eV and 1,000 eV) is very small. This is the direct consequence of the point explained above that it is low-frequency lasers that are more efficient in delivering energy to the fusion system. It is difficult to absorb energy from a high-frequency laser field.

Substantial enhancements can be seen with the low-frequency 1.55 eV lasers. For E = 1 keV, the enhancement is 3 orders of magnitude at intensity 10^{20} W/cm^2 and 9 orders of magnitude at intensity 5×10^{21} W/cm^2. The ratio of enhancement decreases as E increases. This is because the cross section function $\sigma(E)$ of Eq. (1) depends more sensitively on E for smaller E values. The effect of the laser field is more pronounced for smaller E values.

The possibility of using lasers to reduce the temperature requirement of controlled fusion research can be seen. For example, without laser fields the DT fusion cross section at E = 1 keV is $1.37 \times 10^{-11} barn. With an 800-nm (1.55-eV) laser field of intensity 10^{20} W/cm^2, the effective fusion cross section becomes 1.02 \times 10^{-8}$ barn, which is equal to the cross section value at E = 1.6 keV without laser fields. If the laser intensity is 5×10^{21} W/cm^2, the effective cross section becomes 0.027 barn, which is equal to the cross section value at E = 10 keV without laser fields. Putting it another way, the enormous gap in the DT fusion cross section between 1 keV $(1.16 \times 10^7 K) and 10 keV (1.16 \times 10^8$ K) is filled, or compensated, completely by the intense laser field.

IV. Discussions

A. Applicability to other fusion reactions

Although the results presented above are for the DT fusion reaction, our analyses apply to other fusion reactions as well, for example, the neutronless proton-boron (p-11B) fusion reaction [65–69], which is limited in its practical applica-

tion due to its significantly lower fusion cross section compared to that of DT fusion under low bare energy conditions.

The p-11B interaction exhibits a larger effective charge, as q = 0.5, in contrast to the value of q = 0.2 for the deuteron-triton (DT) system. Furthermore, the reduced mass for the p-11B pair is notably smaller, with μ   1683.2 a.u., compared to μ   2203.4 a.u. for the DT pair. Consequently, under identical laser parameters, the ponderomotive energy, given by $U\_p = q^{\{\{\{2A0\}\}}\{2/4\mu,\}}$ is higher for the p-11B system than for the DT system. The parameters u = qpA0 cos /μω and $v = q^{\{\{\{2A0\}\}}\{2/8\mu\omega\}\}}$ indicate a greater number of photons absorbed or emitted by the p-11B pair compared to the DT pair in laser fields.

Fig. 6(a) illustrates that the energy distribution of the Volkov state for the p-11B system is broader than that of the DT system, as depicted in Fig. 3. The laser is characterized by a photon energy of 1.55 eV and a peak intensity of 10^{20} W/cm^2.

Fig. 6(b) presents the laser-assisted effective cross section for p-11B fusion as a function of bare collision energy, assuming a laser photon energy of 1.55 eV. Various laser intensities are considered, as indicated in the figure. The cross section in the absence of laser fields, as given in Ref. [70], is indicated by the black dashed line. Notably, under the influence of a strong laser field with an intensity of 5×10^{21} W/cm^2, the effective cross section of p-11B fusion can reach the order of 10^{-4} barn in the bare energy range of 1–10 keV, representing a significant enhancement in fusion reactivity. One can also see that the resonant peak of the p-11B fusion cross section, originally located around 150 keV, shifts toward a lower and broader bare energy range under the influence of the laser field. This further indicates an effective enhancement in fusion reactivity. These results highlight the potential of laser-assisted p-11B fusion as a promising pathway for future research and development in clean, radiation-free energy generation.

However, laser fields have no effect, at least within the approximations adopted in the current article, on fusion reactions with the two nuclei having the same charge-to-mass ratio. An example of this kind is the deuteron-deuteron fusion reaction. If so, one can find from Eq. (12) that the charge of the relative-motion virtual particle is zero and the laser ceases to have an effect on the relative-motion degree of freedom. This is easy to understand because the motion of a charged particle in a laser field is determined by its charge-to-mass ratio. If two nuclei have the same charge-to-mass ratio, then their motion in the laser field will be the same and the laser field has no tendency of separating them apart or pressing them closer, i.e., the laser field has no effect on the relative motion between the two nuclei.

B. The long-wavelength approximation

We have assumed the validity of the long-wavelength approximation by neglecting the spatial dependency of the laser vector potential or the electric field. This

is justified if the spatial range relevant to fusion is much smaller than the laser wavelength. The former range can be estimated by a quiver motion amplitude $z_0 = qA_0/\mu\omega$, which is the amplitude of spatial oscillations of a free charged particle in a laser field.

The required validity condition of the long-wavelength approximation is $z_0 \ll \lambda = 2\pi c/\omega$, or $A_0 \ll 2\pi c\mu/q$. Putting in the values of μ and q for the DT fusion we get $A_0 \ll 9.5 \times 10^6$ a.u., or the amplitude of the laser electric field $E_0 = A_0\omega \ll 9.5 \times 10^6\omega$ in a.u. Let us give two examples. For photon energy 1,000 eV, $\omega = 36.8$ a.u., so $E_0 \ll 3.5 \times 10^8$ a.u. or equivalently the intensity $I \ll 4.3 \times 10^{33}$ W/cm^2. For photon energy 1.55 eV, $\omega = 0.0569$ a.u., so $E_0 \ll 5.4 \times 10^5$ a.u. or equivalently the intensity $I \ll 1.0 \times 10^{28}$ W/cm^2. The validity of the long-wavelength approximation can be guaranteed for both examples.

C. The quasi-static approximation

We mentioned in Section II.D that in writing Eq. (15) we have implicitly used the quasi-static approximation. This approximation is valid when the laser period is much longer than the time scale of the tunneling process, which can be estimated using a classical picture. For example, for E = 5 keV, the velocity of relative motion is $v = \sqrt{(2E/\mu)} = 0.41$ a.u. The tunneling length (from the tunneling entrance point to the tunneling exit point) is 5.38×10^{-3} a.u. (285 fm). The estimated time for the tunneling process is therefore 0.013 a.u. or 0.32 as (1 as $= 10^{-18}$ s). As long as the laser period is much longer than 0.32 as, or the photon energy much lower than 13 keV, the quasi-static approximation is valid.

We mention here that the quasi-static approximation is also an important concept in strong-field atomic ionization [71, 72]. The ratio between the classically estimated electron tunneling time and the laser period is called the Keldysh parameter [71]. The quasi-static approximation is valid when the Keldysh parameter $\gamma \ll 1$, that is, the time scale of tunneling is much shorter than the laser period.

D. The Coulomb Volkov state

For simplicity we have been using the (plane-wave) Volkov state to describe the relative-motion virtual particle. The Volkov state is the solution of the time-dependent Schrödinger equation for a free charged particle in the presence of a laser field. The Coulomb potential between the two nuclei is neglected.

The quantum state of the full Coulomb-plus-laser system does not have general analytical solutions. An approximate solution is the so-called Coulomb-Volkov state [73, 74], which has the same form as the plane-wave Volkov state except to replace the $e^{ip\cdot r}$ in Eqs. (19) or (21) with a Coulomb wave function $\_p(r)$. The temporal part, hence the energy distribution, remains the same. Therefore the results and discussions presented above are not affected if Coulomb Volkov states are used.

E. Elliptical or circular polarization

Extension of the above formalism to elliptically or circularly polarized laser fields is straightforward, so we only outline a few steps here. Assume the vector potential is in the z-y plane: $A(t) = \hat{z} A\_ \sin\omega t + \hat{y} A\_ \cos\omega t$, with ellipticity and amplitude $A\_ = A0/\sqrt{(1 + \hat{}2)}$. Then the Volkov wavefunction can be written in the same form as Eq. (21) except that $F\_n = F\_n(u, w, v)$ having an additional argument w due to the additional polarization direction. Here $u = u() = qpA\_ \cos /\mu\omega$, $w = w(,) = qp A\_ \sin \sin /\mu\omega$, and $v = q^{2A}\_ ^2(1 - ^2)/8\mu\omega$. $(,)$ are the direction angles of the momentum p. The coefficient F_n is obtained via the integral:

$$F\_n(u, w, v) = (1/2\pi) \_{-\pi}^{\pi} \exp(-iu \cos + iw \sin + iv \sin2 + in) d$$

Because elliptical polarization breaks the cylindrical symmetry of linear polarization, the effective fusion cross section σ_L needs to be averaged over both and . Fig. 7 shows how the angle-averaged σ_L changes with for intensity 1×10^{20} W/cm^2. Elliptical or circular polarization does not lead to further enhancements in the fusion yield. The efficiency is (slightly) lower than linear polarization, mainly due to the reduction of the field amplitude from A0 to $A\_$.

F. Effects of plasma screening

In laser-induced plasmas, the range of charge screening effects can be characterized by a Debye sphere, with a radius given by the Debye length λ_D $740\sqrt{(T\_e/n\_e)}$, where λ_D is in cm, T_e in eV, and n_e in units of cm^{-3}. Plasma screening becomes significant when the inter-nuclear separation exceeds the Debye length, as the Coulomb interaction is substantially shielded by the surrounding plasma electrons. Conversely, when the inter-nuclear distance is shorter than the Debye length, the screening effect is relatively weak, and the bare nuclear interaction predominates.

Assuming a plasma electron temperature of 2 keV and an electron density of 10^{17} cm^{-3}, the corresponding Debye length is approximately 1 µm. This characteristic length scale is significantly larger than the amplitude of the quiver motion between two nuclei in a laser field with an intensity of 5×10^{21} W/cm^2 and a photon energy of 1.55 eV, which is approximately 0.56 nm. As a result, the inter-nuclear separation remains well below the Debye length, and the influence of plasma screening is negligible. Therefore, the two-nuclei interaction model employed in this study is well justified in disregarding the effects of plasma screening.

The plasma environment may also involve additional effects, such as collisional interactions and radiation processes, which could be further investigated through advanced laser-plasma simulations in future work.

V. Summary and Outlook

In summary, we have considered the process of nuclear fusion in the presence of a laser field. Without laser fields, a nuclear fusion process is usually treated as a three-region process, and we have analyzed the effects of laser fields on each of the three regions. Our analysis is physically oriented, aiming to provide a clear physical understanding of the laser-assisted nuclear fusion process. We show that major effects of the laser field on the nuclear fusion process come from influencing the collision energy before tunneling. We explain why this influence in the collision energy leads to enhanced fusion yields. By treating lasers with different frequencies on the same footing, we are able to draw conclusions on optimal laser parameters to enhance fusion. We show that intense low-frequency lasers are most efficient in delivering energy to the fusion system and enhancing the fusion yield.

The possibility is pointed out that lasers may be used to reduce the temperature requirement of controlled fusion research. The vast difference between fusion cross sections at different temperatures shrinks in the presence of laser fields. Controlled fusion experiments might be performed at lower temperatures with the aid of intense laser fields.

In this article we only consider the (pure) system of two nuclei plus a laser field. We have not considered a more complicated plasma environment with nuclei, electrons, laser-plasma interactions, etc. These complications are important, but outside the scope of the current article. As the first step we need to understand the laser-assisted nuclear fusion process and this is the goal of the current article. Adding the above-mentioned complications into the picture and evaluating the effect of laser fields is the next step to go.

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