I. INTRODUCTION

Investigations to study the effects of screened Coulomb interaction among the charged particles in plasmas on the atomic collision properties are of utmost importance in the sense that such studies can provide fruitful information regarding the radiation and transport properties of the plasmas, and also help us to interpret various phenomena associated with astrophysics, hot plasma physics, and experiments performed with charged ions. As a result, atomic collision in plasma has been the subject of extensive study for the last couple of years (and further references therein). Of late, the authors have investigated the effects of screening of weakly coupled plasmas on the scattering of positrons from the hydrogen atoms. The interactions among the charged particles in weakly coupled plasmas have been taken care of by Debye-Huckel potentials or screened Coulomb potentials (SSCP) Have been investigated by applying a formulation of the three-body collision problem in the form of coupled multi-channel two-body Lippmann-Schwinger equations. The interactions among the charged particles in dense quantum plasma have been represented by exponential cosine-screened Coulomb potentials. Variationally determined hydrogenic wave function has been employed to calculate the partial-wave scattering amplitude. Plasma screening effects on various possible mode of fragmentation of the system $e^+ + H(1s)$ during the collision, such as $1s \rightarrow 1s$ and $2s \rightarrow 2s$ elastic collisions, $1s \rightarrow 2s$ excitation, positronium formation, elastic proton-positronium collisions, have been reported in the energy range 13.6-350 eV. Furthermore, a comparison has been made on the plasma screening effect of a dense quantum plasma with that of a weakly coupled plasma for which the plasma screening effect has been represented by the Debye model. Our results for the unscreened case are in fair agreement with some of the most accurate results available in the literature.

\[ V(r) = (1/r)e^{-\mu r} \text{ (in a.u.),} \]

where $\mu$ is the screening parameter. In the Debye model of screening, the screening parameter $\mu$ is related to the Debye length $\lambda_D$ by means of the relation $\lambda_D = 1/\mu = [KT/4\pi e^2 N_e]^{1/2} = v_T/\omega_p$, where $e$ is the electronic charge, $N_e$ is the plasma-electron density, $K$ is the Boltzmann constant, $T$ is the electron temperature, $v_T$ is the thermal velocity, and $\omega_p$ is the plasma frequency. In weakly coupled plasmas, long-range self-consistent interactions (described by the Poisson equation) dominate over short-range two-particle interactions (collisions). In other words, the ratio of the potential energy to the average kinetic energy, called the coupling parameter ($\Gamma$), is much less than unity. The Coulomb interaction screening in weakly coupled plasmas is a collective effect of the correlated many-particle interactions; and in the lowest particle correlation order, it reduces to the Debye-Huckel potential, or SSCP of the form (1). So, the Debye-Huckel model can be employed to adequately describe the screened interaction potential in classical weakly coupled plasmas.

However, the Debye-Huckel model cannot be reliable for investigating the physical properties of quantum plasmas with an increase of the plasma density due to the multi-particle co-operative interactions. If the de Broglie wavelength of the charge carriers is comparable to the Debye length $\lambda_D = 1/\mu$ of the plasma, an electron inside the Debye sphere is screened by the potential outside of the sphere. This can happen if a plasma is cooled to an extremely low temperature. In such occasions, the ultracold plasma behaves as a Fermi gas, and the quantum mechanical effects play a vital role in the behaviour of the collective interactions of the charged particles. Recently, Shukla and Eliasson have shown that in a dense quantum plasma, the effective potential of a test charge of mass $m$ can be modeled by a modified Debye-Huckel potential or exponential cosine-screened Coulomb potential (ECSCP)

\[ V(r) = (1/r)e^{-\mu r} \cos(\mu r) \text{ (in a.u.),} \]

rather than Debye-Huckel potential in the form of Eq. (1). Here, $\mu$ is the screening parameter which is related to the plasma frequency $\omega_p$ by means of the relation $\mu = \omega_p/(\hbar \omega_p/m)^{1/2}$, which, in turn, is related to the quantum wave number $k_q$ by the relation $k_q = \sqrt{2\mu}$. Quantum plasmas are usually characterized by a low-temperature and a high number density and are found to exist in various nano-scale...
objects, such as nano-wires, quantum dots, semiconductor devices, and laser produced dense plasmas. In quantum plasmas, the ranges of the electron number density \( N_e \) and temperature \( T \) are, approximately, \( 10^{18} - 10^{23} \text{ cm}^{-3} \) and \( 10^2 - 10^5 \text{ K} \), and the coupling plasma parameter \( \Gamma > 1 \). However, in deriving Eq. (2), it was assumed that the quantum force acting on the electrons dominated over the quantum statistical pressure. It is to be noted that the existence of 'cos' function in the potential (2) leads to a long range oscillatory behaviour. As a result, collision dynamics in dense quantum plasma characterised by ECSCP is expected to be different from the collision dynamics in weakly coupled plasma.

In this paper, our objective is to study the effects of screening of dense quantum plasma on the following scattering processes related to the scattering of positron from the ground state of hydrogen atom:

(i) \( e^+ + H(1s) \to e^+ + H(1s) \)
(ii) \( e^+ + H(2s) \to e^+ + H(2s) \)
(iii) \( e^+ + H(1s) \to e^+ + H(2s) \)
(iv) \( e^+ + H(1s) \to p + Ps(1s) \)
(v) \( p + Ps(1s) \to p + Ps(1s) \)

To study the above processes, we use coupled multi-channel two-body Lippmann-Schwinger equations, developed by Sloan and Moore. We represent plasma screening effect of dense quantum plasma by exponential cosine-screened Coulomb potentials in the form of Eq. (3). Also we intend to investigate the changes in the collisional properties when the background plasma environment changes from weakly coupled to dense quantum. Such a study is of the utmost importance in determining the radiation and transport properties of plasmas, and also plays a role in the research into inertial confinement fusion, and astrophysics (stellar atmospheres and interiors).

The paper is organized as follows. We shall briefly discuss coupled multi-channel two-body Lippmann-Schwinger equations in Sec. II. In Sec. III, we shall discuss our computed results; and in Sec. IV, we shall make our concluding remarks. Atomic units (a.u.) will be used in the remaining part of this paper unless explicitly indicated otherwise.

II. THEORY

The non-relativistic Hamiltonian (in a.u.) of the positron-hydrogen system in dense quantum plasma characterized by the exponential cosine-screened potential (2) is given by

\[
H = -\frac{1}{2} \nabla_1^2 - \frac{1}{2} \nabla_2^2 + \frac{e^{-\mu r_1}}{r_1} \cos(\mu r_1) - \frac{e^{-\mu r_2}}{r_2} \cos(\mu r_2)
- \frac{e^{-\mu r_{12}}}{r_{12}} \cos(\mu r_{12}),
\]

where \( r_1 \) and \( r_2 \) are, respectively, the coordinates of particle 1 (positron) and particle 2 (electron) relative to the particle 3 (proton assumed to be at rest), and \( r_{12} \) is their relative distance. Let \( V_{12} \) or \( V_3 \) denotes the interaction between the pair of particles 1 and 2, then

\[
V_3 = V_{12} = -\frac{e^{-\mu r_{12}}}{r_{12}} \cos(\mu r_{12}),
\]

\[
V_1 = V_{23} = -\frac{e^{-\mu r_2}}{r_2} \cos(\mu r_2)
\]

and

\[
V_2 = V_{31} = -\frac{e^{-\mu r_1}}{r_1} \cos(\mu r_1).
\]

Therefore, \( H \) can be written as

\[
H = H_0 + V_1 + V_2 + V_3,
\]

such that

\[
(H_0 + V_1)\Phi_{\mathbf{k}n} = E_{\mathbf{k}n}\Phi_{\mathbf{k}n},
\]

where \( \Phi_{\mathbf{k}n} \) denotes the state in which particle \( \gamma \) is free and other pair are in the bound state \( \phi_{\mathbf{k}n} \) with channel momentum \( \mathbf{k} \), and

\[
E_{\mathbf{k}n} = \frac{k^2}{2M_\gamma} + \epsilon_{\mathbf{k}n}.
\]

Here, \( \epsilon_{\mathbf{k}n} \) is the energy of the bound state \( \phi_{\mathbf{k}n} \), and \( M_\gamma \) is the reduced mass of particle \( \gamma \) and the bound pair. Following Sloan and Moore, the formally exact three-body scattering amplitude for a transition from bound state \( n \) in channel \( \alpha \) to bound state \( n' \) in channel \( \beta \) is given by

\[
(k' n'|Y_{\beta\alpha} | \mathbf{k}n) = (k' n'|Y_{\beta\alpha} ^{(1)} | \mathbf{k}n)
+ \sum_{\gamma} \sum_{n''} \left( \int dk'' \frac{\langle k' n'|k'' n''|k'' n''|Y_{\gamma\alpha} | \mathbf{k}n \rangle}{E - E''_{\gamma} + i\epsilon} \right),
\]

where \( Y_{\beta\alpha} \) and \( Y_{\beta\alpha} ^{(1)} \) denote the three-body and two-body transition operators, and \( k' \) and \( k \) are the momenta in the channel \( \beta \) and \( \alpha \), respectively, with \( E = E_{\beta\alpha} = E_{\mathbf{k}n} \). The summation over \( \gamma \) runs over different channels and summation over \( n'' \) runs over different bound states in a particular channel. The pole term \( (E - E''_{\gamma} + i\epsilon)^{-1} \) appearing in Eq. (4) can be expressed as a sum of \( \delta \)-function and principal-value part as

\[
\frac{1}{E - E''_{\gamma} + i\epsilon} = 2M_\gamma(k^2 - k''^2 + i\epsilon)^{-1}
- i\pi\delta(E - E''_{\gamma}) + P \frac{1}{E - E''_{\gamma}}.
\]

In the present formalism, only the physical amplitudes are taken into account, so that the pole term is approximated by retaining only the delta function. This approximation is quite justified for the high-energy region. With this approximation, Eq. (4) takes the form

\[
(k' n'|Y_{\beta\alpha} | \mathbf{k}n) = (k' n'|Y_{\beta\alpha} ^{(1)} | \mathbf{k}n)
- i\pi \sum_{\gamma} \sum_{n''} \left( \int dk'' \frac{\langle k' n'|k'' n''| \mathbf{k}n \rangle \delta(E - E''_{\gamma})}{E - E''_{\gamma}} \times \langle k'' n''|Y_{\gamma\alpha} | \mathbf{k}n \rangle \right).
\]
This is the expression for the three-body amplitude that is applied to study the processes (i)–(v) mentioned above.

We denote the elastic channel $e^+ + H(1s)$ by channel 1, the inelastic channel $e^+ + H(2s)$ by channel 2, and the rearrangement channel $p + Ps(1s)$ by channel 3. Then, taking $\alpha, \beta = 1, 2$, and retaining only 1s and 2s states in the summation over \( \nu \), we obtain the following set of coupled equations for 1s–1s, 1s–2s, 2s–1s, and 2s–2s transitions:

\[
\langle k'1s | Y_{11} | k1s \rangle = \langle k'1s | Y_{11} | k1s \rangle \\
- i \pi \int dk'' \langle k'1s | Y_{11} | k''1s \rangle \delta(E - E'') \times \langle k''1s | Y_{11} | k1s \rangle \\
- i \pi \int dk'' \langle k'1s | Y_{11} | k''1s \rangle \delta(E - E'') \times \langle k''2s | Y_{21} | k1s \rangle,
\]

\[
\langle k'2s | Y_{21} | k1s \rangle = \langle k'2s | Y_{21} | k1s \rangle \\
- i \pi \int dk'' \langle k'2s | Y_{21} | k''1s \rangle \delta(E - E'') \times \langle k''1s | Y_{11} | k1s \rangle \\
- i \pi \int dk'' \langle k'2s | Y_{21} | k''1s \rangle \delta(E - E'') \times \langle k''2s | Y_{22} | k2s \rangle,
\]

\[
\langle k'1s | Y_{12} | k2s \rangle = \langle k'1s | Y_{12} | k2s \rangle \\
- i \pi \int dk'' \langle k'1s | Y_{12} | k''1s \rangle \delta(E - E'') \times \langle k''1s | Y_{12} | k2s \rangle \\
- i \pi \int dk'' \langle k'1s | Y_{12} | k''1s \rangle \delta(E - E'') \times \langle k''2s | Y_{22} | k2s \rangle,
\]

Neglecting the higher order terms of the multiple scattering series, the matrix elements of the two-body operator $Y_{2ij}$ and the matrix elements of the three-body operator $Y_{3ij}$ can be approximated as

\[
\langle k'1s' | Y_{11} | k1s \rangle = - \frac{\mu_{ij}^B}{4\pi^2} f_{ij}^B(k', k) \]  
and \ \[\langle k'1s' | Y_{11} | k1s \rangle = - \frac{\mu_{ij}^B}{4\pi^2} f_{ij}^B(k', k),\]  respectively, where $f_{ij}^B(k', k)$ denotes the corresponding first Born amplitude and $\mu_{ij}^B$ is the reduced mass in the channel $\beta$. Now performing partial wave analysis on Eq. (6), such as

\[
f_{ij}^B(k', k) = \frac{1}{\sqrt{k'k}} \sum_{l=0}^{\infty} (2l + 1) T_l^j(\beta \rho) P_l(\hat{\rho} \cdot \hat{k}),
\]

we obtain the following set of coupled algebraic equation:

\[
T_l^j(11) = T_l^j(11) + i T_l^j(12) T_l^j(21) \\
1 - i T_l^j(21) + i T_l^j(21) T_l^j(11) \\
1 - i T_l^j(21) T_l^j(11),
\]

\[
T_l^j(21) = T_l^j(21) + i T_l^j(12) T_l^j(11) \\
1 - i T_l^j(12) T_l^j(11) \\
1 - i T_l^j(12) T_l^j(11),
\]

\[
T_l^j(12) = T_l^j(12) + i T_l^j(12) T_l^j(11) \\
1 - i T_l^j(12) T_l^j(11) \\
1 - i T_l^j(12) T_l^j(11),
\]

\[
T_l^j(22) = T_l^j(22) + i T_l^j(21) T_l^j(11) \\
1 - i T_l^j(21) T_l^j(11) \\
1 - i T_l^j(21) T_l^j(11),
\]

Solving this set of coupled algebraic equations, we obtain the partial-wave scattering amplitudes for the transitions 1s–1s, 1s–2s, 2s–1s, and 2s–2s are obtained. Similarly taking $\alpha, \beta = 1, 3$, and retaining only $H(1s)$ and $Ps(1s)$ states in the summation over \( \nu \), we obtain the following set of coupled algebraic equations:

\[
T_l^j(11) = T_l^j(11) + i T_l^j(13) T_l^j(31) \\
1 - i T_l^j(31) T_l^j(11) \\
1 - i T_l^j(31) T_l^j(11),
\]

\[
T_l^j(31) = T_l^j(31) + i T_l^j(13) T_l^j(33) \\
1 - i T_l^j(33) T_l^j(11) \\
1 - i T_l^j(33) T_l^j(11),
\]

\[
T_l^j(13) = T_l^j(13) + i T_l^j(31) T_l^j(33) \\
1 - i T_l^j(33) T_l^j(11) \\
1 - i T_l^j(33) T_l^j(11),
\]

\[
T_l^j(33) = T_l^j(33) + i T_l^j(31) T_l^j(33) \\
1 - i T_l^j(33) T_l^j(11) \\
1 - i T_l^j(33) T_l^j(11),
\]

Solving this set of coupled algebraic equations, we obtain the partial-wave scattering amplitudes for $e^+ + H(1s)$ elastic collisions, positronium formation and $p + Ps(1s)$ elastic collisions.

In order to determine the bound state energies and wave functions of hydrogenic atom in dense quantum plasma characterized by the exponential cosine-screened potential (2), we solve the Schrodinger equation $h_1 \phi = E \phi, E < 0$, with

\[
h_1 = - \frac{1}{2} \nabla^2 - \frac{e^{-\mu r}}{r^2} \cos(\mu r_2)
\]

within the framework of Ritz's variational principle by employing the trial wave function in the form

\[
\phi_{nlm}(r_2) = R_{nl}(r_2) Y_{lm}(\theta, \phi),
\]

where $Y_{lm}$ is the spherical harmonics and the radial function $R_{nl}$ satisfies the radial equation

\[
h_2 R_{nl} = \left[ - \frac{d}{dr} \left( r^2 \frac{d}{dr} \right) - \frac{e^{-\mu r}}{r^2} \cos r_2 + l(l+1) \right]
\]

R_nl = E_n R_{nl}.

In our present investigation, we choose the following simple radial wave function:
where $\lambda$ is a variational parameter and $L_{n+l}^{2l+1}$ is the associated Laguerre polynomial of degree $(n+l)$ and order $(2l+1)$. It is interesting to note that $R_{\alpha l}(r_2)$ in Eq. (14) corresponds to the exact radial wave function of the free hydrogenic atom if $\lambda = a_0$, the first Bohr’s radius. Choice of this type of hydrogenic wave function makes our calculation for the partial-wave amplitudes $T_l$ easy. Only the calculations of the partial-wave amplitudes $T_l^0(31), T_l^0(33)$ require to invoke the Lewis integral.\(^{34}\)

Using the expression for scattering amplitude $Y_{\beta s}$, we calculate the differential cross section as a function of scattering angle and energy for positron-hydrogen collisions as

$$
\frac{d\sigma}{d\omega} = \frac{\mu_{ik}}{\mu_{ik}} |Y_{\beta s}(\mathbf{k}', \mathbf{k})|^2 = \frac{k_f}{k_i} |Y_{\beta s}^{(r)}(\mathbf{k}', \mathbf{k})|^2 + |Y_{\beta s}^{(i)}(\mathbf{k}', \mathbf{k})|^2.
$$

(15)

where $Y_{\beta s}^{(r)}$ and $Y_{\beta s}^{(i)}$ are the real and imaginary parts of $Y_{\beta s}$, respectively.

### III. RESULTS AND DISCUSSION

The variationally determined hydrogenic wave function (14) is reasonably accurate. Its efficiency and accuracy have been checked in many calculations.\(^{16-18,35}\) It has been seen that the eigen energies corresponding to this wave function are in close agreement with the most accurate eigen energies computed by using extensive Slater-type orbital\(^{36}\).

### TABLE I. Partial wave contribution (in $\pi a_0^2$) to the total $e^+ - H(1s)$ elastic cross section in dense quantum plasma and weakly coupled plasma at 54.4 eV incident positron energy. Column A stands for dense quantum plasma and column B stands for weakly coupled plasma. $\sigma^d$ denotes the total elastic cross section in units of $\pi a_0^2$: (a) Results of Ghoshal and Ho,\(^{15}\) a1: Results of Ghoshal and Ho,\(^{34}\) b: Integral approach of Basu et al.,\(^{37}\) c: Schwinger variational results of Ghoshal and Mandal,\(^{38}\) d: convergent close-coupling approach of Kadyrov and Bray\(^{39}\) (data taken from the graph within 5% accuracy), e: 18-state approximation of Kernoghan et al.,\(^{40}\) (data taken from the graph within 5% accuracy.)

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A. 1s–1s elastic collisions

In Table I, we present the results of partial-wave contributions to the total elastic cross section of the hydrogen atom in $e^+–H$ collisions in both dense quantum plasma and weakly coupled plasma at 54.4 eV of incident positron energy for some values of the screening parameter $\mu$. We have included all the significant partial-wave contributions to calculate total cross sections. These results are fairly accurate, as can be seen from the comparison of our results, for the unscreened case, with some of the most accurate results available in the literature.\(^{13,14,37-40}\) It is worthy to mention here that the present approach is most appropriate for intermediate and high incident positron energies. From Table I, we see that the trend of elastic cross section as a function of partial wave is the same for both dense quantum plasma and weakly coupled plasma; but for each partial wave, cross section is little larger in dense quantum plasma than in weakly coupled plasma for small screening effect. However, for strong screening effect, cross section is little larger in weakly coupled plasma than in dense quantum plasma. We present our results for elastic differential cross sections in dense quantum plasma as a function of scattering angle at 100 eV for different values of the screening parameter in Figure 1. It is to be mentioned here that the total elastic cross section obtained by integrating the differential cross section over the entire solid angle is the same as that obtained by summing the partial-wave contributions. This ensures the accuracy of the values of the differential cross section. This has been
done for every processes. From Figure 1, we see that for the unscreened case, differential cross section is peaked at $0^\circ$; but with increasing screening strength, this peak slowly moves forward. We also note that the trend of the elastic differential cross sections is alike in both dense quantum plasma and weakly coupled plasma. In Table II, we present the results for total elastic cross section in dense quantum plasma for different screening effects at different incident positron energies. Also in Figure 2, we present the total elastic cross section for both dense quantum plasma and weakly coupled plasma. From these table and figure, we see that cross section decreases with increasing screening strength as well as increasing incident positron energy; but for small screening effect, cross section is little larger in dense quantum plasma than in weakly coupled plasma.

B. $1s\rightarrow 2s$ transition

In Table III, we present the results of partial-wave contributions to the total inelastic cross section for $1s\rightarrow 2s$ transition processes in both dense quantum plasma and weakly coupled plasma at 100 eV of incident positron energy for some values of the screening parameter $\mu$. In this table, we also present the total inelastic cross section for the unscreened case obtained by using some of the accurate calculations, such as distorted-wave approximation, $^{16}$ 18-state approximation, $^{40}$ Pseudomodel calculation, $^{41}$ close-coupled pseudostate approximation, $^{42}$ and two-center convergent close-coupling approach. $^{39}$ Close agreement of our results with these results indicates the accuracy of our present calculation. From Table III, we notice that the contribution of the $p$–wave is larger than the contributions of other partial waves. This is an indication of the existence of maxima and minima in the surface of the differential cross section. We also notice that the trend of inelastic cross section as a function of partial wave is almost same for both dense quantum plasma and weakly coupled plasma. But for a small screening effect, cross section is little larger in dense quantum plasma.

TABLE II. Total $e^+ + H(1s) \rightarrow e^+ + H(1s)$ elastic cross section, in units of $a_0^2$, for $e^+ + H(1s) \rightarrow e^+ + H(1s)$ collisions in dense quantum plasma. (a: Results of Ghoshal and Ho.$^{14}$)

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FIG. 1. Differential cross section (a.u.) as a function of scattering angle (degree) for elastic $e^+ + H(1s)$ collisions in dense quantum plasma for different screening effect at 100 eV of the incident positron energy.

FIG. 2. Total elastic cross section for $e^+ + H(1s) \rightarrow e^+ + H(1s)$ as a function of screening parameter $\mu$ at 100 eV of the incident positron energy. WCP stands for weakly coupled plasma and DQP stands for dense quantum plasma.
plasma than in weakly coupled plasma. We present our results of differential cross sections as a function of scattering angle at 100 eV for different plasma screening effects in dense quantum plasma in Figure 3. From this figure, we note that the inelastic differential cross section is highly peaked at 0°; and with increasing screening effect, this peak moves forward. Furthermore, the nature of the differential cross section is almost the same in both dense quantum plasma and weakly coupled plasma, but the differential cross section in a particular direction is slightly larger in dense quantum plasma than in weakly coupled plasma. Also we note that most of the variations in the differential cross sections are due to the interference effects of scattered waves of different angular momentum states. At a peak (or maxima), these waves constructively interfere; whereas at a minima, the waves destructively interfere. Furthermore, orbiting of the positron around the target occurs in positron-impact excitation processes. In Figure 4, the total 1s–2s inelastic cross sections are shown as a function of screening parameter μ at 10 eV of incident positron energy in both dense quantum plasma and weakly coupled plasma. In this case, we see that effect of plasma screening is significant.

### C. 2s–2s elastic collisions

In Figure 5, we present the 2s–2s elastic cross sections as a function of incident positron energy for different

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**TABLE III. Partial wave contributions to the 1s–2s excitation cross section (in units of \( \sigma_{tot} \)) of hydrogen atom by positron impact at 100 eV in dense quantum plasma and weakly coupled plasma for various values of the screening parameter μ. Column A stands for dense quantum plasma and column B stands for weakly coupled plasma. \( \sigma^{sl} \) denotes the total 1s–2s inelastic cross section in units of \( \sigma_{tot} \). \( \alpha \) stands for \( \times 10^{-5} \).**

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<th>( \mu = 0.15 )</th>
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**FIG. 3. Differential cross section (a.u.) as a function of scattering angle (degree) for \( e^+ + H(1s) \rightarrow e^+ + H(2s) \) in dense quantum plasma for different screening effect at 100 eV of the incident positron energy.**

**FIG. 4. Total inelastic cross section for \( e^+ + H(1s) \rightarrow e^+ + H(2s) \) as a function of screening parameter μ at 50 eV of the incident positron energy. WCP stands for weakly coupled plasma and DQP stands for dense quantum plasma.**
screening effects for both dense quantum plasma and weakly coupled plasma. We notice that \( 2s-2s \) elastic cross section is high at low incident energies and decreases with increasing positron energy. Moreover, cross section decreases slightly with increasing screening effect. In this case, we notice the screening effect of plasma is not substantial.

**TABLE IV.** Partial wave contribution (in ps cm\(^{-2}\)) to the total positronium formation cross section for \( e^+ + H(1s) \rightarrow \gamma + Ps(1s) \) in dense quantum plasma and weakly coupled plasma for various values of the screening parameter \( \mu \). Column A stands for dense quantum plasma and column B stands for weakly coupled plasma. \( \sigma^b \) denotes the total positronium formation cross section in units of \( \text{ps cm}^{-2} \). Table entry \( x \times 10^y \) stands for \( x \times 10^y \).

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| \( \phi^b \) | \( 0.3035 \) | \( 0.3035 \) | \( 0.3035 \) | \( 0.3035 \) | \( 0.3035 \) |
| \( \delta^b \) | \( 0.300 \) | \( 0.300 \) | \( 0.300 \) | \( 0.300 \) | \( 0.300 \) |
D. Positronium formation

In Table IV, we present the results of partial-wave contributions to the Ps(1s) formation cross section for different screening effect along with the total Ps formation cross sections at 13.6 eV and 50 eV of incident positron energies respectively for both plasmas. In Table IV, we also present the total Ps(1s) cross section for the unscreened case obtained by using some of the accurate calculations, such as distorted wave method, Schwinger variational method, method taking effect of polarization, and 33-state approximation. Further, a comparison of our present results with theoretical studies and experimental observations is also made in Fig. 6 for the unscreened case. Close agreement of our results with these theoretical and experimental results indicates the accuracy of our present calculation. From Table IV, we see that there is a variation in the partial wave cross sections with the partial waves \( l = 0, 1, 2, \ldots \). Also we see that the trend of partial wave cross section as a function of partial wave is the same for both dense quantum plasma and weakly coupled plasma. In Figure 7, we show our results of differential cross sections at 50 eV of incident positron energy in dense quantum plasma for some values of the screening parameter \( \mu \). We note that the trend of the differential cross section is almost the same in both dense quantum plasma and weakly coupled plasma for each screening effect—differential cross section is highly peaked at \( 0^\circ \), and gradually falls down from forward to large scattering angles. In Table V, we present our results for total Ps formation cross sections in dense quantum plasma for some values the screening parameter \( \mu \) at different incident positron energies. In this table, we have also included the corresponding results for weakly coupled plasma. We would like to state here that the rows and columns of Table V in our previous paper were inadvertently interchanged. Also in Figure 8, we present the Ps formation cross sections as a function of the incident positron energy.

![Fig. 6. Positronium formation cross section (in units of \( a_0^2 \)) in the ground state as a function of incident positron energy (in eV) for \( e^+ + H(1s) \rightarrow Ps(1s) + p \) for \( \mu = 0 \). Solid line: present result, solid circle: Kadyrov and Bray, star: Kernaghan et al., asterisk: experimental results of Zhou et al., diamond: experimental results of Weber et al.](image)

![Fig. 7. Differential cross section (a.u.) as a function of scattering angle (degree) for \( e^+ + H(1s) \rightarrow p + Ps(1s) \) in dense quantum plasma for different screening effect at 50 eV of the incident positron energy.](image)

**TABLE V. Total positronium formation cross section for \( e^+ + H(1s) \rightarrow p + Ps(1s) \) in dense quantum plasma for different screening effects. \( x[y] \) stands for \( x \times 10^y \). Row a: stands for the corresponding results in weakly coupled plasma.**

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screening parameter $\mu$ for both plasmas. From this figure, we note that cross section decreases with increasing screening strength; but at 13.6 eV, energy cross section in dense quantum plasma first increases slowly and then decreases. Moreover, from Table V, we notice that that positronium formation cross section is peaked around ionization threshold.

**E. Elastic $p + Ps(1s)$ collisions**

In Table VI, we present the partial wave contribution to the elastic $p + Ps(1s)$ cross section for different plasma screening effects and at various incident proton energies for dense quantum plasma. Here, we see that screening has very little effect on the change in cross section. Appreciable change in the cross section is observed for strong screening effect. Moreover, cross section in dense quantum plasma is smaller than that of in weakly coupled plasma.

**IV. CONCLUSIONS**

In this paper, we have made an investigation to study the plasma screening effect of a dense quantum plasma on

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$\sigma^e_l$ denotes the total elastic cross section in units of $\pi a_0^2$, $x[y]$ stands for $x \times 10^y$.  

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</table>

$\sigma^e_l$ denotes the total elastic cross section in units of $\pi a_0^2$, $x[y]$ stands for $x \times 10^y$.  

**TABLE VI.** Partial wave contribution (in $\pi a_0^2$) to the total elastic cross section for $p + Ps(1s) \rightarrow p + Ps(1s)$ in dense quantum plasma.  

FIG. 8. Total positronium formation cross section in the ground state for positron-hydrogen collisions at (a) 13.6 eV and (b) 50 eV. WCP stands for weakly coupled plasma and DQP stands for dense quantum plasma.
The three-body collision problem in the form of coupled multi-channel two-body Lippmann-Schwinger equation. Making use of a simple variationally determined hydrogenic wave function scattering amplitudes has conveniently been calculated by solving a set of coupled algebraic equations. We report the results of differential and total cross sections for various mode of fragmentation of positron-hydrogen system in dense quantum plasma for a wide range of screening parameter. To the best of our knowledge, the results for 2s–2s elastic collision and $p - Ps(1s)$ elastic collisions in dense quantum plasma are reported for the first time in the literature. Our results for the unscreened case are in fair agreement with the most accurate results available in the literature. We have also made a comparative study on the positron-hydrogen collisions in dense quantum plasma and weakly coupled plasmas. We hope that our present investigation will provide useful information to the research community of atomic physics, astrophysics, plasma physics, and few-body systems in atomic physics.

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