

Research Article

Exploring QGP-Like Phenomena with Charmonia in p+p Collisions at $\sqrt{s}=13$ TeV

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In ultra-relativistic collisions of nuclei at the Large Hadron Collider, the created QCD environment rapidly changes, leading to a non-adiabatic evolution of the quantum states involved. Considering this, we first examine the pre-equilibrium state of QCD matter and its effect on the initially produced charmonium using a temperature-independent Hamiltonian. As the QCD matter reaches local thermal equilibrium, this Hamiltonian transforms to its finite temperature counterpart. To model the pre-equilibrium stage, we use the bottom-up thermalization approach to determine the effective temperature of the QCD matter, followed by a Gubser-type expansion for the thermalized medium. Additionally, we consider collisional damping, gluonic dissociation, and regeneration mechanisms, which specifically modify the charmonium yield in the thermalized medium. Mainly, the gluonic dissociation and collisional damping cause a reduction in the yield conversely, regeneration through gluonic de-excitation enhances the yield of charmonium. Further, we explore the combined effects of these mechanisms on the collective yield of charmonium states with transverse momentum (p_T) and event multiplicity in the proton-proton collisions at $\sqrt{s} = 13$ TeV. Based on our findings, we contend that the combined effects of these mechanisms can serve as a robust probe for determining the possible existence of a thermalized QCD medium in such a small collision system.

I. Introduction

The suppression of quarkonia has been proposed as an efficient probe for the creation of the transient phase of quark-gluon plasma (QGP) in heavy-ion collisions^{[1][2]}. In QGP-like scenarios, quarkonia suppression arises from the breaking of the heavy quark-antiquark pair ($Q - \bar{Q}$) and the screening of the QCD potential, the transition from color neutral to a colored state^{[1][3][4]}. In the heavy-ion collisions, even in the absence of a QGP-like medium, the quarkonia production itself gets suppressed

to a certain extent due to the presence of a nuclear environment in the colliding ions, such phenomena are incorporated in the cold nuclear matter (CNM) effects^[5]. However, the CNM effect and the QGP effect were separately unable to explain the experimental data for quarkonia suppression from heavy-ion collisions at RHIC and LHC energies. The inclusion of these two effects jointly helped to explain the quarkonia suppression data qualitatively in Au–Au, Pb–Pb, p –Pb collisions at RHIC and LHC energies^{[4][6][7][8][9]}. In these studies, gluonic dissociation, collisional damping, and color screening are the main effects reducing the effective yield of the quarkonia in the QGP medium. Under CNM effects, authors considered the shadowing effect, which modifies the initial production of the quarkonia in heavy-ion collisions. Besides the suppression or mechanism of yield reduction, an enhancement in the yield due to secondary production or regeneration of quarkonia in the QGP medium is also considered. Meanwhile, quarkonia suppression in heavy-ion collisions is an interplay of various phenomena including cold and hot nuclear matter effects.

The quarkonium suppression in heavy-ion collisions is studied by considering results from $p + p$ collisions as a baseline^{[2][10][11][12]}. The $p + p$ collision is used as a benchmark because it is assumed that such collisions lack the nuclear environment and are also unable to achieve the critical conditions to create a thermal QCD medium. But over the decades, a significant increase in the center of mass collision energies at the LHC has been achieved and results from $p + p$ collisions have changed this perception. The data from high multiplicity $p + p$ collisions at $\sqrt{s} = 7$ and 13 TeV have shown the phenomena which resemble the conditions of heavy ion collision^{[13][14][15][16]}. However, the existence of QGP in such a small system is still unclear and requires more investigation. In this direction, experimental data of the normalized charmonium yield observed in $p + p$ collisions have been quantitatively explained using the Unified Model of Quarkonia Suppression (UMQS)^[8]. The UMQS model is based on the QGP phenomenology and it was successful in explaining the quarkonia suppression in A–A and p –A collisions at various center of mass energies. And in the given conditions it predicts a QGP-like scenario in the $p + p$ collisions at the LHC energy.

The quarkonia suppression in heavy-ion collisions due to the production of QGP was first proposed by Matsui and Satz^[1] based on the color screening mechanism. In their work, it was considered that when the temperature of the medium exceeds the dissociation/melting temperature of the quarkonium, the quark-antiquark potential gets screened, resulting in the suppression of these states. It infers that quarkonia have adequate time to adjust to the evolving medium, thereby undergoing adiabatic evolution. The adiabatic evolution involves gradual changes, enabling the system to adapt its

configuration over time. As a consequence, the change in the Hamiltonian of the system must occur slowly to prevent transitions to different eigenstates. However, such conditions may not be satisfied in small collision systems like $p + p$, where the temperature of the fireball is extremely high but the system size is very small. Such a system is expected to cool down rapidly and consequently, the rapid evolution of the plasma dynamics can challenge the conditions required for adiabatic evolution. This necessitates a theoretical framework that accounts for non-adiabatic evolution, where quarkonia states make a transition to other bound states or continuum states due to rapid changes in temperature^{[17][18][19][20][21][22]}. A similar scenario may exist in non-central heavy-ion collisions, where a transient magnetic field can contribute to the non-adiabatic evolution of medium viz quarkonia^{[23][24][25][26]}. As the evolution of quarkonia depends on the QGP lifetime, particularly the temperature decay rate and the initial temperature of the medium, a rapid decrease in the temperature may not allow sufficient time for quarkonia to dissociate, even if the initial temperature surpasses the dissociation threshold. By extending the concept of adiabatic evolution, it is argued that the effective temperature determines the fate of quark-antiquark bound states in $p + p$ collisions. If the effective temperature exceeds the dissociation temperature, bound states dissolve; otherwise, dissociation is minimal. As discussed, in the $p + p$ collisions, the rapid temperature reduction can abbreviate the lifespan of the deconfined QCD medium, leading to abrupt alterations in the Hamiltonian of the quarkonia. Consequently, it permits non-adiabatic evolution to take place.

In this study, we consider the Gubser-like expansion of the medium created in ultra-relativistic $p + p$ collision, which predicts that the thermalized medium gets exhausted in a very brief time, say ≤ 1 fm. It allows us to delve into the suppression of charmonium by incorporating non-adiabatic evolution, showcasing how it can extend the persistence of quark-antiquark bound states even amidst heightened multiplicities^[27]. Following this, we study the yield modification of the J/ψ , χ_c and $\psi(2S)$ in $p + p$ collision via incorporating non-adiabatic evolution of quarkonia along with the collision damping, gluonic dissociation, and regeneration mechanisms as the QGP effects^[4].

The paper is organized as follows: In Section II, we discuss the dynamics of the fireball by modeling temperature evolution in the pre-hydrodynamic or pre-equilibrium phase followed by Gubser flow for the thermalized/hydrodynamic phase. In this section, we also discuss the modification in the temperature in the particle rest frame caused by the Relativistic Doppler Shift (RDS). Next, section III, incorporates the dissociation probability of quark-antiquark bound states as well as transitions to other states under non-adiabatic evolution using time-dependent perturbation theory. Further, it

briefly describes the regeneration of charmonium during QGP evolution. It takes us to the next Section IV, which presents the main outcomes of the study, demonstrating the yield modification of different charmonium states against charged particle multiplicity ($\frac{dN_{ch}}{d\eta}$) and transverse momentum (p_T). We also observe the yield modification of the χ_c and $\psi(2S)$ with respect to J/ψ in terms of double ratio as well as number ratios. Finally, Section V concludes and summarizes the results, providing an outlook on future research.

II. System Dynamics

This section is divided into three main parts. Firstly, we will discuss the solution of the time-dependent Schrödinger equation, focusing on the effects of rapid changes in potential on bound states. Subsequently, we explore the evolution of temperature during the pre-equilibrium stage. Lastly, we analyze the temperature evolution in the late stage after hadronization or during the near-equilibrium stage.

A. Pre-equilibrium Kinematics

The key quantity that controls the evolution of wave function and specifically modifies the survival probability is the Hamiltonian, which carries the temporal dependence originating from the time dependence of the temperature. Modeling the time evolution of temperature is non-trivial for the entire evolution of the plasma in heavy ion collisions. Fortunately, hydrodynamic evolution plays a crucial role in the space-time evolution of the QCD medium after the partonic medium thermalizes. The hydrodynamics unambiguously describes the bulk evolution of the medium. Therefore, we can choose a hydrodynamic model to study the temperature evolution which governs the evolution of the Hamiltonian. However, to model the pre-equilibrium stages, one may rely on the effective QCD kinetic theory description as discussed within the framework of bottom-up thermalization^{[28][29][30]}. Qualitatively, in this approach, it has been argued that in non-expanding systems, gauge bosons can rapidly achieve equilibrium (kinetic) among themselves, followed by the equilibration of the fermions. On the other hand, if the system undergoes rapid longitudinal expansion, partons may remain out of equilibrium, but the system can be effectively described by fluid dynamics. Without going into the details of the model, we consider the following ansatz for the proper time evolution of the pseudo temperature (T_p). Note that it is a pseudo-temperature because the temperature is strictly defined only in equilibrium:

$$\frac{T_p}{T_{Hydro}} = \left(\frac{\tau}{\tau_{Hydro}} \right)^{\frac{1}{7} \frac{\alpha-1}{(\alpha+3)}} \quad (1)$$

The physical picture that prompts us to explore the above scaling is that the initial out-of-equilibrium partons scatter with each other to achieve kinetic/thermal equilibrium. We also identify the thermalization time scale as the time when we can apply the hydrodynamic description (τ_{Hydro}). In principle, all these different time scales can form a hierarchy, but we expect that if thermalization is achieved very fast, then the difference between different scales may not be too large, not affecting the system dynamics significantly. The parameter α enters the above equation because the pseudo temperature can be defined through the α -th moment of the fermionic or bosonic distribution function^[28]. Physically, the parameter α determines how fast the system achieves hadronization (onset of hydrodynamic description) or thermalization. In the subsequent discussion, we appropriately choose T_{Hydro} , τ_{Hydro} , and α to model the pre-equilibrium dynamics. Instead of going into the microscopic description, naively, one can also assume the temperature starts at zero at some initial time and increases linearly until it reaches a value T_{Hydro} at time τ_{Hydro} .

B. Thermal Evolution with Gubser flow

Once we have a description of the pre-equilibrium pseudo temperature that also quantifies the pre-equilibrium dynamics of the Hamiltonian, we can look into the temperature evolution due to the flow dynamics. To solve the hydrodynamic equations, defining both the initial conditions and the equation of state (EoS) is essential. In the absence of a first-principle method for estimating the initial temperature ($T_0 = T_{Hydro}$), the following relation has been employed in this study to constrain T_0 using available data^[31];

$$T_0 = \left[\frac{90}{g_k 4\pi^2} C' \frac{1}{A_T \tau_0} 1.5 \frac{dN_{ch}}{dy} \right]^{1/3} \quad (2)$$

where A_T is the transverse area of the system obtained using the IP-Glasma model^[32], and g_k is the statistical degeneracy of the QGP phase. In Eq. 2, $C' = \frac{2\pi^4}{45\zeta(3)} \approx 3.6$. Additionally, we assume $\frac{dN_{ch}}{dy} \cong \frac{dN_{ch}}{d\eta}$, which holds true in the massless limit. Given the lack of a first-principle approach to determine thermalization time ($\tau_0 = \tau_{Hydro}$), it is reasonable to hypothesize that the thermalization time decreases with increasing center-of-mass collision energy, i.e., $\tau_0 \propto 1/\sqrt{s}$ ^{[33][31]}. In this study, we assume $\tau_0 = 0.1$ fm for pp collisions at $\sqrt{s} = 13$ TeV.

Notably, in $p + p$ collisions, the size of the produced medium is expected to be relatively small, and the maximum size for high multiplicity $p + p$ collision can approximately be 1.5 fm^[32]. Consequently, transverse expansion must be addressed, unlike in the Bjorken flow solution for large systems. To account for the transverse expansion of the system in this calculation, we examine the Gubser flow, first explored by Gubser and Yarom^{[34][35]}. This approach combines a “boost-invariant” longitudinal flow, akin to the Bjorken flow, with consideration for transverse flow. The evolution of thermodynamic quantities, including energy density (ϵ) and shear stress (π), within the framework of Gubser flow with third-order viscous corrections, is detailed in^{[36][37]}.

$$\frac{d\hat{\epsilon}}{d\rho} = -\left(\frac{8}{3}\hat{\epsilon} - \hat{\pi}\right)\tanh(\rho) \quad (3)$$

$$\frac{d\hat{\pi}}{d\rho} = -\frac{\hat{\pi}}{\hat{\tau}_\pi} + \tanh(\rho)\left(\frac{4}{3}\hat{\beta}_\pi - \hat{\lambda}\hat{\pi} - \hat{\chi}\frac{\hat{\pi}^2}{\hat{\beta}_\pi}\right) \quad (4)$$

The dimensionless quantities, $\hat{\epsilon}$ and $\hat{\pi}$, are expressed as $\hat{\epsilon} = \hat{T}^4 = \epsilon\tau^4 = 3\hat{P}$ and $\hat{\pi} = \pi\tau^4$ where τ is the proper time and $\hat{T}$ is related to temperature. The parameters are chosen^[36] as $\epsilon = \frac{3}{\pi^2}T^4$, $\hat{\tau}_\pi (= c/\hat{T})$ is related to relaxation time, where $c = 5\frac{\eta}{s}$, $\hat{\beta}_\pi = 4\hat{P}/5$, $\hat{\lambda} = 46/21$ and the third order correction parameter $\hat{\chi} = 72/245$.

The conformal time ρ can be written as

$$\rho = -\sinh^{-1}\left(\frac{1 - q^2\tau^2 + q^2x_T^2}{2q\tau}\right) \quad (5)$$

where q is an arbitrary energy scale, which is related to the transverse size of the medium (r_T) like $q = 1/r_T$, x_T is the position in the transverse plane. One can retrieve the Bjorken flow solution by taking the limit $r_T \rightarrow \infty$ or $q \rightarrow 0$. One can also use the $(3 + 1)$ dimensional hydrodynamic description for a more accurate description of non-boost invariant flow with nontrivial rapidity dependence. But considering the possible boost invariance in ultra-relativistic collisions, we restrict ourselves to the analytically solvable hydrodynamic description with transverse expansion.

Previously, it was demonstrated in Ref. ^[27] that the effect of the transverse expansion on the temperature evolution by Eqs.(3) and (4) with initial conditions $T = T_{Hydro} = 350$ MeV and $\hat{\pi} = \frac{4}{3}\hat{\beta}_\pi\hat{\tau}_\pi$ at $\tau = \tau_{Hydro} = 0.3$ fm for various system sizes (r_T). The results indicate that, as r_T increases, the lifetime of the Quark-Gluon Plasma (QGP) increases. At sufficiently large r_T , the variation of temperature (T) with proper time (τ) for Gubser flow closely resembles that of Bjorken flow. The variation of temperature with system size clearly indicates that for small systems, the time

evolution of the system can be rapid as compared to large systems, allowing us to explore the scenario of non-adiabatic evolution.

C. In-medium Implicit Temperature for Quarkonia

Heavy quarkonia like charmonia do not attain thermalization with the medium, and such mesons have a different velocity than the medium. The velocities of the medium and charmonium are denoted by v_m and $v_{J/\psi(nl)}$, respectively. This lack of integration between charmonium and the surrounding medium gives rise to acquiring an effective temperature for charmonium. The effective temperature of charmonium in the medium is obtained using the Relativistic Doppler Shift (RDS), which arises from the velocity difference between the charmonium and the thermalized QCD medium. The RDS leads to an angle-dependent effective temperature (T_{eff}), expressed as^[38]:

$$T_{\text{eff}}(\theta, |v_r|) = \frac{T(\tau, b) \sqrt{1 - |v_r|^2}}{1 - |v_r| \cos \theta} \quad (6)$$

where θ is the angle between the relative velocity v_r and the direction of the free-flowing light partons. The $T(\tau)$ in Eq. 6, represents the medium cooling rate obtained using Gubser flow. Now we have averaged T_{eff} over the solid angle, which leads to an angle-independent effective temperature, given as^[6],

$$T_{\text{eff}}(\tau, p_T) = T(\tau) \frac{\sqrt{1 - |v_r|^2}}{2|v_r|} \ln \left[\frac{1 + |v_r|}{1 - |v_r|} \right] \quad (7)$$

Using Eq. 7, we have obtained the effective temperature for J/ψ , χ_c and $\psi(2S)$ charmonium states and its variation with τ along comparison with medium temperature T_{medium} is shown Fig. 1. For $p_T \leq 3$ GeV, all the charmonia resonances are found to be thermalized with medium, as T_{eff} corresponding to J/ψ , χ_c and $\psi(2S)$ is almost same as T_{medium} . While T_{eff} obtained for $3 < p_T \leq 9$ GeV comes out less than T_{medium} , following the argument that quarkonia moving with high p_T is incapable of being in thermal equilibrium with the medium. As $\psi(2S)$ and χ_c masses are higher than J/ψ , traverse through the medium with relatively slower speed, and consequently feel slightly higher temperatures at given $p_T > 3$ GeV. This mass ordering on the T_{eff} for J/ψ , χ_c and $\psi(2S)$ is preserved at $9 < p_T \leq 30$ GeV while T_{eff} is further reduced respective to the p_T -range; $3 < p_T \leq 9$. Our findings suggest that T_{eff} plays a crucial role in the modification of charmonium yield in the QGP medium.

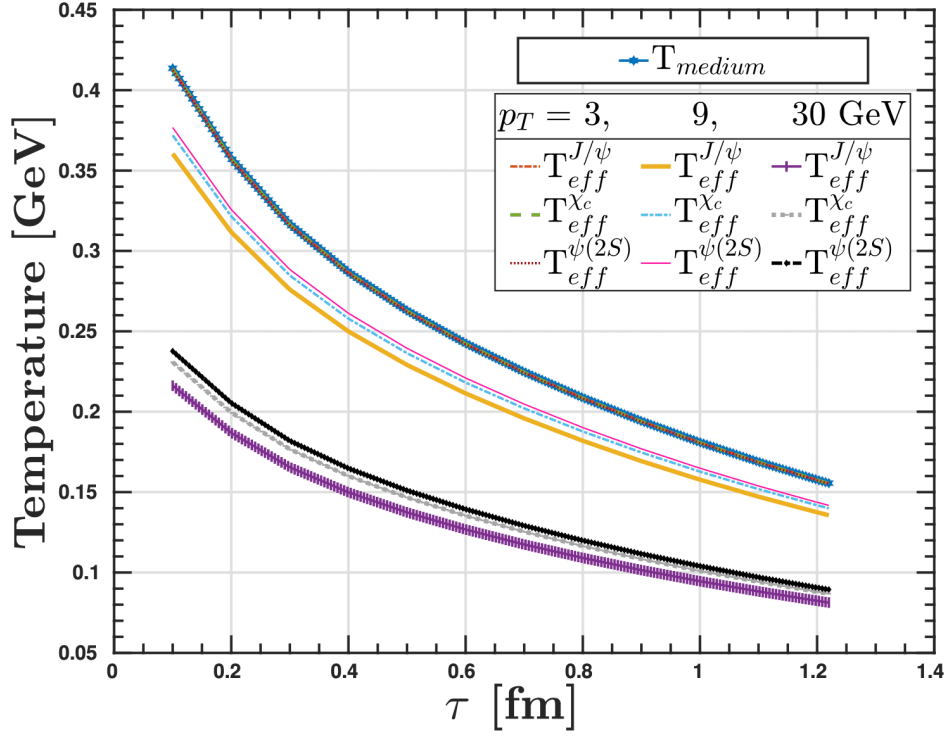


Figure 1. (Color online) Medium Temperature evolution with time (τ) corresponding to Gubser flow is represented through T_{medium} . Along with it, the effective temperature (T_{eff}) respective J/ψ , $\chi_c(1P)$ and $\psi(2S)$ for $p_T = 3, 9, 30$ GeV are also shown here.

III. Yield Modification Mechanisms

There are several phenomena related to cold nuclear matter and hot QCD matter that may influence the production/suppression of charmonium in heavy-ion collisions. However, unlike heavy-ion collisions, the nuclear environment is absent in p+p collisions. Therefore, any changes in charmonium yield in these collisions can be attributed solely to the effects of a hot partonic medium. In the context of hot QCD matter effects, we have considered factors such as collisional damping, gluonic dissociation, and the regeneration of charmonium states through gluonic de-excitation. Additionally, we have formulated the non-adiabatic approximation for charmonium evolution and studied its impact on charmonium yield.

A. Collisional Damping

The collisional energy loss of charmonium within the quark-gluon plasma (QGP) is characterized by “Collisional Damping.” The charmonium dissociation due to this damping effect is evaluated using the complex singlet potential. In this study the singlet potential for the $c\bar{c}$ bound state in the QGP medium is defined as follows [39][40][41]:

$$V(r, m_D) = \frac{\sigma}{m_D} (1 - e^{-m_D r}) - \alpha_{\text{eff}} \left(m_D + \frac{e^{-m_D r}}{r} \right) - i\alpha_{\text{eff}} T_{\text{eff}} \int_0^\infty \frac{2zdz}{(1+z^2)^2} \left(1 - \frac{\sin(m_D r z)}{m_D r z} \right) \quad (8)$$

In Eq. (8), the first two terms on the right-hand side represent the string and Coulombic contributions, respectively. The third term is the imaginary component of the heavy-quark potential, which accounts for collisional damping. Here, σ is the string tension for the $c\bar{c}$ bound state, and m_D denotes the Debye mass. The running coupling constant at the hard scale, α_s^T ; $\alpha_s^T = \alpha_s(2\pi T)$. The effective coupling constant, α_{eff} , is defined at the soft scale as $\alpha_s^s = \alpha_s$, and is expressed as $\alpha_{\text{eff}} = \frac{4}{3} \alpha_s^s$. The collisional damping, $\Gamma_{c,nl}$, describes the decay of charmonium caused by the imaginary part of the complex potential. It is computed using first-order perturbation theory by integrating the imaginary part of the potential with the radial wave function:

$$\Gamma_{c,nl}(\tau, p_T) = \int [g_{nl}(r)]^\dagger [Im(V)] g_{nl}(r) dr, \quad (9)$$

where $g_{nl}(r)$ is the charmonia singlet wave function. We have obtained the wave functions by solving the Schrödinger equation for J/ψ , $\chi_c(1P)$, $\psi(2S)$.

B. Glutonic Dissociation

In the quark-gluon plasma (QGP) medium, quarkonium states can make a transition from a color singlet state to a color octet state, which gradually dissociates within the medium. The cross-section for this process is given by [8]:

$$\sigma_{d,nl}(E_g) = \frac{\pi^2 \alpha_s^u E_g}{N_c^2} \sqrt{\frac{m_c}{E_g + E_{nl}}} \times \left(\frac{l |J_{nl}^{q,l-1}|^2 + (l+1) |J_{nl}^{q,l+1}|^2}{2l+1} \right) \quad (10)$$

where m_c is the mass of the charm quark and α_s^u is the coupling constant, scaled as $\alpha_s^u = \alpha_s(\alpha_s m_c^2/2)$. The E_{nl} is the energy eigenvalue corresponding to the charmonium wave function, $g_{nl}(r)$. Here, $J_{nl}^{q,l}$ is

the probability density, derived using both the singlet and octet wave functions as follows;

$$J_{nl}^{q'} = \int_0^\infty dr r g_{nl}^*(r) h_{q'}(r) \quad (11)$$

The octet wave function $h_{q'}(r)$ has been obtained by solving the Schrödinger equation with the octet potential $V_8 = \alpha_{\text{eff}}/8r$. The value of q is determined by using the conservation of energy, $q = \sqrt{m_c(E_g + E_{nl})}$.

The gluonic dissociation rate, $\Gamma_{gd,nl}$, is calculated by taking the thermal average of the dissociation cross-section [8]:

$$\Gamma_{gd,nl}(\tau, p_T, b) = \frac{g_d}{4\pi^2} \int_0^\infty \int_0^\pi \frac{dp_g d\theta \sin \theta p_g^2 \sigma_{d,nl}(E_g)}{e^{\left\{ \frac{\gamma E_g}{T_{\text{eff}}} (1 + v_{J/\psi} \cos \theta) \right\}} - 1} \quad (12)$$

The p_T is the transverse momentum of the charmonium, and g_d represents the degeneracy factor of the gluons.

Now taking the sum of the decay rates associated with collisional damping and gluonic dissociation, the combined effect is expressed in terms of the total decay width, given as;

$$\Gamma_{D,nl}(\tau, p_T) = \Gamma_{c,nl} + \Gamma_{gd,nl} \quad (13)$$

C. Time-dependent Schrödinger equation: non-adiabatic evolution of quantum states

First, to study the non-adiabatic behavior of charmonium states, it is essential to ensure that the evolution rate exceeds the transition time. We define the evolution time scale, τ_{ev} , as $T \times \frac{d\tau}{dT}$, while the transition time scale, τ_{tr} , represents the time associated with transitions between different energy states, specifically given by $1/\Delta E$ fm. Using this, we estimate τ_{ev} to be approximately 0.3 fm during the thermalization phase, where the temperature evolution follows a Gubser-type profile with viscous correction. This estimation assumes an initial system size of ~ 1.5 fm, which is relevant to high-multiplicity $p + p$ collisions. In contrast, the transition time scale τ_{tr} for charmonium states is calculated to be around 0.7 fm.

Now, coming back to the charmonium evolution; the charmonia are hypothesized to form during the initial phases of collision. Utilizing a bottom-up thermalization approach rooted in QCD kinetic theory, it can be argued that from an initially interacting out-of-equilibrium state, a thermalized medium is reached at a subsequent time marked as τ_{Hydro} . Right after the collision (well before

thermalization). The evolution of the initial state of charmonia from $\tau = 0$ to τ_{Hydro} , can be determined by solving the *zero-temperature* Hamiltonian [41],

$$H_0 = \frac{\vec{p}^2}{2M} + \sigma r - \frac{4}{3} \frac{\alpha_s}{r} \quad (14)$$

Here, M denotes the reduced mass of the quark-antiquark system. However, as thermalization occurs in the medium, the zero-temperature Hamiltonian evolves into its finite temperature counterpart [42],

$$H(\tau) = \frac{\vec{p}^2}{2M} + \text{Re}(V) \quad (15)$$

here $\text{Re}(V)$ is the real part of the potential given in Eq. 8.

The time-dependent nature of the Hamiltonian arises from the temporal variation of temperature. As the system expands further, the medium temperature eventually drops below the threshold for hadronization, causing the Hamiltonian to revert to a zero-temperature state. It has been contended in [27] that the adiabatic approximation for the evolution of the quantum bound state of charmonia may not hold, as the Hamiltonian evolves quite rapidly for medium produce in $p + p$ collision.

Due to the rapid evolution of the medium, the initial quarkonia states experience non-adiabatic evolution, which may cause transitions to states orthogonal to their initial configurations. Consequently, the survival probability of the initial state is affected. To determine this survival probability, we solve the time-dependent Schrödinger equation, which can be expressed as follows:

$$-\frac{1}{2M} \nabla^2 \psi + V(r)\psi = i \frac{\partial \psi}{\partial \tau} \quad (16)$$

In spherical polar coordinate ∇^2 can be written as,

$$\nabla^2 = \frac{1}{r^2} \left[\frac{\partial}{\partial r} \left(r^2 \frac{\partial}{\partial r} \right) + \frac{1}{\sin \theta} \frac{\partial}{\partial \theta} \left(\sin \theta \frac{\partial}{\partial \theta} \right) + \frac{1}{\sin^2 \theta} \frac{\partial^2}{\partial \phi^2} \right] \quad (17)$$

and $V(r) = \frac{\sigma}{\mu} (1 - \exp(-\mu r)) - \frac{4}{3} \alpha_s \exp(-\mu r)/r$. As for simplicity, we are considering potential is spherically symmetric, one can write $\psi = R(r)\psi(\theta, \phi)$ and separate the radial part of Schrödinger equation as

$$\left[-\frac{1}{2M} \frac{d^2}{dr^2} + V_{\text{eff}}(r) \right] u(r) = E u(r) \quad (18)$$

where $V_{\text{eff}}(r) = \frac{\hbar^2}{2mr^2} l(l+1) + V(r)$, $u(r) = rR(r)$ and E is energy Eigenvalue, represents the binding energy.

Here we solve the time-dependent Schrödinger equation (Eq. 18) for the time-dependent Hamiltonian shown in Eq. 15 using the Crank-Nicolson Method ([43]) to obtain the survival probability of a particular initial state. The initial states, charmonia bound states (J/ψ , $\psi(2S)$, and $\chi_c(1P)$), evolves with time until the temperature drops below the QGP threshold temperature T_c . The survival probability of particular charmonium states can be calculated by taking the overlap integration of the final wave function with the initial zero-temperature charmonium state.

If we consider $|J/\psi\rangle$, $|\psi(2S)\rangle$, and $|\chi_c(1P)\rangle$ states to represent bound states of the initial zero-temperature Hamiltonian (Eq. 14), and $\psi(\tau)$ represents the evolving wave function, the survival probability of $|J/\psi\rangle$, $|\psi(2S)\rangle$, and $|\chi_c(1P)\rangle$ at $\tau = \tau_c$ can be represented as:

$$P_{J/\psi} = |\langle\psi(\tau_c)|J/\psi\rangle|^2 \quad (19)$$

$$P_{\psi(2S)} = |\langle\psi(\tau_c)|\psi(2S)\rangle|^2 \quad (20)$$

$$P_{\chi_c(1P)} = |\langle\psi(\tau_c)|\chi_c(1P)\rangle|^2 \quad (21)$$

D. The Regeneration Factor

In addition to the gluonic excitation of a color-neutral state to a color-octet state, gluonic de-excitation from the color-octet to the neutral state is also feasible. Consequently, charmonia gets regenerated in the QGP medium through this process. The regeneration is significant in heavy-ion collisions at LHC energies due to the abundant production of $c\bar{c}$ pairs in the hot QGP medium, which regenerate charmonia through recombination of $c\bar{c}$. While, in smaller systems like $p + p$ collisions, the production of $c\bar{c}$ pairs is relatively low, making the regeneration due to the coalescence less probable. However, regeneration due to gluonic de-excitation plays an important role in estimating charmonium production in such a small collision system (discussed in detail in Ref. [8]). This de-excitation is calculated in terms of the regeneration cross-section $\sigma_{f,nl}$ for charmonium by employing the detailed balance of the gluonic dissociation cross-section $\sigma_{d,nl}$ [6]:

$$\sigma_{f,nl} = \frac{48}{36} \sigma_{d,nl} \frac{(s - M_{nl}^2)^2}{s(s - 4m_c^2)} \quad (22)$$

where s is the Mandelstam variable, related with the center-of-mass energy of $c\bar{c}$ pair, given as; $s = (p_c + p_{\bar{c}})^2$, where p_c and $p_{\bar{c}}$ are four momenta of c and $\bar{c}$, respectively.

Finally, we have obtained the recombination factor Γ_F by taking the thermal average of the product of recombination cross-section and relative velocity v_{rel} between c and $\bar{c}$:

$$\Gamma_{F,nl} = \langle \sigma_{f,nl} v_{rel} \rangle_{p_c}, \quad (23)$$

$$v_{rel} = \sqrt{\frac{(\mathbf{p}_c^\mu p_{\bar{c}\mu})^2 - m_c^4}{p_c^2 p_{\bar{c}}^2 + m_c^2 (p_c^2 + p_{\bar{c}}^2 + m_c^2)}} \quad (24)$$

Since the gluonic dissociation increases with the increase in temperature, it leads to the production of a substantial number of octet states in a high multiplicity events. Such that the de-excitation of $c\bar{c}$ octet states to J/ψ enhances the regeneration of J/ψ in high multiplicity events conferred with relatively low multiplicities.

E. The Quantified Yield

Modification of the charmonium yield in the medium due to collisional damping, gluonic dissociation, and regeneration is obtained by combining all these mechanisms in one transport equation^{[44][45][46][6][8]}:

$$\frac{dN_{J/\psi(nl)}}{d\tau} = \Gamma_{F,nl} N_c N_{\bar{c}} [V(\tau)]^{-1} - \Gamma_{D,nl} N_{J/\psi(nl)} \quad (25)$$

The first term on the right-hand side of Eq. 25 is a gain term, and the second is the loss term. Here, $V(\tau)$ is the dynamic volume of the evolving medium. We assume that initially, the number of charms (N_c) and anti-charm quarks ($N_{\bar{c}}$) are produced in equal numbers, $N_c = N_{\bar{c}} = N_{c\bar{c}}$. The Eq. 25 can be solved analytically under the assumption of $N_{J/\psi(nl)} \ll N_{c\bar{c}}$ at τ_0 ^{[4][6][8]}:

$$N_{J/\psi(nl)}^f(\tau_{QGP}, p_T) = \epsilon_1(\tau_{QGP}, p_T) \left[N_{J/\psi(nl)}^i + N_{c\bar{c}}^2 \int_{\tau_0}^{\tau_{QGP}} \Gamma_{F,nl}(\tau, p_T) [V(\tau) \epsilon_2(\tau, p_T)]^{-1} d\tau \right] \quad (26)$$

Here, $N_{J/\psi(nl)}^f(\tau_{QGP}, p_T)$ is the net number of charmonium formed during the QGP evolution period τ_{QGP} . The quantities $N_{J/\psi(nl)}^i$ and $N_{c\bar{c}}$ represent the number of J/ψ and $c\bar{c}$ pairs formed during the initial hard scattering, respectively. Input for $N_{J/\psi(nl)}^i$ and $N_{c\bar{c}}$ are taken from Ref. ^[8] corresponding to p-p collisions at $\sqrt{s} = 13$ TeV.

In Eq. 26, $\epsilon_1(\tau_{QGP})$ and $\epsilon_2(\tau)$ are decay factors for the meson due to gluonic dissociation and collisional damping in the QGP with a lifetime of τ_{QGP} , and τ represents the evolution time. These factors are calculated using the following expressions:

$$\epsilon_1(\tau_{QGP}, p_T) = \exp \left[- \int_{\tau_{nl}}^{\tau_{QGP}} \Gamma_{D,nl}(\tau, p_T) d\tau \right], \quad (27)$$

and

$$\epsilon_2(\tau, p_T) = \exp \left[- \int_{\tau_{nl}}^{\tau} \Gamma_{D,nl}(\tau', p_T) d\tau' \right] \quad (28)$$

After obtaining the $N_{J/\psi(nl)}^f(\tau_{QGP}, p_T)$ from Eq. 26 we took the ratio to the initially produced charmonium, $N_{J/\psi(nl)}^i$ to quantifying the medium effect and called it survival probability S_P . The survival probability of charmonium, due to gluonic dissociation and collisional damping along with the regeneration effect is defined as:

$$S_{cgr}^{J/\psi}(p_T, mc) = \frac{N_{J/\psi(nl)}^f(p_T, mc)}{N_{J/\psi(nl)}^i(mc)} \quad (29)$$

here “ mc ” stands for Multiplicity Class defined as $\langle dN_{ch}/d\eta \rangle$.

It is assumed here that from the initial collision to the QGP end-point (at $\tau = 0$ to $\tau = \tau_{qgp}$), the non-adiabatic evolution of charmonia states is a completely independent process with the other suppression mechanisms in QGP. The net yield in terms of survival probability, $S_P(p_T, mc)$ is expressed as:

$$S_P^{J/\psi}(p_T, mc) = S_{gc}^{J/\psi}(p_T, mc) P_{J/\psi}(p_T, mc). \quad (30)$$

Further, we incorporate the feed-down correction from excited charmonium states, such as $\chi_c(1P)$ and $\psi(2S)$, into J/ψ . To determine the feed-down, we calculate the ratio between the net initial and final numbers of J/ψ . The net initial number is derived by accounting for the feed-down from higher resonances into J/ψ in the absence of the QGP medium. This is expressed as $N_{J/\psi}^{in} = \sum_{J \geq I} C_{IJ} N(J)$, where C_{IJ} represents the branching ratio for the decay of state J into state I . The net final number of J/ψ incorporates medium effects, represented by the survival probability ($S_P(p_T, b)$), along with feed-down: $N_{J/\psi}^{fi} = \sum_{J \geq I} C_{IJ} N(J) S_P(J)$. The overall generalized survival probability, including the feed-down correction, is given as^[6]:

$$S_P(I) = \frac{\sum_{J \geq I} C_{IJ} N(J) S_P(J)}{\sum_{J \geq I} C_{IJ} N(J)}. \quad (31)$$

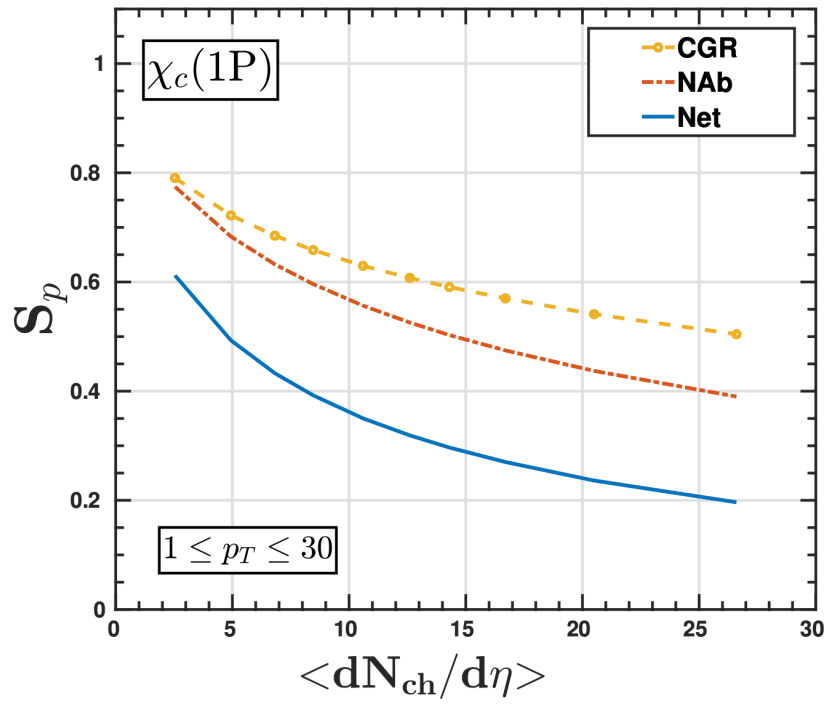
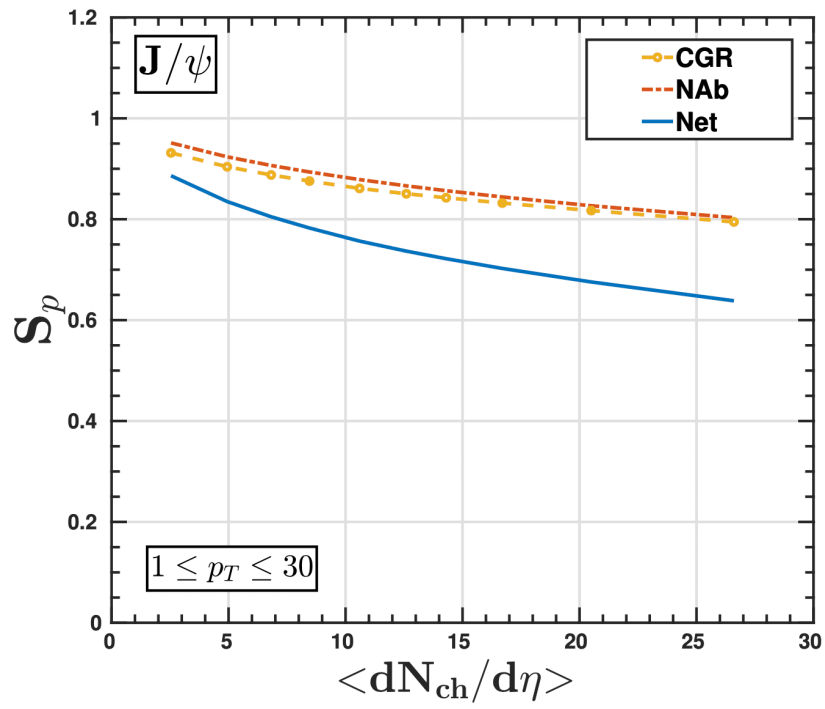
IV. Results and Discussions

In our investigation of charmonium yield modification ultra-relativistic proton-proton ($p + p$) collisions at $\sqrt{s}=13$ TeV under the above-mentioned circumstances, we have obtained the survival probability (S_P), the double ratio (used for the direct comparison of two probabilities), and the variation of particle ratios with respect to both charged particle multiplicity ($\langle dN_{ch}/d\eta \rangle$) and

transverse momentum (p_T) at the mid-rapidity. This study analyzes a spectrum of suppression mechanisms alongside the regeneration process. We have methodically categorized these mechanisms into two distinct groups for clarity and detailed analysis: CGR and NAb. The “CGR” group encapsulates mechanisms such as Collisional damping, Gluonic Dissociation, and Recombination processes, highlighting the interactions that directly involve gluonic exchanges. On the other hand, “NAb” focuses on the non-adiabatic evolution of charmonium states, considering the temporal evolution under the scenario when the reaction time is so short that the transition amplitude is described as the overlap of these states. The combined effects, both CGR and NAb, are presented in a dataset labeled “Net”, showcasing the intertwined relationship and net impact of these complex mechanisms on the charmonium yield in such high-energy collisions.

A. Multiplicity Dependent Yield

The p_T -integrated charmonium yield modification in terms of S_P with event multiplicity at mid-rapidity has been explored using the charmonium distribution function $1/E_T^4$ as discussed in the Ref. [6][8]. The yield modification of J/ψ , χ_c and $\psi(2S)$ shown in Fig. 2, predicts that suppression due to CGR is relatively large for $\chi_c(1P)$ than J/ψ and $\psi(2S)$, and it further increases with increasing multiplicity. The CGR and NAb independently predict about 20% suppression for J/ψ at highest multiplicity, as shown in Fig. 2. However, J/ψ suppression due to NAb is slightly less than CGR at low multiplicity events. Conversely, NAb predicts substantial suppression for $\chi_c(1P)$ compared to CGR. In contrast to J/ψ and χ_c , $\psi(2S)$ experiences an enhancement due to the NAb approach, increasing with multiplicity. Normally, the average radius of $\chi_c(1P)$ is larger than that of J/ψ , and the average radius of $\psi(2S)$ is even larger than both. As the average radius increases, dissociation due to non-adiabatic evolution also increases. Consequently, $\chi_c(1P)$ undergoes more dissociation compared to J/ψ , and $\psi(2S)$ should, in principle, experience even greater dissociation than both J/ψ and $\chi_c(1P)$. Despite this, $\psi(2S)$ exhibits an enhancement due to non-adiabatic transitions from J/ψ to $\psi(2S)$, as described in Ref. [46]. The combined effects of CGR and NAb, represented through “Net”, lead to up to 40% suppression for J/ψ and 80% for $\chi_c(1P)$ with increasing multiplicity. For $\psi(2S)$, the combined effects reduce the enhancement to some extent but are unable to transform it into suppression.



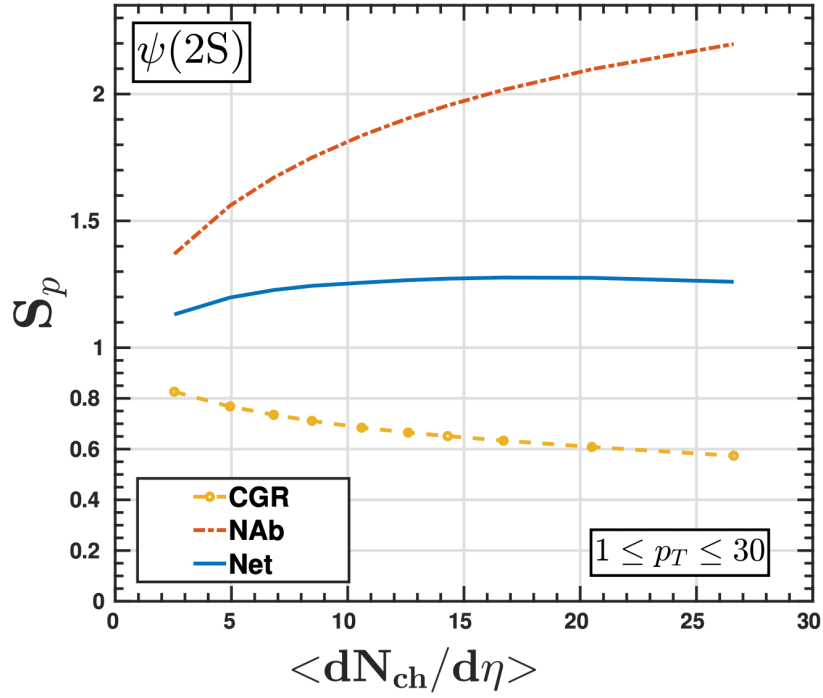


Figure 2. (Color online) Survival probability S_p as a function of multiplicity is shown for J/ψ , $\chi_c(1P)$ and $\psi(2S)$ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV.

In Fig. 3, the feed-down of the $\chi_c(1P)$ and $\psi(2S)$ into J/ψ further increases the suppression for J/ψ at all the multiplicity classes. The feed-down corresponding to the NAb only predicts maximum suppression up to 20% for J/ψ , which is almost the same as its prediction in the absence of the feed-down correction. While considering feed-down correction, only incorporating the CGR process increases the suppression up to 40% while twice as its earlier prediction (CGR without feed-down correction shown in Fig. 2). Finally, the combined effects of CGR and NAb, with feed-down of higher resonances, lead to 50% suppression for J/ψ at high multiplicity events in ultra-relativistic $p + p$ collisions at $\sqrt{s} = 13$ TeV.

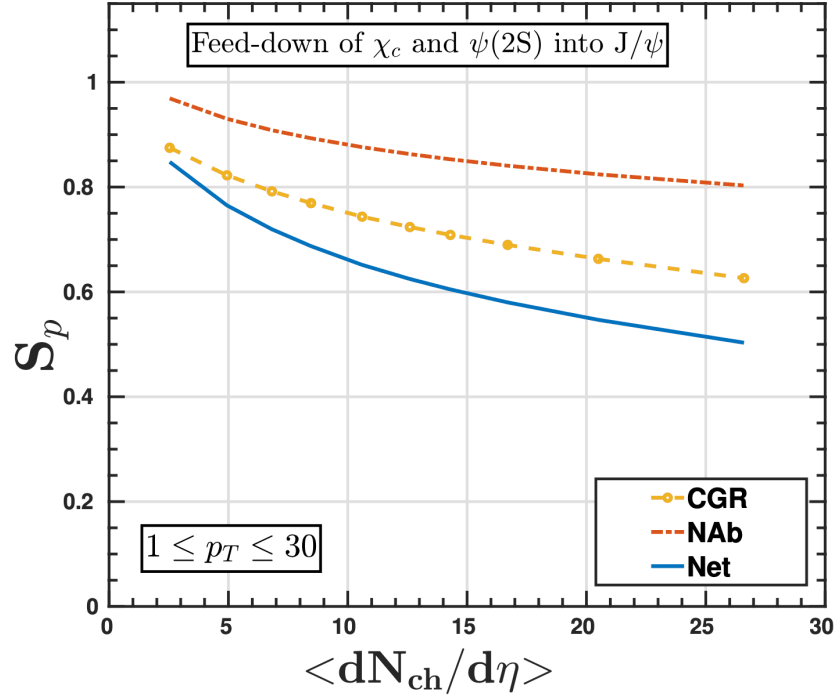
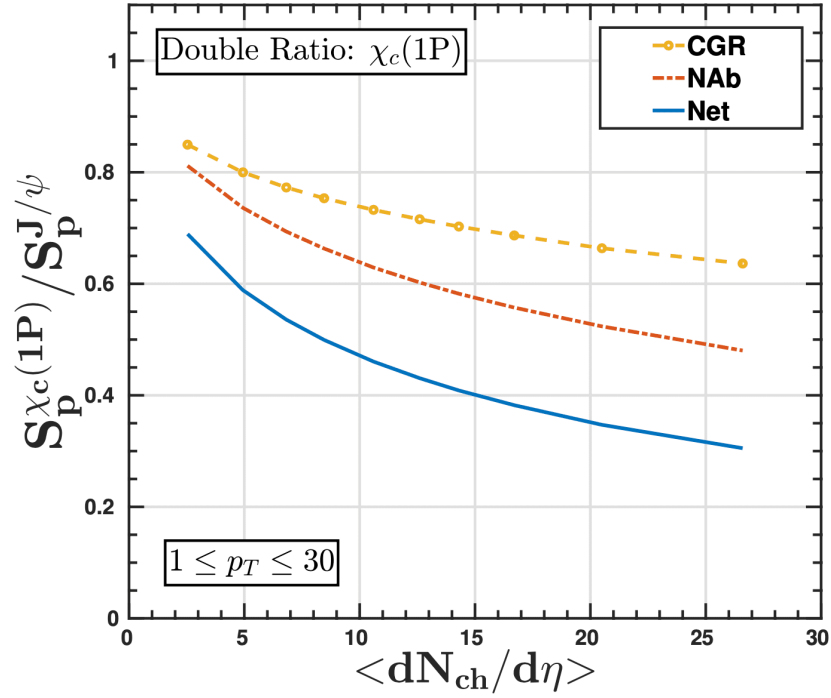


Figure 3. (Color online) Survival probability S_P as a function of multiplicity is shown for J/ψ , considering the feed-down of $\chi_c(1P)$ and $\psi(2S)$ into J/ψ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV.

In Fig. 4 the survival probability ratios or double ratio between $\chi_c(1P)$ to J/ψ and $\psi(2S)$ to J/ψ are shown to quantify the relative yield modification of $\chi_c(1P)$ and $\psi(2S)$ with respect to J/ψ . Experimental observations employ double ratios to ascertain that the medium, which may have existed in ultra-relativistic collisions whether, affects the $\psi(2S)$ and J/ψ yields differentially or the same. Notably, due to the technical difficulties in the observation of $\chi_c(1P)$, the yield modification and double ratio for $\chi_c(1P)$ with J/ψ has not been reported in any of the ultra-relativistic heavy-ion collision experiments. However, the present study explored the J/ψ , $\chi_c(1P)$ and $\psi(2S)$ dynamics in the medium and the impact on their yield imposed by the medium. Fig. 4 depicts that $\chi_c(1P)$ experiences significant suppression compared to J/ψ at high multiplicity, whereas the suppression magnitude for $\chi_c(1P)$ to J/ψ is relatively small at low multiplicity. On the other hand, the yield of $\psi(2S)$ is considerably enhanced compared to J/ψ due to the NAb mechanism, leading to a populated transition to $\psi(2S)$ in the final state, which increases with multiplicity. However, CGR predicts that $\psi(2S)$ is more suppressed than J/ψ but less than χ_c at the high multiplicity classes. While at low multiplicity $\psi(2S)$

and $\chi_c(1P)$ are almost equally suppressed. The cumulative influence of CGR and NAb predicts a multiplicity-dependent 30% to 70% suppression for χ_c compared to J/ψ and similarly estimates an enhancement of approximately 130% to 200% for $\psi(2S)$ relative to J/ψ .



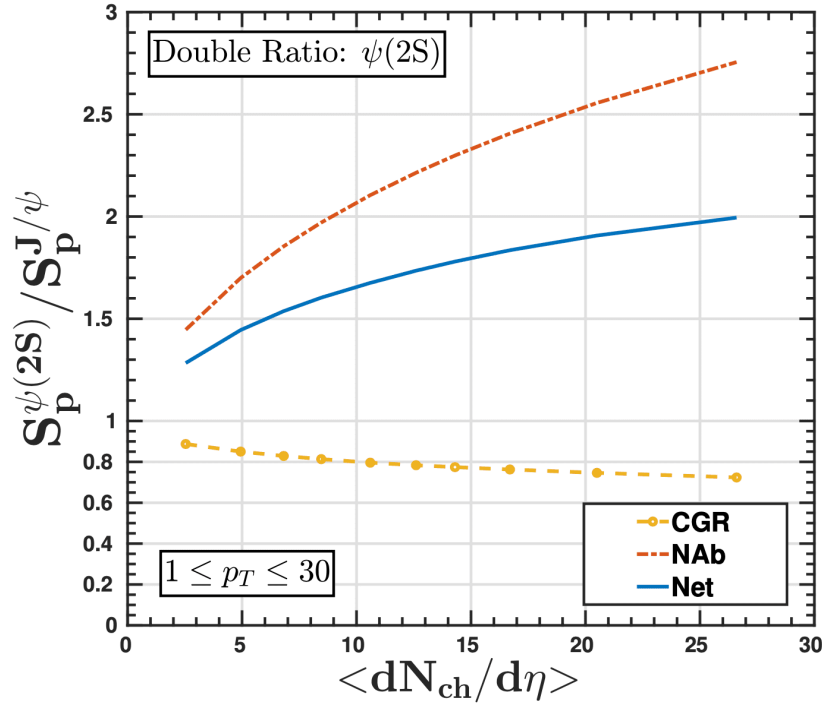


Figure 4. (Color online) Double ratio as a function of multiplicity is shown for $\frac{\chi_c(1P)}{J/\psi}$, and $\frac{\psi(2S)}{J/\psi}$, at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV.

Furthermore, we extend the examination of survival probability ratios of the charmonium states in ultra-relativistic collisions to quantitatively assess the final numbers of J/ψ , $\chi_c(1P)$, and $\psi(2S)$ at the chemical freezeout boundary, as depicted in Fig. 5. These results align with previous observations, indicating a notably higher production and relatively less suppression of J/ψ compared to $\chi_c(1P)$ and $\psi(2S)$ during the transportation from initial production to the QGP endpoint, i.e., $T = T_c$. Additionally, Fig. 5 suggests that $\chi_c(1P)$ dissociation due to the CGR mechanism is relatively smaller in comparison with $\psi(2S)$. While the NAb mechanism effectively reduces the $\chi_c(1P)$ and comparably predicts a large production for $\psi(2S)$. The χ_c yield decreases with increasing multiplicity for both CGR and NAb processes. The $\psi(2S)$ yield is almost steady at all the multiplicities corresponding to CGR while it increases with multiplicity. Meanwhile, the combined effects of CGR and NAb predict the survived production for $\chi_c(1P)$ around 8% to 2% and approximately 0.5% to 1% for $\psi(2S)$ with respect to J/ψ depending on the multiplicity classes.

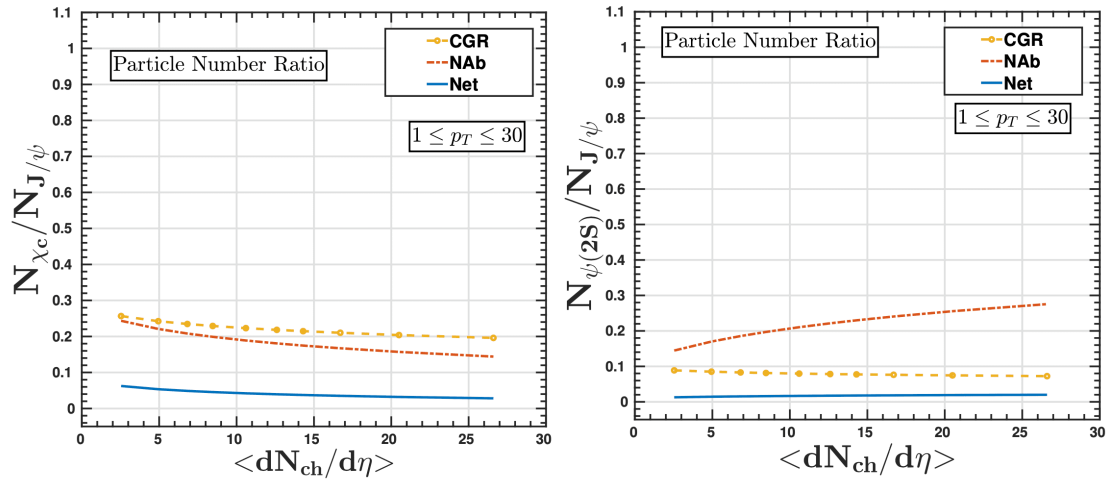


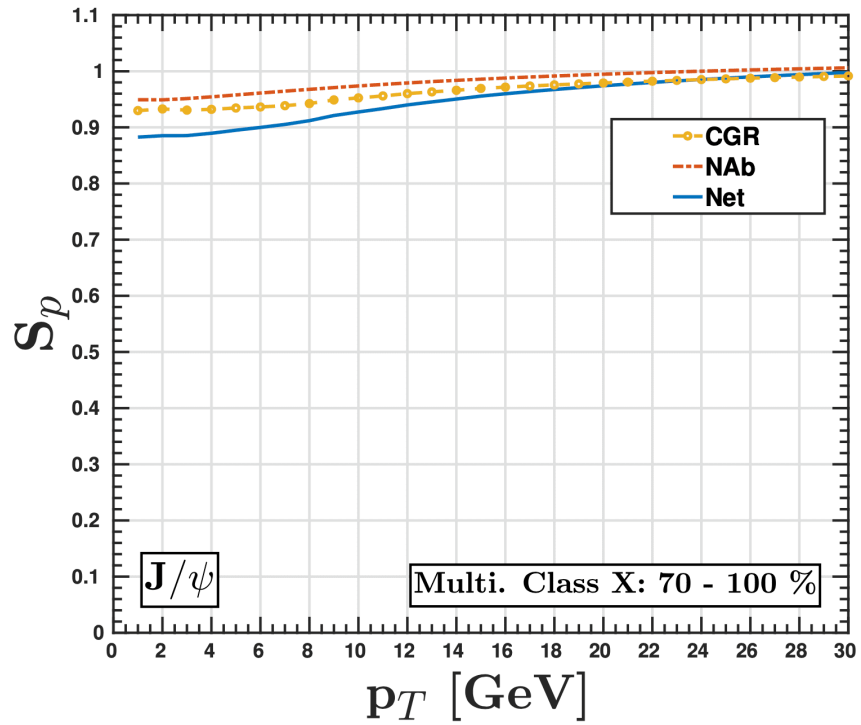
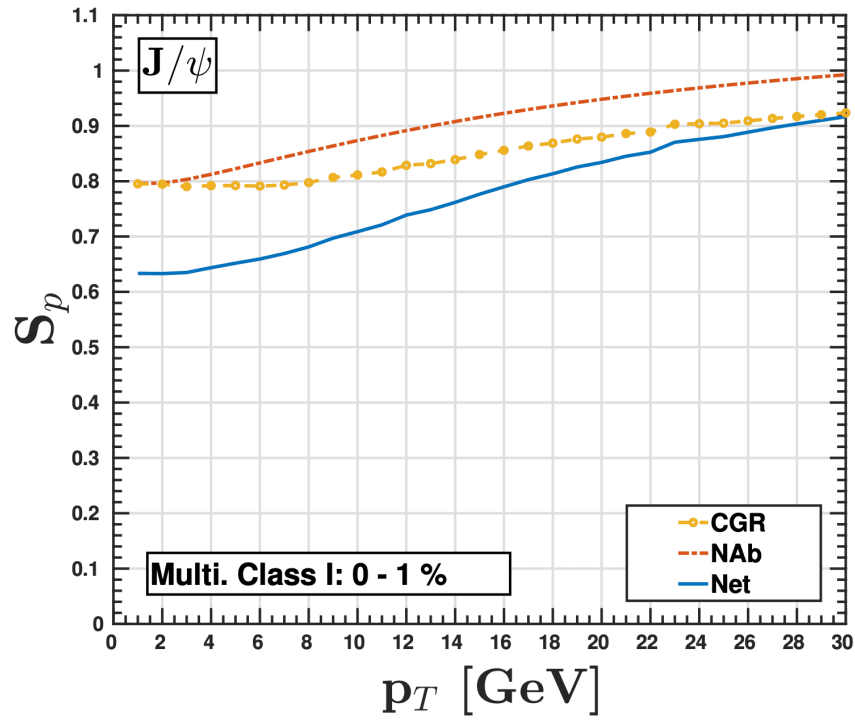
Figure 5. (Color online) Particle number ratio as a function of multiplicity is shown for $\frac{\chi_c(1P)}{J/\psi}$, and $\frac{\psi(2S)}{J/\psi}$ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV.

B. Yield Modification with Transverse Momentum

The production of charmonia as a function of transverse momentum (p_T) provides valuable insights into the physics at both low and high p_T . We have also examined the impact of system size or, in this case, multiplicity classes on the charmonium yield over the considered p_T -range. The chosen multiplicity classes include the lowest multiplicity class (Multi. Class X: 70 - 100%), the highest multiplicity class (Multi. Class I: 0 - 1%), and minimum bias (Min. Bias: 0 - 100%). To this end, we have computed the survival probability (S_P) as a function of p_T by averaging over the range of the corresponding multiplicity bins. The expression for the weighted average of S_P is given by^[6]:

$$S_P(p_T) = \frac{\sum_i S_P(p_T, \langle b_i \rangle) W_i}{\sum_i W_i} \quad (32)$$

The index $i = 1, 2, 3, \dots$ represents the multiplicity bins. The weight function W_i is defined as $W_i = \int_{b_{min}}^{b_{max}} N_{coll}(b) \pi b db$. The number of binary collisions N_{coll} is determined using a Glauber model for $p + p$ collisions, which incorporates an anisotropic and inhomogeneous proton density profile to calculate N_{coll} ^[47]. Also, we have obtained the impact parameter, b , for $p + p$ collisions corresponding to multiplicity bins at $\sqrt{s} = 13$ TeV using the above-mentioned Glauber model.



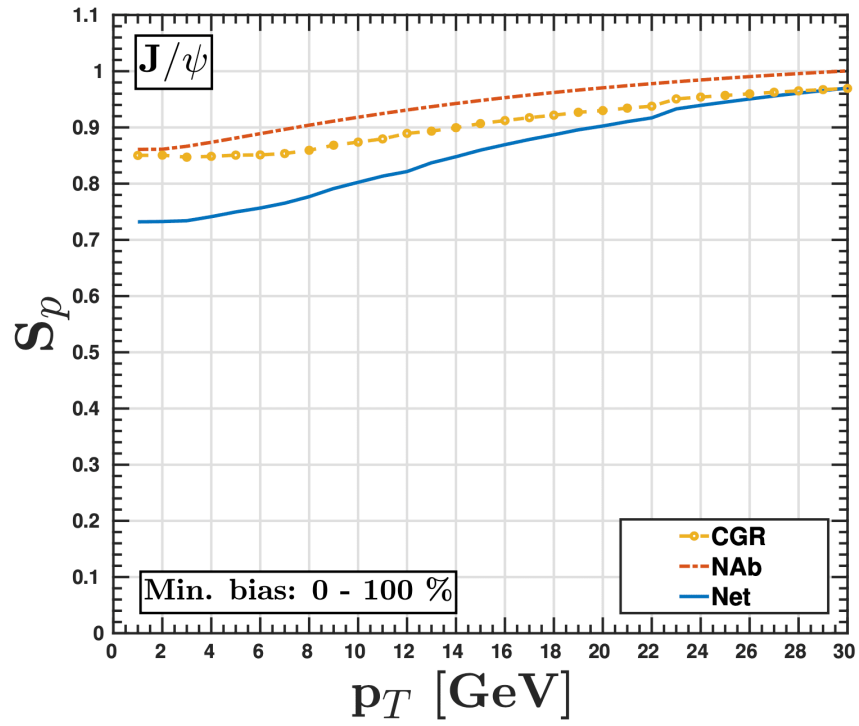
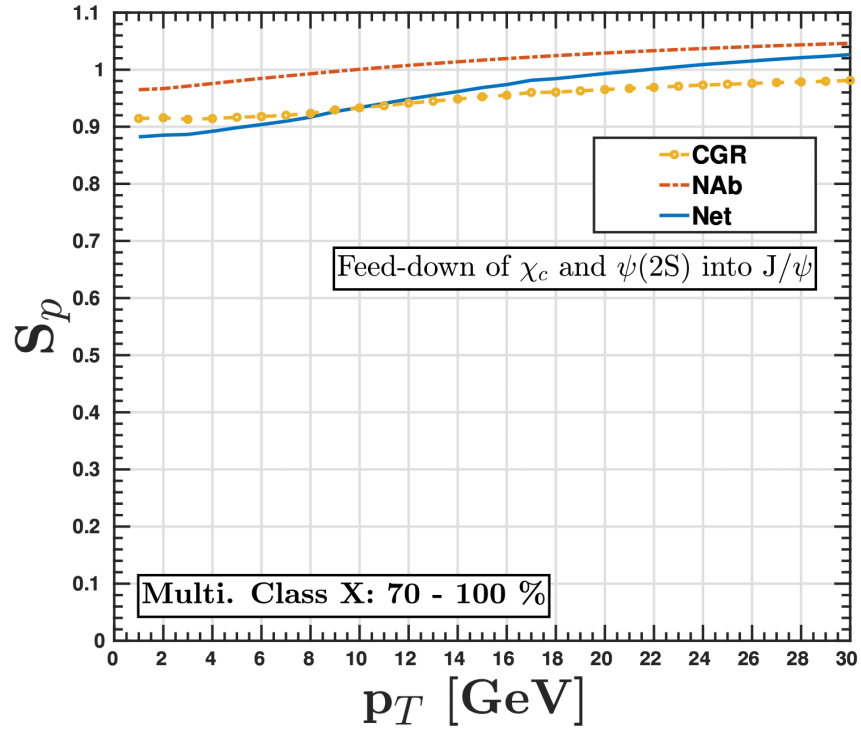
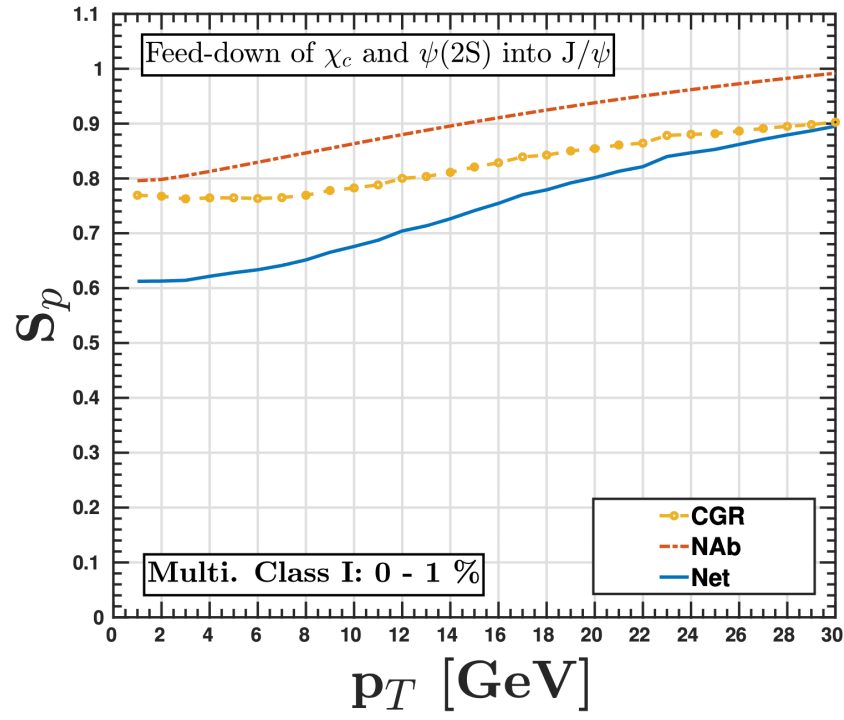


Figure 6. (Color online) Survival probability S_P as a function of p_T is shown for J/ψ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. From top to bottom, results are shown for high-multiplicity, low-multiplicity, and minimum bias events, respectively.



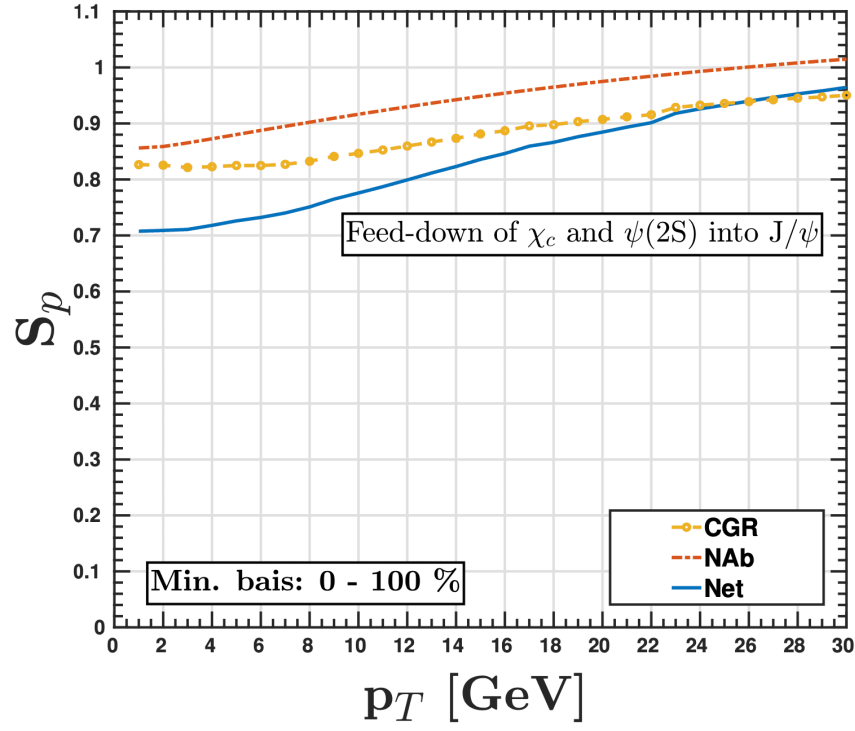
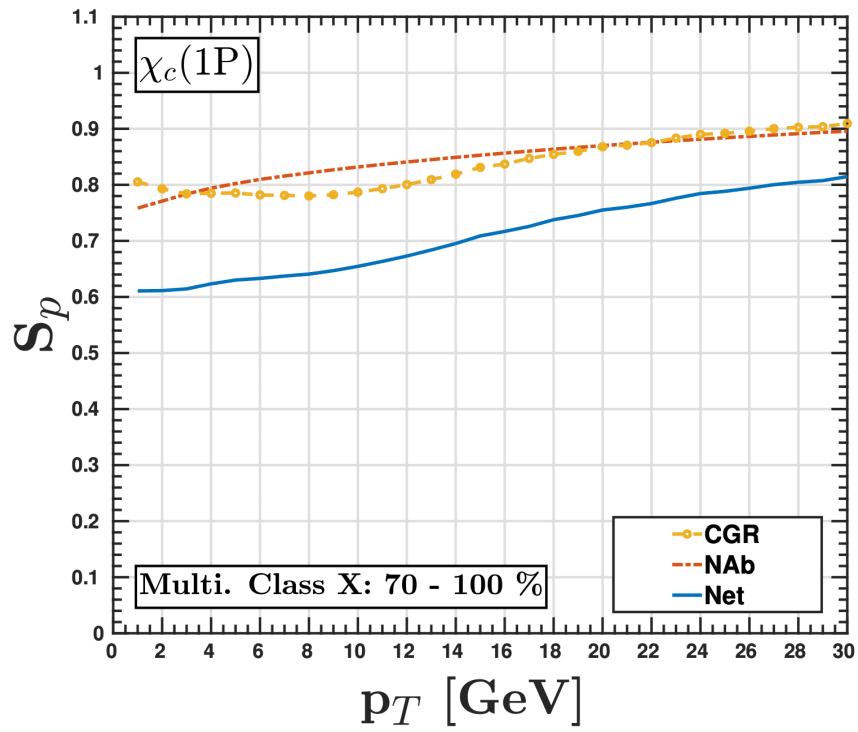
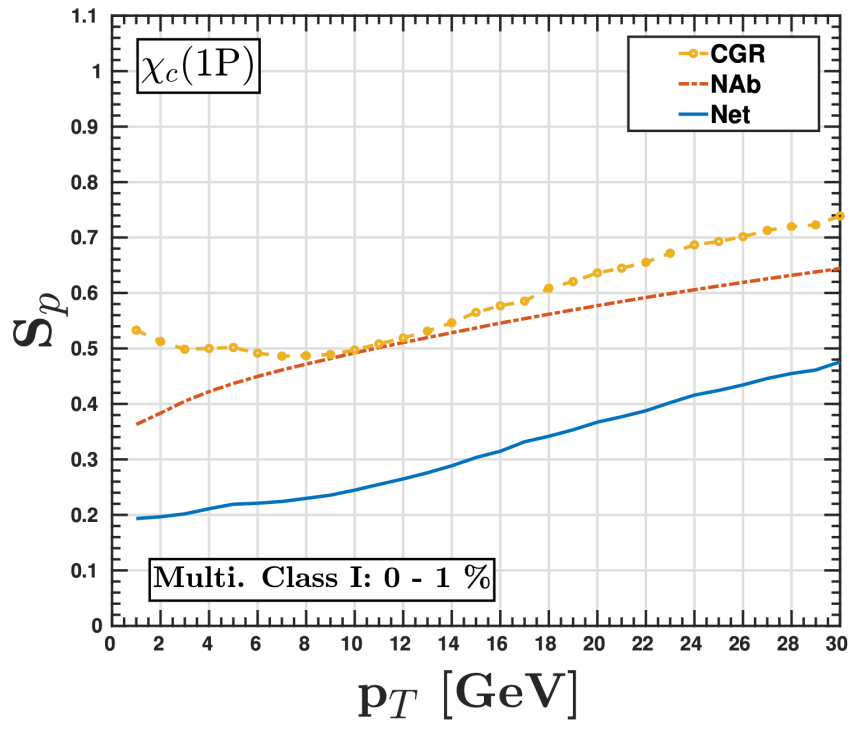


Figure 7. (Color online) Survival probability S_P as a function of p_T is shown for J/ψ considering the feed-down of $\chi_c(1P)$ and $\psi(2S)$ into J/ψ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. From top to bottom, results are shown for high-multiplicity, low-multiplicity, and minimum bias events, respectively.



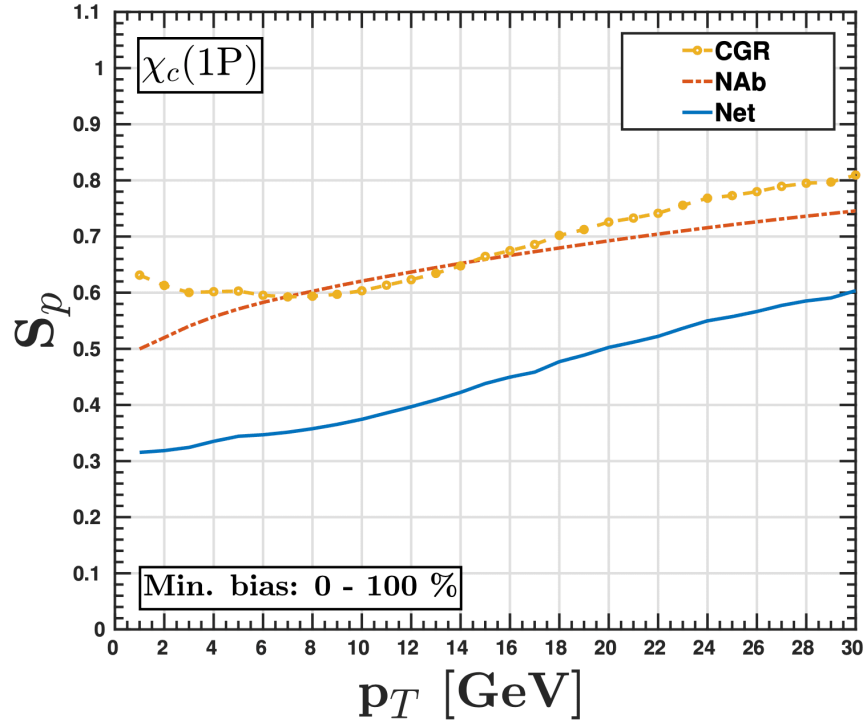
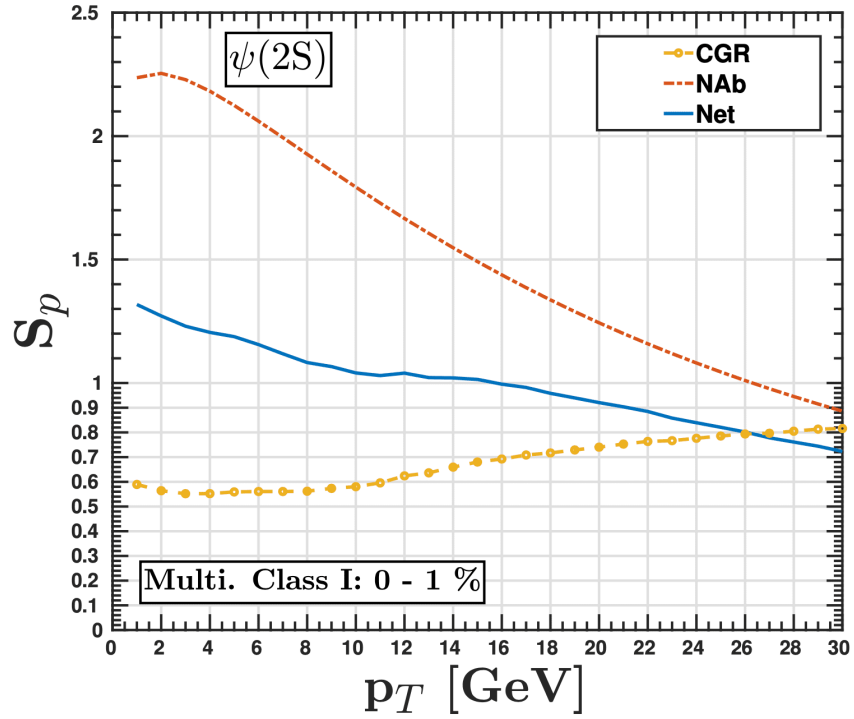
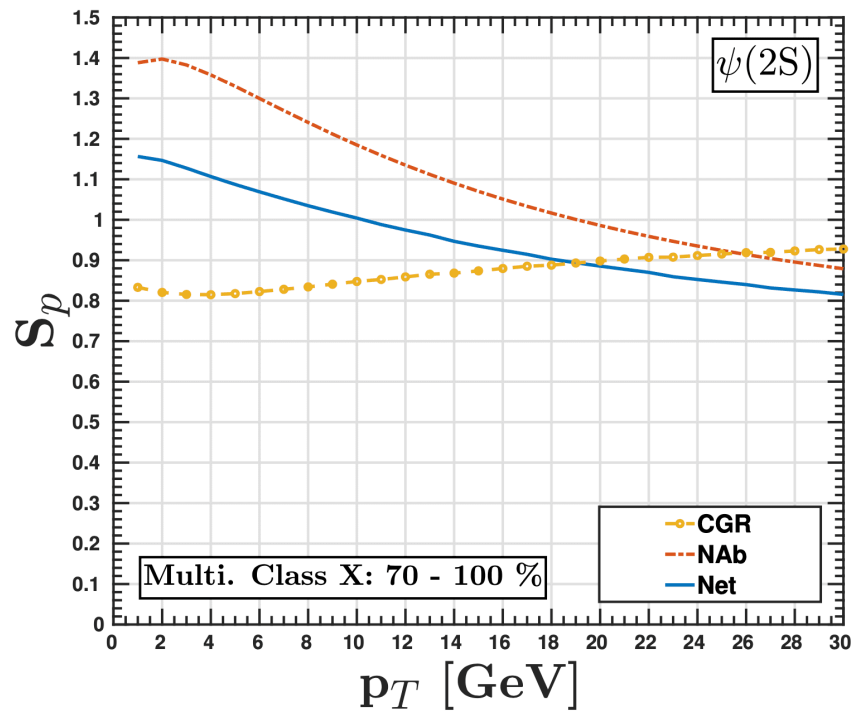


Figure 8. (Color online) Survival probability S_P as a function of p_T is shown for χ_c at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. From top to bottom, results are shown for high-multiplicity, low-multiplicity, and minimum bias events, respectively.

The p_T -dependent suppression of J/ψ corresponds to Multi. Class I, shown in Fig. 6, predicts around 20% suppression at low p_T for both NAb and CGR mechanisms. However, as p_T increases, the suppression due to NAb rapidly decreases and becomes negligible at $p_T \gtrsim 30$ GeV. On the other hand, the suppression due to CGR also decreases with increasing p_T , but at a slower rate, still predicting around 10% suppression at $p_T \simeq 30$ GeV. When these two mechanisms are combined, the suppression increases to around 40% at low p_T and 10% at high p_T , primarily due to the CGR mechanism. Fig. 7 shows that considering the feed-down corrections of $\chi_c(1P)$ and $\psi(2S)$ into J/ψ at Multi. Class I provides marginal changes in the results compared to the case without the feed-down. It suggests that high-multiplicity $\chi_c(1P)$ is largely suppressed, and as $\psi(2S)$ contribution in feed-down is relatively small, mainly J/ψ dynamics in the medium dominate the feed-down correction. The results for Multi. Class X, depicted in Fig. 6, indicates a significant decrease in suppression due to NAb, with its impact on J/ψ suppression being smaller than that of CGR at $p_T \lesssim 3$ GeV. The yield modification caused by

CGR is also reduced at the lowest multiplicity and around $p_T \simeq 30$ GeV, where its effect nearly vanishes, while NAb deactivates at $p_T \simeq 30$ GeV. When these mechanisms are combined, they predict less than 10% suppression at low p_T , which almost disappears at high $p_T \gtrsim 30$ GeV. However, the feed-down correction shown in Fig. 7 for Multi. Class X predicts a non-zero suppression at high p_T due to CGR stemming from the larger suppression of higher resonances. In contrast, the corresponding NAb effect provides a slight enhancement for J/ψ at high p_T . In the minimum bias case illustrated in Fig. 6, the prediction indicates a significant suppression of J/ψ . When both mechanisms are combined, this suppression is approximately 30% at $p_T \lesssim 3$ GeV, decreasing to 5% at higher p_T . The feed-down correction shown in Fig. 7 for the minimum bias case slightly enhances the suppression for J/ψ . In this case, the enhancement of J/ψ is found to be absent at high p_T , unlike Multi. Class X.





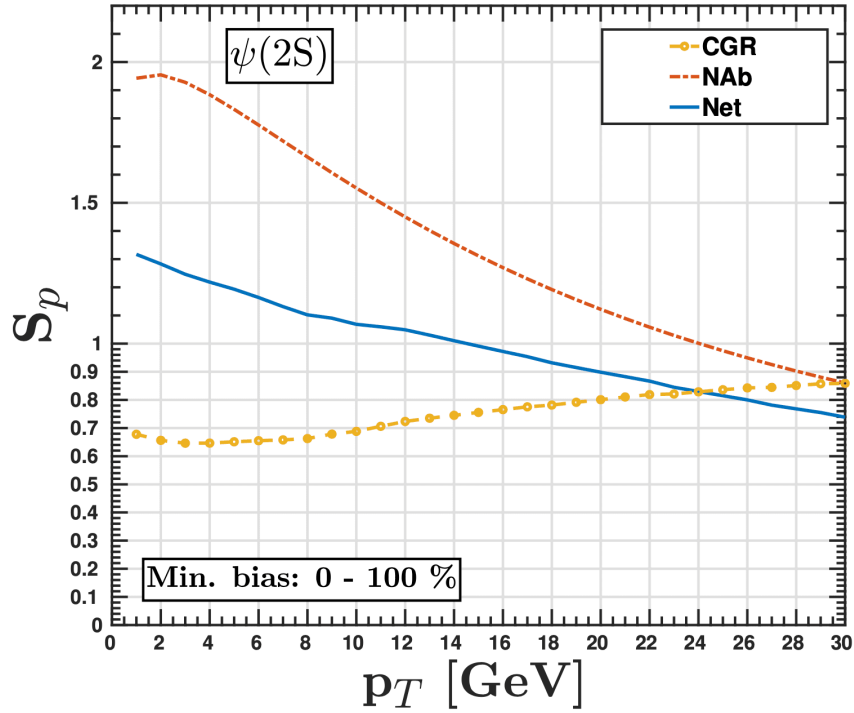


Figure 9. (Color online) Survival probability S_p as a function of p_T is shown for $\psi(2S)$ at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. From top to bottom, results are shown for high-multiplicity, low-multiplicity, and minimum bias events, respectively.

These results indicate that the non-adiabatic (NAb) evolution of the J/ψ state is predominant at low transverse momentum (p_T) and high multiplicity. At these conditions, the system size is maximized compared to lower multiplicities, and particles with low p_T moving slowly through the medium, prolong the transition from $\psi(\tau_c)$ to J/ψ at $\tau = \tau_c$. As a result, the yield of J/ψ decreases under high multiplicity and low p_T , however, this reduction diminishes rapidly with increasing p_T and a decrease in system size.

In parallel, the combined effects of collisional damping, gluonic dissociation, and regeneration (CGR) significantly influence suppression. Within CGR, collisional damping and gluonic dissociation substantially lower the yield at low p_T , whereas the regeneration mechanism tends to increase the yield at high p_T , rendering suppression less impactful for J/ψ . Conversely, for other charmonium resonances like $\chi_c(1P)$ and $\psi(2S)$, the regeneration the effect is marginal. Given their excited state nature, these resonances experience greater suppression compared to J/ψ due to CGR mechanisms.

Interestingly, NAb affects these states differently. For example, in Multi. Class I, $\chi_c(1P)$ exhibits dominant suppression due to the NAb mechanism as shown in Fig. 8. This suppression results in a yield reduction of approximately 65% at $p_T \lesssim 2$, while survival probability at high p_T increases to about 55%. In contrast, the CGR suppression for $\chi_c(1P)$ in Fig. 8 at $p_T \lesssim 2$ is around 45%, rising to 50% in the range of $2 \lesssim p_T \lesssim 12$ GeV before declining at $p_T \gtrsim 12$ GeV. The interplay of CGR and NAb predicts a suppression range of 50% to 80% for $\chi_c(1P)$ across high to low p_T in high multiplicity (Multi. Class I) events.

Further, Fig. 8 illustrates for low multiplicity events (Multi. Class X), the CGR and NAb mechanisms exhibit a complex relationship regarding $\chi_c(1P)$ suppression from low to high p_T . At $p_T \lesssim 3$ GeV, NAb is the primary suppression mechanism; however, in the range of $3 < p_T < 18$ GeV, dissociation of $\chi_c(1P)$ is largely driven by CGR processes. At high transverse momenta ($p_T \gtrsim 20$), CGR and NAb equally affect the yield of $\chi_c(1P)$ in low multiplicity events, resulting in net suppression between 20% and 40%, depending on the p_T region. Similar to Multi. Class I, NAb primarily drives suppression mechanisms for $\chi_c(1P)$ in a minimum bias scenario, except in the p_T range of 8 – 14 GeV, where CGR predicts greater suppression. The overall suppression of $\chi_c(1P)$ in minimum bias (Min. bias: 0 – 100%) lies between the extremes of multiplicity classes, ranging from approximately 70% suppression at low p_T to about 40% at high p_T .

So far, observations indicate that non-adiabatic evolution tends to reduce the yields of quarkonia, as has been predicted for J/ψ and χ_c . However, the results shown in Fig. 9 for $\psi(2S)$ is on the contrary. Instead of suppression, non-adiabatic evolution leads to a significant enhancement of $\psi(2S)$ yields across both low and high multiplicity classes. In Multi. Class I, Fig. 9 reveals a substantial enhancement of $\psi(2S)$ at low transverse momentum (p_T), which diminishes as p_T increases. At very high p_T (around 26 GeV and above), there is a noticeable suppression pattern for $\psi(2S)$. A similar phenomenon is observed in Multi. Class X, though the magnitude of enhancement is smaller compared to Class I, with the yield of $\psi(2S)$ starting to decrease at $p_T \gtrsim 10$ GeV.

These findings suggest that the non-adiabatic evolution of charmonium states facilitates the transition to excited states characterized by larger principal quantum numbers (n) and smaller azimuthal quantum numbers (l). This transition is particularly dominant when the lifetime of the medium is sufficiently long, allowing the continuum state to evolve into a discrete charmonium state. Given that $\psi(2S)$ is a higher excited state with relatively high eigen energy, it is particularly conducive to the formation during this transition from continuum to discrete eigenstates. Consequently, the

non-adiabatic mechanism predicts a significant enhancement of $\psi(2S)$ yields at high multiplicity, which then decreases due to changes in eigenenergy in lower multiplicity events.

On the other hand, the CGR mechanisms significantly reduce the yield of $\psi(2S)$ across all chosen multiplicity classes, as illustrated in Fig. 9. For Multi. Class I, the suppression is around 40–45% at $p_T \lesssim 12$ GeV, further with increasing p_T suppression reduces to 20%. The combined effects of CGR and NAb lead to a decrease in the “Net” survival probability (S_p) for $\psi(2S)$ at high p_T , which in contrast with the behavior observed for J/ψ and χ_c . A similar trend reflects the influence of NAb and CGR on the $\psi(2S)$ yield with p_T , is also seen in Multiplicity Class X; however, the magnitude of enhancement and suppression is less pronounced due to change in charged-particle multiplicity density. Under the Min. bias scenario, the $\psi(2S)$ yield falls within the range set by Multi. Class I and Multi. Class X and results are consistent with the behavior observed in other multiplicity classes.

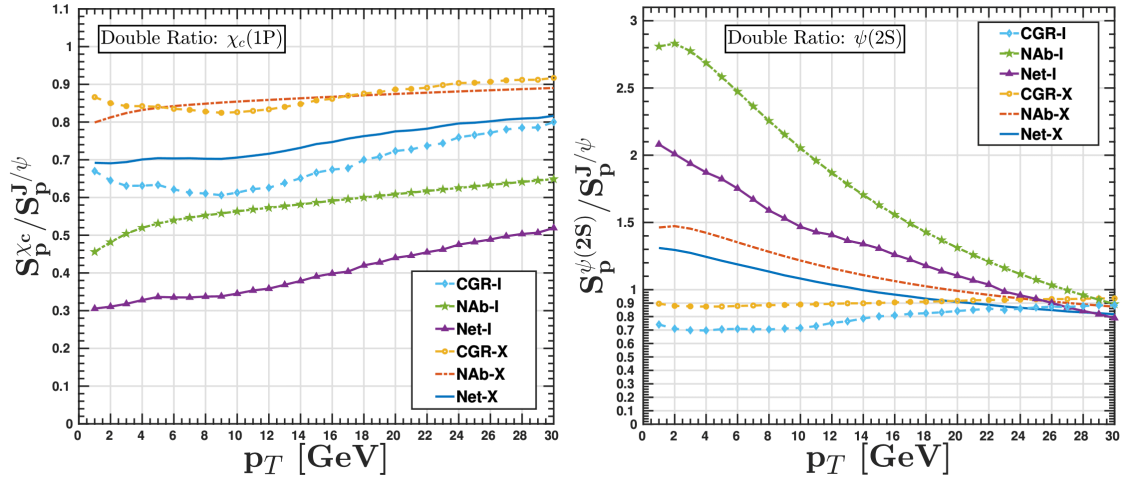


Figure 10. (Color online) Double ratio as a function of p_T is shown for $\frac{\chi_c(1P)}{J/\psi}$, and $\frac{\psi(2S)}{J/\psi}$, at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. Legends shown with “I” and “X”, stand for High and Low multiplicity events, respectively.

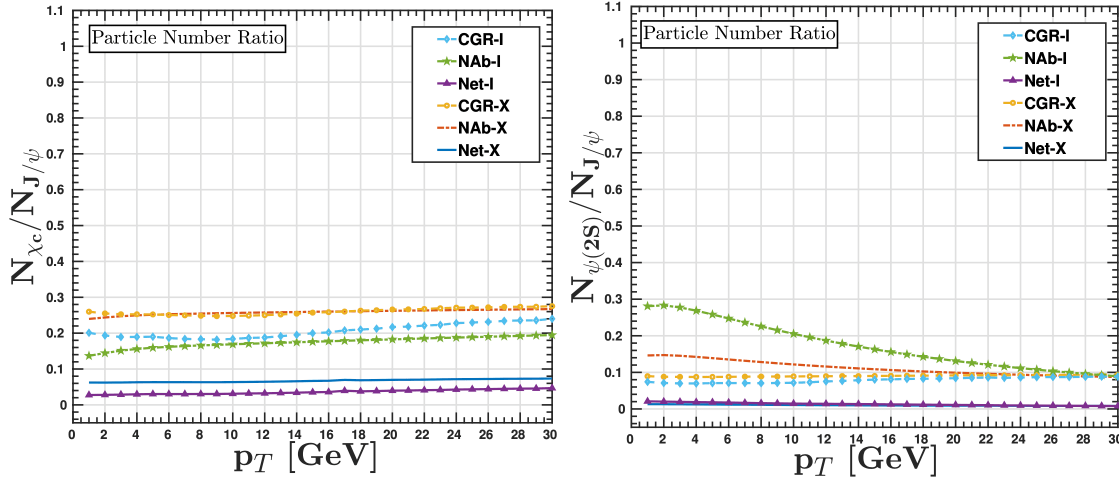


Figure 11. (Color online) Particle number ratio as a function of p_T is shown for $\frac{\chi_c(1P)}{J/\psi}$, and $\frac{\psi(2S)}{J/\psi}$, at mid-rapidity corresponding to $p + p$ collision at $\sqrt{s} = 13$ TeV. Legends shown with “I” and “X”, stand for High and Low multiplicity events, respectively.

The p_T -dependent double ratios depicted in Fig. 10 provide important insights into the relative suppression of $\chi_c(1P)$ and $\psi(2S)$ in comparison to J/ψ . At low multiplicity and p_T , the NAb mechanism demonstrates a higher degree of suppression than the CGR mechanism for $\chi_c(1P)$ relative to J/ψ . Notably, at $p_T \gtrsim 4$ GeV, the suppression for $\chi_c(1P)$ predicted by both mechanisms align closely, indicating a convergence in their outcomes under certain conditions.

In high multiplicity scenarios, the NAb mechanism is identified as a significant factor in the reduction of the $\chi_c(1P)$ yield, leading to a suppression approximately 60% to 50% greater than that of J/ψ across the selected p_T range. While the CGR mechanism predicts substantial suppression for $\chi_c(1P)$, its estimates are notably lower than the NAb mechanism. Collectively, these findings suggest a net suppression of $\chi_c(1P)$ with respect to J/ψ is around 70% to 50% at high multiplicity and from 30% to 20% at low multiplicity within the p_T range of $1 \leq p_T \leq 30$ GeV.

Moreover, the left panel of Fig. 10 illustrates that the relative yield of $\psi(2S)$ experiences a suppression of approximately 10% at low multiplicity and 30% at high multiplicity due to the CGR mechanism. Particularly, the $\psi(2S)$ considerable enhancement through the NAb mechanism underscores the distinct roles played by each mechanism in particle dynamics. The survival probability of $\psi(2S)$ in comparison with J/ψ increases at both low and high multiplicities. However, it is noteworthy that at $p_T \gtrsim 14$ GeV in low multiplicity scenarios, $\psi(2S)$ is observed to be more suppressed than J/ψ . In

higher multiplicity conditions, the onset of suppression for $\psi(2S)$ is shifted to higher p_T values, as evidenced by results indicating suppression at $p_T \gtrsim 22$ GeV. This subtle understanding enhances our comprehension of particle behavior across varying system dynamics depending on the charged-particle multiplicity density.

We have conclusively observed that the χ_c survival probability (S_p) in the QGP medium is significantly lower than that of J/ψ and $\psi(2S)$. While $\psi(2S)$ is indeed suppressed due to the CGR mechanisms, that too is less pronounced than that of χ_c . Additionally, $\psi(2S)$ experiences enhancement from the NAb mechanism. This leads us to critical questions about whether the net production of $\psi(2S)$ actually exceeds that of J/ψ or if its survival probability is merely bolstered under these conditions.

To resolve this, we have estimated the relative production numbers of χ_c and $\psi(2S)$ with respect to J/ψ , as illustrated in Fig. 11 through particle number ratios. At low multiplicity, the ratio of χ_c to J/ψ shows that approximately 30% of the final production of χ_c is conceded when factoring in CGR and NAb mechanisms independently. Notably, this final yield is nearly independent of the transverse momentum (p_T) of the particles. The combined effects of CGR and NAb definitively reduce the net χ_c yield by up to 8% at low multiplicity. At high multiplicity, the production of χ_c under CGR mechanisms accounts for around 20% to 25%, whereas predictions based solely on NAb yield estimates of 12% to 20%, fluctuating from low to high p_T . The cumulative impacts of CGR and NAb conspicuously diminish the net χ_c production by approximately 2% in high multiplicity events.

Furthermore, when examining the net yield of $\psi(2S)$ relative to J/ψ , it is found that the final yield of $\psi(2S)$ is substantially lower than that of J/ψ . At high p_T , the production levels of $\psi(2S)$ are comparable to those of χ_c , driven by the NAb mechanism. However, at low p_T , the yield of $\psi(2S)$ is significantly less than that of χ_c across both multiplicity classes. With CGR mechanisms taken into account, the yield of $\psi(2S)$ is estimated at approximately 8% to 10% relative to J/ψ for both high and low multiplicity. The combined effects of CGR and NAb indicate that the production of $\psi(2S)$ is roughly 1% of J/ψ and which is smaller than χ_c . This trend remains consistent across low and high multiplicity as well as throughout the selected p_T ranges.

These results strongly suggest that while the sequential suppression of charmonium may appear inconsistent in this context, the sequential production of charmonium states is upheld. Even when accounting for the complexities of medium dynamics and charmonium evolution in ultra-relativistic $p + p$ collisions, it is evident that $\psi(2S)$ may experience enhancements; however, the net number of J/ψ will invariably surpass that of $\psi(2S)$.

V. Summary and Outlook

This work explored the charmonium yield modification under various mechanisms that could possibly exist in ultra-relativistic proton-proton ($p + p$) collisions at $\sqrt{s} = 13$ TeV. The study considers both pre-equilibrium and thermalized QCD medium effects, modeling the temperature evolution through the bottom-up thermalization approach and Gubser flow. Under the “in-medium suppression effects” for QGP, it incorporates collisional damping, which arises because of the energy loss due to interactions between charmonium and the medium, and gluonic dissociation as the consequence of quarkonium states into a color octet lead interactions with gluons. It also includes the regeneration of charmonium states within the medium due to the transition from the color octet state to the color singlet state. Additionally, the non-adiabatic evolution of charmonium states is considered, recognizing that rapid temperature changes in small systems like $p + p$ collisions can challenge the adiabatic assumption and significantly affect the charmonium yield. At last, feed-down corrections from higher resonances into J/ψ have been incorporated for more realistic predictions.

The findings conclude that charmonium suppression is driven by these mechanisms, and their combined effect is modeled in terms of survival probabilities (S_p) as a function of transverse momentum (p_T) and charged-particle multiplicity ($dN_{ch}/d\eta$). The study finds that while J/ψ and χ_c experience significant suppression, $\psi(2S)$ shows enhancement at higher multiplicities due to non-adiabatic evolution at low p_T and high multiplicities.

- These results indicate that the QGP evolution time scale is significantly smaller than the charmonium transition time scale in ultra-relativistic $p + p$ collisions, thereby invalidating the use of the adiabatic approximation for the state evolution in the medium. This discrepancy necessitates considering a non-adiabatic evolution of charmonium, especially in small systems such as those formed in ultra-relativistic $p + p$ and even in ultra-peripheral heavy-ion collisions.
- The results suggest that J/ψ suppression and/or $\psi(2S)$ enhancement in small systems, such as $p + p$ collisions can be a valuable probe for understanding the presence of a thermalized QCD medium. This investigation suggests that ultra-relativistic $p + p$ collisions may also exhibit QGP-like behavior under specific conditions.

This study presented a holistic approach that reinforces our understanding of quark-gluon plasma characteristics and enhances our grasp of the intricate dynamics within ultra-relativistic collisions from large to small systems.

The future scope of research on charmonium yield modification in $p + p$ collisions at $\sqrt{s} = 13$ TeV offers various promising directions:

- **Non-Adiabatic Evolution:** The study shows that the evolution of quarkonia in smaller systems, such as $p + p$ collisions, may not adhere to adiabatic assumptions. Future work can further explore the non-adiabatic evolution of charmonium states, particularly in different system sizes, like peripheral heavy-ion collisions, where rapid cooling influences their behavior.
- **Comparison with Heavy-Ion Collisions:** Our findings are contrary to the $\psi(2S)$ suppression observed in heavy-ion collisions and that controversy arises because the evolution of charmonium states is considered to behave differently depending on the system size and its cooling rate. Future research should investigate this phenomenon more deeply to understand how charmonium states behave across various collision systems.
- **QGP Characteristics in Small Systems:** The findings suggest the potential for using charmonium suppression as a probe to detect thermalized QCD matter, even in small systems like $p + p$ collisions. Further experimental studies could focus on developing the methodology for such observations to probe the existence of quark-gluon plasma in such a small collision system.

These avenues can help connect theory with experimental observations, enhancing the understanding of QGP properties and charmonium dynamics in ultra-relativistic high-energy collisions.

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Declarations

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