Measurement of the $\chi_{c1}$ and $\chi_{c2}$ polarizations in proton-proton collisions at $\sqrt{s} = 8$ TeV

The CMS Collaboration

Abstract

The polarizations of promptly produced $\chi_{c1}$ and $\chi_{c2}$ mesons are studied using data collected by the CMS experiment at the LHC, in proton-proton collisions at $\sqrt{s} = 8$ TeV. The $\chi_c$ states are reconstructed via their radiative decays $\chi_c \to J/\psi \gamma$, with the photons being measured through conversions to $e^+e^-$, which allows the two states to be well resolved. The polarizations are measured in the helicity frame, through the analysis of the $\chi_{c2}$ to $\chi_{c1}$ yield ratio as a function of the polar or azimuthal angle of the positive muon emitted in the $J/\psi \to \mu^+\mu^-$ decay, in three ranges of $J/\psi$ transverse momentum. While no differences are seen between the two states in terms of azimuthal decay angle distributions, they are observed to have significantly different polar anisotropies. The measurement favors a scenario where at least one of the two states is strongly polarized along the helicity quantization axis, in agreement with nonrelativistic quantum chromodynamics predictions. This is the first measurement of significantly polarized quarkonia produced at high transverse momentum.

Submitted to Physical Review Letters
Quarkonium production is a benchmark for understanding how quarks combine into hadrons. The relative heaviness of c and b quarks makes it possible to describe the production process in the nonrelativistic quantum chromodynamics (NRQCD) factorization approach \cite{1-8}, a rigorous framework valid when the quark velocities are small. This theory successfully described quarkonium production cross sections measured \cite{9} at high transverse momentum, $p_T$, by complementing the earlier color-singlet model \cite{10,11} with subprocesses where the bound state originates from colored $Q\bar{Q}$ pairs, which lose their color via gluon emissions, changing the angular momentum quantum numbers $L$, $S$, and $J$.

In contrast to this complex picture of quarkonium production as a superposition of several color-singlet and multi-step color-octet processes, ATLAS and CMS measurements \cite{12-21} reveal unexpectedly simple patterns \cite{22-24}. First, the five S-wave ($L = 0$, $S = 1$, $J = 1$) charmonium and bottomonium states ($J/\psi$, $\psi(2S)$, $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$) are produced with indistinguishable mass-rescaled $p_T$ spectra ($p_T/M$). This observation, indicating a universal trend independent of the bound-state excitation level and of the constituent quark mass, is also followed by the $\chi_{c1}$ and $\chi_{c2}$ P-wave ($L = 1$, $S = 1$) states, at the current level of experimental precision \cite{17-19}. Second, the measurements of the polarizations of the five S-wave quarkonia do not show indications, over a wide $p_T$ range, of any deviation from unpolarized production, where all angular momentum projections $J_z$ are equally probable. This is a peculiar condition for vector particles, not seen for the Z and W bosons \cite{25-32}, nor for Drell–Yan dileptons \cite{33-38}, which are always significantly polarized, as are the quarkonia produced at low $p_T$ (i.e., $p_T < M$) \cite{39,40}. The lack of polarization of vector quarkonia was a long-standing challenge for NRQCD \cite{41}, but recent global-fit analyses \cite{4,22-24,42} have shown that cross sections and polarizations can be consistently described. This progress revealed a distinct set of fine-tuned long-distance parameters \cite{43} and triggered another strong prediction, this time regarding the remaining gap in the set of LHC quarkonium cross section and polarization measurements mentioned above: the polarizations of the $\chi_{c1}$ and $\chi_{c2}$ states should be opposite and almost maximal \cite{44}.

This Letter reports the first measurement of the polarizations of promptly produced $\chi_{c1}$ and $\chi_{c2}$ mesons, using proton-proton (pp) data collected at the LHC by the CMS experiment at a center-of-mass energy of $\sqrt{s} = 8$ TeV, corresponding to an integrated luminosity of 19.1 $fb^{-1}$. The central feature of the CMS apparatus is a superconducting solenoid of 6 m internal diameter, providing a magnetic field of 3.8 T. Within the solenoid volume are a silicon pixel and strip tracker, a lead tungstate crystal electromagnetic calorimeter, and a brass and scintillator hadron calorimeter, each composed of a barrel and two endcap sections. Forward calorimeters extend the pseudorapidity ($\eta$) coverage provided by the barrel and endcap detectors. Muons are detected in gas-ionization chambers embedded in the steel flux-return yoke outside the solenoid. A detailed description of the CMS detector, together with a definition of the coordinate system used and relevant kinematic variables, can be found in Ref. \cite{45}.

The event sample was collected with a two-level trigger system \cite{46}. At level-1, custom hardware processors select events with two muons. The high-level trigger requires an opposite-sign muon pair of invariant mass 2.8–3.35 GeV, a dimuon vertex fit $\chi^2$ probability larger than 0.5%, and a distance of closest approach between the two muons smaller than 0.5 cm. The trigger also requires that the dimuon has $p_T > 7.9$ GeV and rapidity $|y| < 1.25$. The offline reconstruction requires two oppositely charged muons matching those that triggered the detector readout. The muon tracks must pass high-purity track quality requirements \cite{47}, have $p_T > 3.5$ GeV, $|\eta| < 1.6$, and fulfill the soft muon identification requirements \cite{48}, which imply, in particular, more than five hits in the silicon tracker, of which at least one is in the pixel layers. The muons are combined to form $J/\psi$ candidates, which are kept for further processing if $|y| < 1.2$.
and $8 < p_T < 30$ GeV. Promptly produced $J/\psi$ mesons are selected by requiring the distance between the dimuon vertex and the interaction point be smaller than 2.5 times its uncertainty.

The analysis uses $\chi_c \rightarrow J/\psi \gamma$ decays, with the $J/\psi$ decaying to a dimuon. The photons are detected through their conversions to $e^+e^-$ in the beam pipe and in the material of the silicon tracker, starting from two oppositely charged tracks, of which one has at least four tracker hits and the other at least three. The tracks must have a small angular separation, a small distance of closest approach, a conversion vertex at least 1.5 cm away from the beam axis, and a $\chi^2$ probability of the kinematic fit imposing zero mass and a common vertex that exceeds 0.05%.

A more detailed account of the reconstruction and selection procedures is given in Refs. [18],[19]. The photons must have $p_T > 0.4$ GeV and $|\eta| < 1.5$. If the distance along the beam axis between the dimuon vertex and the extrapolated photon trajectory is smaller than 5 mm, a $\chi_c$ candidate is formed through a kinematic fit of the dimuon-photon system, constraining the dimuon mass to the $J/\psi$ mass [19], the dielectron mass to zero, and requiring that the two muons and the photon have a common vertex. Only $\chi_c$ candidates with a fit $\chi^2$ probability larger than 1% and invariant mass between 3.2 and 3.75 GeV are kept for the evaluation of the $\chi_{c1}$ and $\chi_{c2}$ yields. After all selection criteria, the event samples used to perform the analysis, in the $J/\psi$ $p_T$ ranges 8–12, 12–18, and 18–30 GeV, contain around 103 000, 106 000, and 45 000 $\chi_c$ candidates, respectively.

The seemingly natural way to measure the $\chi_{c1}$ and $\chi_{c2}$ Polarizations is to determine the angular distribution of the considered $\chi_c$ decay; in the present case, this means the distribution of the photon direction in the $\chi_c$ rest frame. However, that distribution depends not only on the $\chi_c$ angular momentum composition, but also, and possibly in a very significant way, on the (poorly known) contributions of photons with large orbital angular momentum ($J^T > 1$). A cleaner determination of the $\chi_c$ polarization is obtained by measuring the dilepton angular decay distribution in the rest frame of the daughter $J/\psi$ [50]. It is crucial to choose as polarization axis for the $J/\psi$ decay not the $J/\psi$ direction in the $\chi_c$ rest frame, as usually done in cascade decays, but rather any axis (center-of-mass helicity or Collins–Soper [51], for instance) defined in terms of the beam momenta in the $J/\psi$ rest frame and ignoring its origin, as if it were observed inclusively. Contrary to what happens when measuring the photon distribution in the $\chi_c$ rest frame, the latter choice leads to measurements that are insensitive to the higher-order multipoles and reflect only the polarization state of the mother $\chi_c$. The present analysis is performed in the center-of-mass helicity frame [52] and does not use the measured photon momentum, except to select, through the $J/\psi \gamma$ invariant mass distribution, the $J/\psi$ mesons resulting from $\chi_{c1}$ or $\chi_{c2}$ decays. The dimuon angular decay distribution is parametrized with the function [53]

$$W(\cos \theta, \varphi | \bar{\lambda}) = \frac{3}{4\pi(3 + \lambda_\theta)} \left( 1 + \lambda_\theta \cos^2 \theta + \lambda_\varphi \sin^2 \theta \cos 2\varphi + \lambda_{\theta\varphi} \sin 2\theta \cos \varphi \right),$$

(1)

where $\theta$ and $\varphi$ are the polar and azimuthal coordinates of the positive lepton direction in the $J/\psi$ rest frame, the system of axes being defined with $z$ in the direction of the polarization axis and $y$ perpendicular to the production plane. The $\chi_c$ angular momentum composition is encoded in the shape parameters $\lambda_\theta$, $\lambda_\varphi$, and $\lambda_{\theta\varphi}$, whose values depend on the choice of polarization frame but must always be within certain physical domains [50], narrower than the parameter space of inclusive vector-particle production [54]. The relation between the shape parameters and the polarization configuration depends on the quarkonium state. For example, $\lambda_\theta = +1$ indicates $I_z = \pm 1$ for the $J/\psi$, $I_z = 0$ for the $\chi_{c1}$, and $I_z = +2$ for the $\chi_{c2}$; conversely, states in the $I_z = 0$ angular momentum configuration lead to $\lambda_\theta = -1$ for the $J/\psi$, $\lambda_\theta = +1$ for the $\chi_{c1}$, and $\lambda_\theta = -0.6$ for the $\chi_{c2}$. Simulation studies have shown that all physically allowed differences between the $\chi_{c1}$ and $\chi_{c2}$ Polarizations can be reliably determined through the measurement of
the $\chi_{c2}/\chi_{c1}$ yield ratio as a function of $|\cos \theta|$ or $\phi$. Detector effects (trigger, data reconstruction, event selection) cancel to a large extent in the ratio and have a negligible impact on the results.

The analysis is independently performed in three $J/\psi$ $p_T$ ranges: 8–12, 12–18, and 18–30 GeV. In each range, the events are split into subsamples corresponding to six equidistant $\phi$ bins between 0 and 90°. Folding $\phi$ into the first quadrant reduces the effect of the statistical fluctuations without any loss of information, given the four-fold $\phi$ symmetry that the angular distributions obey. For each $p_T$ bin, the six $J/\psi \gamma$ invariant mass distributions are simultaneously fitted with an unbinned maximum likelihood fit. In the mass fit model, identical for all $\phi$ bins, each of the $\chi_{c1}$ and $\chi_{c2}$ signal peaks is represented by a double-sided Crystal Ball (CB) function [55], which complements a Gaussian core distribution with lower and upper power-law tails. The underlying combinatorial background, reflecting uncorrelated $J/\psi \gamma$ associations, is parametrized by an exponential function multiplied by a term that provides a low-mass turn-down, $(1 + \text{erf}((m - \mu_{bg}) / \sigma_{bg})) \exp(-m/\lambda_{bg})$, where $m$ is the $J/\psi \gamma$ invariant mass and $\mu_{bg}$, $\sigma_{bg}$, and $\lambda_{bg}$ are shape parameters. Although the results of this analysis are insensitive to the presence of a small peak reflecting the $\chi_{c0}$ decays, the fit model includes this background term, represented by a Breit–Wigner convolved with a Gaussian resolution function. To minimize fit instabilities, the $\chi_{c0}$ shape and yield parameters are determined from the corresponding parameters of the $\chi_{c1}$ term. The simultaneous fit has the advantage of reducing by a factor of six the number of free parameters defining the shapes of the signal and background mass models, by requiring that those parameters are independent of $\phi$, an assumption validated by studies of simulated and measured event samples. The parameters of interest resulting from this fit are the six $\chi_{c2}/\chi_{c1}$ yield ratios.

The analysis is also performed as a function of $|\cos \theta|$, splitting the events in 6, 7, or 5 $|\cos \theta|$ bins, depending on the $p_T$ range. The $|\cos \theta|$ coverage is smaller in the lowest $p_T$ range (up to 0.45 instead of up to 0.625) because those events are the ones most affected by the single-muon $p_T$ cut. Analogously to the procedure just described for the $\phi$ dimension, the $\chi_{c2}/\chi_{c1}$ yield ratios are obtained as a function of $|\cos \theta|$ through a simultaneous fit of the $J/\psi \gamma$ invariant mass distributions. In this case, however, some of the shape parameters (such as the mass resolutions) show a slight correlation with $|\cos \theta|$, so that they are not required to be independent of $|\cos \theta|$. While some parameters are unconstrained, others are required to depend linearly on $|\cos \theta|$, as an intermediate way of minimizing the number of free parameters.

Figure [1] shows one of the simultaneously fitted $J/\psi \gamma$ invariant mass distributions. The two signal peaks are well resolved, with widths around 6 MeV, consistent with the predictions from simulation. All of the fitted $\chi_c$ mass distributions show good fit qualities, as judged from the $\chi^2$ between the binned distributions and the fitted functions, the worst case giving $\chi^2 = 601$ for 569 degrees of freedom (ndf).

The $\chi_{c2}/\chi_{c1}$ yield ratios provided by the fits of the $\chi_c$ mass distributions are corrected for the slightly different acceptances and efficiencies for the detection of the two states, using fine-grained acceptance times efficiency three-dimensional maps, $A(|\cos \theta|, \phi, p_T)$, computed with large samples of simulated events. The corrected ratios are reported in Tables A.1 and A.2 of Appendix [A] and shown in Fig. [2], where it can be seen that the uncorrected and corrected values are almost identical, apart from normalization factors irrelevant for the determination of the polar and azimuthal anisotropies.

Several sources of potential systematic effects have been considered, by redoing the analysis with different inputs and comparing the obtained results with the nominal ones. The results are insensitive to variations of the thresholds used to reject the nonprompt contamination from $b$ hadron decays, estimated to be around 5%, or events with a poor kinematic vertex fit quality.
in the reconstruction of the $\chi_c$ candidates. The fits of the mass distributions were redone using alternative options for the low- and high-mass tails of the double-sided CB functions, and by varying the combinatorial background description, both by changing the floating parameters of the nominal function and by using the alternative function $(x-x_0)^\lambda \exp(\nu(x-x_0))$, where $\nu$ is left free, $\lambda$ is fitted to a constant, and $x_0 = 3.2$ GeV, a value determined in fits to the background-only mass distributions obtained by excluding the 3.37–3.6 GeV region. The sensitivity of the results to the acceptance and efficiency corrections was evaluated by redoing the analysis with maps computed with alternative single-muon and photon detection efficiencies, as well as with simulated samples generated with different $p_T/M$ shapes for each of the two $\chi_c$ states. All effects lead to similar variations in the yields of the two states and cancel, to a large extent, in the $\chi_{c2}/\chi_{c1}$ ratio, apart from a normalization shift that has no impact on the angular anisotropies.

The $\chi_{c2}$ to $\chi_{c1}$ yield ratios as a function of $\varphi$, shown in Fig.2 (left), are compatible with being flat, excluding large differences in azimuthal anisotropy, as exemplified by the two curves compared to the data points in the second $p_T$ range. These curves represent the simplest conceivable polarization hypotheses leading to large azimuthal effects in the helicity frame: $\chi_{c1}$ and $\chi_{c2}$ have maximally different polar anisotropies in the Collins–Soper frame, corresponding to specific alignments of their angular momentum vectors along the collision direction ($f_{x_{c1}}^z = f_{x_{c2}}^z = 0$ and $f_{x_{c1}}^\perp = \pm 1$, $f_{x_{c2}}^\perp = \pm 2$, for the dotted and dash-dotted curve, respectively). In fact, the change from the Collins–Soper to the helicity quantization axis is almost a $90^\circ$ rotation, transforming polarized distributions into azimuthally anisotropic ones. This uniform $\varphi$ behavior confirms the choice of the helicity axis as the one that, as expected in this kinematic regime, should reflect most closely the natural alignment of the angular momentum vector, maximizing the polar anisotropy effects.

In Fig.2 (right) the measured $|\cos \theta|$ dependence of the $\chi_{c2}/\chi_{c1}$ ratio is compared to the analytic expression $(1 + \lambda_\varphi^{x_{c2}} \cos^2 \theta) / (1 + \lambda_\varphi^{x_{c1}} \cos^2 \theta)$, derived from Eq.1 integrating over $\varphi$. Two scenarios are considered. The “unpolarized scenario”, $\lambda_\varphi^{x_{c1}} = \lambda_\varphi^{x_{c2}} = 0$ independently of $p_T$, represented in Fig.2 (right) by the dashed flat lines, gives a poor description of the data. A fit with free normalizations leads to a $\chi^2/\text{ndf} = 31/15$, corresponding to a $\chi^2$ probability of only 0.9%. The “NRQCD scenario” [44], where $\lambda_\varphi^{x_{c1}} = 0.72$, 0.65, and 0.56, and $\lambda_\varphi^{x_{c2}} = -0.48$, -0.35, and -0.19, for the average $p_T$ values in each of the three ranges, agrees well with the data, with a fit $\chi^2/\text{ndf} = 13/15$, corresponding to a $\chi^2$ probability of 58%.
Figure 2: The $\chi_{c2}/\chi_{c1}$ yield ratio vs. $\varphi$ (left) and $|\cos \theta|$ (right), in the helicity frame, for the three $J/\psi$ $p_T$ ranges. The grey markers (slightly shifted horizontally) show the values before acceptance and efficiency corrections, scaled vertically for an easier shape comparison. The vertical bars represent the statistical uncertainties and the horizontal bars the bin widths. The solid and dashed curves show, respectively, the “NRQCD” and “unpolarized” scenarios. The dotted and dash-dotted curves illustrate maximally different natural polarizations in the Collins–Soper frame, leading to large differences in azimuthal anisotropy.

Figure 3: Two-dimensional $\lambda^X_{\varphi}$ vs. $\lambda^X_{\chi_{c1}}$ contours, at 68.3, 95.5, and 99.7% confidence levels (CL), measured combining the three $J/\psi$ $p_T$ ranges. The physically allowed region (red rectangle) and six pure angular momentum configurations (markers) are also shown. The crossing of the two dashed lines represents the unpolarized case.

Figure 3 shows the polar anisotropy parameters $\lambda^X_{\chi_{c1}}$ and $\lambda^X_{\chi_{c2}}$ derived from the measured $|\cos \theta|$ dependence of the $\chi_{c2}/\chi_{c1}$ ratio, combining the three $p_T$ ranges. The contours in the $\lambda^X_{\varphi}$ vs. $\lambda^X_{\chi_{c2}}$ plane are obtained by scanning the two $\lambda_{\varphi}$ parameters and the three normalizations to evaluate the $\chi^2$ profiles corresponding to the 68.3, 95.5, and 99.7% confidence levels. The unpolarized scenario ($\lambda^X_{\varphi} = \lambda^X_{\chi_{c2}} = 0$), as well as more than half of the physically allowed region,
The \( \lambda_{c2} \) values (circles) measured when the \( \lambda_{c1} \) values (squares) are fixed to the unpolarized (left) or the NRQCD (right) scenarios, as a function of \( p_T/M \) of the J/\( \psi \). The purple band on the right is the NRQCD prediction for \( \lambda_{c2} \) [44], while in the unpolarized scenario \( \lambda_{c2} = \lambda_{c1} = 0 \). The markers are shown at the average \( p_T/M \) values in each bin, the vertical bars represent the total uncertainties, and the horizontal bars the bin widths. The dashed lines indicate the physically allowed range of \( \lambda_{c2} \).

including all cases where \( \lambda_{c2} \geq \lambda_{c1} \), are outside the 99.7% contour. In terms of specific pure angular momentum configurations, it can be seen that, in particular, the cases \( J_{c2} = \pm 2 \) and \( J_{c1} = J_{c2} = \pm 1 \) are strongly disfavored.

The correlation between the \( \lambda_{c1} \) and \( \lambda_{c2} \) parameters can be accurately expressed through a simple parametrization:

\[
\lambda_{c2} = (-0.94 \pm 0.90 \lambda_{c1}^{\chi_{c1}}) \pm (0.51 \pm 0.05 \lambda_{c1}^{\chi_{c1}}),
\]

\[
(-0.76 \pm 0.80 \lambda_{c1}^{\chi_{c1}}) \pm (0.26 \pm 0.05 \lambda_{c1}^{\chi_{c1}}),
\]

\[
(-0.78 \pm 0.77 \lambda_{c1}^{\chi_{c1}}) \pm (0.26 \pm 0.06 \lambda_{c1}^{\chi_{c1}}),
\]

for the three consecutive \( p_T \) ranges. These expressions can be used for direct comparisons to theoretical scenarios.

Figure 4 shows, as a function of \( p_T/M \) of the J/\( \psi \) (equal on average to the \( p_T/M \) of the \( \chi_{c1} \) and \( \chi_{c2} \) mothers [23]), the \( \lambda_{c2} \) values measured when \( \lambda_{c1}^{\chi_{c1}} \) is fixed to the predictions of the two scenarios already considered in Fig. 2. Setting \( \lambda_{c1}^{\chi_{c1}} = 0 \) leads to \( \lambda_{c2} \) values that are significantly different from zero (and even tend to be outside the physically allowed range). The NRQCD prediction is, instead, in good agreement with the measurement.

In summary, the polarizations of promptly produced \( \chi_{c1} \) and \( \chi_{c2} \) mesons have been measured in pp collisions at \( \sqrt{s} = 8 \) TeV. The analysis uses the J/\( \psi \gamma \) decay channel in three J/\( \psi \) \( p_T \) ranges between 8 and 30 GeV. The measurement, made in the helicity frame, shows a significant difference between the polar anisotropy parameters \( \lambda_{c1}^{\chi_{c1}} \) and \( \lambda_{c2}^{\chi_{c1}} \). The result strongly disfavors, in particular, the unpolarized scenario uniformly observed in the J/\( \psi \), \( \psi(2S) \), and \( \Upsilon \) measurements. Remarkably, the measurement agrees with the NRQCD prediction. This result provides a new piece in the experimental scenario of quarkonium production at mid-rapidity and the first significant indication of kinematic differences between the various quarkonia. It should improve the understanding of hadron formation and of the interplay between the long- and short-distance aspects of the strong interaction.
Acknowledgments

We congratulate our colleagues in the CERN accelerator departments for the excellent performance of the LHC and thank the technical and administrative staffs at CERN and at other CMS institutes for their contributions to the success of the CMS effort. In addition, we gratefully acknowledge the computing centers and personnel of the Worldwide LHC Computing Grid for delivering so effectively the computing infrastructure essential to our analyses. Finally, we acknowledge the enduring support for the construction and operation of the LHC and the CMS detector provided by the following funding agencies: BMBWF and FWF (Austria); FNRS and FWO (Belgium); CNPq, CAPES, FAPERJ, FAPERGS, and FAPESP (Brazil); MES (Bulgaria); CERN; CAS, MoST, and NSFC (China); COLCIENCIAS (Colombia); MSES and CSF (Croatia); RPF (Cyprus); SENESCYT (Ecuador); Academy of Finland, MEC, and HIP (Finland); CEA and CNRS/IN2P3 (France); BMBF, DFG, and HGF (Germany); GSRT (Greece); NKFIA (Hungary); DAE and DST (India); IPM (Iran); SFI (Ireland); INFN (Italy); MSIP and NRF (Republic of Korea); MES (Latvia); LAS (Lithuania); MOE and UM (Malaysia); BUAP, CINVESTAV, CONACYT, LNS, SEP, and UASLP-FAI (Mexico); MESTD (Montenegro); MBIE (New Zealand); PAEC (Pakistan); MSHE and NSC (Poland); FCT (Portugal); JINR (Dubna); MON, RosAtom, RAS, RFBR, and NRC KI (Russia); MESTD (Serbia); SEIDI, CPAN, PCTI, and FEDER (Spain); MOSTR (Sri Lanka); Swiss Funding Agencies (Switzerland); MST (Taipei); ThEPCenter, IPST, STAR, and NSTDA (Thailand); TUBITAK and TAEK (Turkey); NASU (Ukraine); STFC (United Kingdom); DOE and NSF (USA).

References


[55] M. J. Oreglia, “A study of the reactions $\psi' \rightarrow \gamma\gamma\psi$” PhD thesis, Stanford University, 1980. SLAC Report SLAC-R-236, see Appendix D.
A Numerical values of the measured yield ratios

Table A.1: The ratio of the $\chi_{c2}$ to $\chi_{c1}$ yields, corrected for acceptance and efficiencies, vs. $\phi$, in three $J/\psi$ $p_T$ ranges. The average $\phi$ values are also given.

<table>
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<tr>
<th>$J/\psi$ $p_T$ (GeV)</th>
<th>$\phi$ (degrees)</th>
<th>$\langle \phi \rangle$ (degrees)</th>
<th>$\chi_{c2}/\chi_{c1}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>8–12</td>
<td>0–15</td>
<td>7.8</td>
<td>0.451$^{+0.027}_{-0.025}$</td>
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<td></td>
<td>15–30</td>
<td>22.6</td>
<td>0.452$^{+0.026}_{-0.025}$</td>
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<td></td>
<td>30–45</td>
<td>37.6</td>
<td>0.499$^{+0.027}_{-0.026}$</td>
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<td>45–60</td>
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<td>0.472$^{+0.025}_{-0.024}$</td>
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<td>60–75</td>
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<td>0.450$^{+0.023}_{-0.022}$</td>
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<tr>
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<td>75–90</td>
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<td>0.445$^{+0.023}_{-0.022}$</td>
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<td></td>
<td>15–30</td>
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<td>0.393$^{+0.018}_{-0.017}$</td>
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<td>30–45</td>
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<td>0.412$^{+0.019}_{-0.018}$</td>
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<td>0.425$^{+0.030}_{-0.028}$</td>
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<td>75–90</td>
<td>82.5</td>
<td>0.409$^{+0.028}_{-0.027}$</td>
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Table A.2: The ratio of the $\chi_{c2}$ to $\chi_{c1}$ yields, corrected for acceptance and efficiencies, vs. $|\cos \theta|$, in three $J/\psi$ $p_T$ ranges. The average $|\cos \theta|$ values are also given.

| $J/\psi$ $p_T$ (GeV) | $|\cos \theta|$ | $\langle |\cos \theta| \rangle$ | $\chi_{c2}/\chi_{c1}$ |
|-----------------------|----------------|-----------------|-----------------|
| 8–12                  | 0.000–0.075    | 0.037           | $0.453^{+0.018}_{-0.018}$ |
|                       | 0.075–0.150    | 0.111           | $0.463^{+0.021}_{-0.020}$ |
|                       | 0.150–0.225    | 0.185           | $0.489^{+0.025}_{-0.024}$ |
|                       | 0.225–0.300    | 0.259           | $0.439^{+0.024}_{-0.025}$ |
|                       | 0.300–0.375    | 0.332           | $0.388^{+0.035}_{-0.031}$ |
|                       | 0.375–0.450    | 0.404           | $0.411^{+0.056}_{-0.054}$ |
| 12–18                 | 0.000–0.075    | 0.038           | $0.476^{+0.023}_{-0.021}$ |
|                       | 0.075–0.150    | 0.113           | $0.438^{+0.020}_{-0.019}$ |
|                       | 0.150–0.225    | 0.187           | $0.421^{+0.020}_{-0.019}$ |
|                       | 0.225–0.300    | 0.262           | $0.397^{+0.021}_{-0.019}$ |
|                       | 0.300–0.375    | 0.336           | $0.398^{+0.022}_{-0.021}$ |
|                       | 0.375–0.450    | 0.409           | $0.376^{+0.026}_{-0.024}$ |
|                       | 0.450–0.625    | 0.502           | $0.392^{+0.033}_{-0.032}$ |
| 18–30                 | 0.000–0.150    | 0.076           | $0.445^{+0.036}_{-0.032}$ |
|                       | 0.150–0.300    | 0.225           | $0.455^{+0.030}_{-0.027}$ |
|                       | 0.300–0.375    | 0.338           | $0.463^{+0.039}_{-0.036}$ |
|                       | 0.375–0.450    | 0.412           | $0.365^{+0.032}_{-0.030}$ |
|                       | 0.450–0.625    | 0.526           | $0.370^{+0.027}_{-0.025}$ |
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