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Blurred femtoscopy in two-proton decay

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ABSTRACT

We study the effects of final state interactions in two-proton emission by nuclei. Our approach is based on the solution the time-dependent Schrödinger equation. We show that the final relative energy between the protons is substantially influenced by the final state interactions. We also show that alternative correlation functions can be constructed showing large sensitivity to the spin of the diproton system.

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Two-proton emission has been observed for numerous excited states in nuclei, populated both in β decay and in nuclear reactions [1–6]. Although these decays are thought to be sequential one-proton emissions proceeding through states in the intermediate nucleus [7,8], there is an intriguing possibility [9] that the diproton (^{2}He) correlation may play an important role in the mechanism of the two-proton emission. This has been nicely demonstrated in the analysis of the two-proton decay of the 21^+ isomeric state in ^{94}Ag, where 19 events were clearly assigned to the simultaneous emission of two correlated protons [4,6]. The traditional idea of diproton radioactivity is due to the pairing effect. Two protons form a quasiparticle (diproton) under the Coulomb barrier and this facilitates penetration. In a more formal description, one has a system with two valence protons in the same shell and coupled to $J^\pi = 0^+$. This question, being still open, continues to motivate studies in this field. In order to assess this information, it is necessary to understand final state interactions between the protons and between each proton and the daughter nucleus.

In nuclear decays the emission of correlated, identical, particles is sensitive to the geometry of the system. Measurements of correlation functions are often performed with charged particle pairs, which interact via the short-range nuclear interaction and the long-range Coulomb interaction and they also interact with the remaining source. As a result, theoretical corrections are needed to subtract the final state interactions (FSI) before one can extract any useful information about the emitting source from the measurements [10–13]. At first sight, the FSI can be regarded as a contamination of "pure" particle correlations. However, it should be noted that the FSI depend on the structure of the emitting source and thus provide information about source dynamics as well.

Two-proton decay in s-wave states can also be used for testing quantum mechanics versus local realism by means of Bell's inequalities [14]. Since the final state of the two protons can be either in a singlet or in a triplet state, their wavefunction is spin entangled. The identification of the spins of the proton in two detectors separated far away would be useful to test the Einstein–Podolski–Rosen (EPR) paradox [15,16]. In fact, these tests should be performed in different and complementary branches of physics to avoid the loopholes encountered in photon experiments. The advantage of using massive fermions to test Bell-type inequalities is that the particles are well localized and the spin state of the pair can be well established by measuring the internal energy of the two-proton system. However, the validity of this method highly depends on our ability to treat FSI. Coincidence measurements of the two proton momenta require knowledge of FSI in order to extract information about their original wavefunction. Here will propose a new method to calculate FSI based on the numerical solution of the time-dependent Schrödinger equation. We hope with that to get a quantitative estimate of the FSI and how they can be used to address the points raised above. Though tests of Bell's inequality using proton–proton spin correlation in low energy scattering con-

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formed to quantum mechanics (see, e.g., [17]), one would still like to find other means to perform the verification of the complete nature of quantum mechanics.

One should distinguish this work from Hanbury-Brown–Twiss (HBT) studies in high energy nucleus–nucleus collisions. Indeed, in the case of HBT, the whole game is played by FSI. FSI are the means by which one can determine information about the source. In this case FSI is not viewed as a "contamination". The important difference in our case is that we are interested in a case where the protons are not emitted chaotically like in the case of HIC (Heavy-Ion Collisions). We are studying nuclear structure. In HIC (and HBT) protons are assumed to be emitted independently and chaotically from the source (any information about their initial spins is lost and they are assumed to be "evaporated" with a boiling pot). Then, there are no initial state correlations and FSI make the whole physics. In the case of two-proton radioactivity, the emission of the two protons is not chaotic because their correlation function keeps memory of their spin admixture and wave function in the parent nucleus.

We consider first a single proton described at the initial time by a localized wave-packet $\psi_0(\mathbf{r}_1)$. The probability amplitude to find the proton at the detector with momentum $\mathbf{p}_1$ is given by

$$\mathcal{A}(\mathbf{p}_1,\mathbf{r}_1)=\int d\mathbf{r}\,\chi^{(+)}(\mathbf{p}_1,\mathbf{r})K(\mathbf{r},\mathbf{r}_1)\psi_0(\mathbf{r}_1), \tag{1}$$

where $\chi^{(+)}(\mathbf{p}_1,\mathbf{r})$ is an asymptotic outgoing Coulomb wave with energy $E=\mathbf{p}_1^2/2m_p$, and $K(\mathbf{r},\mathbf{r}_1)$ is the propagator which accounts for the time evolution of the particle from the source to the detector.

We now look at the case of two-protons interacting with the residual nucleus and between themselves. We will consider the distortion caused by the Coulomb plus nuclear interaction between each proton i with the nucleus, $V_C(\mathbf{r}_i)+V_N(\mathbf{r}_i)$, and between themselves, $v_C^{12}(\mathbf{r})+v_N^{12}(\mathbf{r})$, where $\mathbf{r}_i$ is the coordinate of proton i, and $\mathbf{r}$ is their relative coordinate. The proton–nucleus interaction, $V_N(\mathbf{r}_i)$ yields smaller final state interaction effects than the Coulomb counterpart.

We adopt a classical description of the center-of-mass motion for the two-protons and solve the time-dependent Schrödinger equation for the relative motion between them. The Coulomb field that distorts the relative motion of the particles is given by

$$V_C(t)=Ze^2\left(\frac{1}{|\mathbf{r}_1-\mathbf{R}(t)|}-\frac{1}{|\mathbf{r}_2-\mathbf{R}(t)|}-\frac{2}{R(t)}\right), \tag{2}$$

where Z is the charge of the daughter nucleus and $\mathbf{r}_1$ and $\mathbf{r}_2$ are the positions of the protons with respect to the center of the nucleus (nuclear recoil is neglected). $V_C(t)$ acts on the relative position $\mathbf{r}=\mathbf{r}_2-\mathbf{r}_1$ through the transformations $\mathbf{r}_1=\mathbf{R}-\mathbf{r}/2$ and $\mathbf{r}_2=\mathbf{R}+\mathbf{r}/2$.

One can perform a multipole expansion of this interaction and for r smaller than $R(t)$ one can express the result in terms of a multipole-dependent effective charge, $e_L=e[(-1/2)^L+(1/2)^L]$ where L is the multipole degree. The dipole field ($L=1$) is only important for particles with different charge-to-mass ratios while the quadrupole field is dominant when these ratios are equal (e.g. for two-proton emission). For the quadrupole interaction, $e_{L=2}=e/2$ and

$$V_C(t)=\frac{Ze^2}{2}\frac{r^2}{R^3(t)}P_2(\cos\theta), \tag{3}$$

where θ is the angle between $\mathbf{R}$ and $\mathbf{r}$ and P_2 is the Legendre polynomial of order 2. We now assume that the protons are produced simultaneously and nearly at rest at position $2a_0$ and time $t=0$. Their center-of-mass follows a radial trajectory described by

$$R(t)=\frac{a_0}{2}(\cosh w+1),\qquad t=\frac{a_0}{2v}(\sinh w+w), \tag{4}$$

where the asymptotic velocity is given by $v=\sqrt{E/m_p}$, E is the two-proton decay energy, and $a_0=e^2/2E$. This assumes that the relative energy between the protons is much smaller than E, which is not a good approximation, as we will show later. It is important to notice that Eqs. (4) only account for the motion of the protons after they emerge from inside the nucleus through the Coulomb barrier and propagate from the closest distance $2a_0$ to infinity. Hence, our calculations neglect what happens during the tunneling process and treat only the external motion. Hence, neglecting the proton–nucleus strong FSI is justified.

We can still use Eq. (1) to calculate the probabilities for relative motion of the protons, with the wavefunction for the relative motion given by $\Psi(\mathbf{r})=K(\mathbf{r},\mathbf{r}_0)\psi_0(\mathbf{r}_0)$. In the time-dependent description, at time t this wave function can be expanded in spherical harmonics

$$\Psi(\mathbf{r})=\frac{1}{r}\sum_{lm}u_{lm}(r,t)Y_{lm}(\hat{\mathbf{r}}), \tag{5}$$

and the Schrödinger equation, describing the time evolution of the relative motion between the protons can be solved by the finite difference method, calculating the wavefunction at time $t+\Delta t$ in terms of the wavefunction at time t, according to the algorithm

$$u_{lm}(t+\Delta t)=\left[\frac{1}{i\tau}-\Delta^{(2)}+\frac{\Delta t}{2\hbar\tau}U\right]^{-1}\times\left[\frac{1}{i\tau}+\Delta^{(2)}+\frac{\Delta t}{2\hbar\tau}U+\frac{\Delta t}{\hbar\tau}S_{l'm';lm}\right]u_{lm}(t), \tag{6}$$

where $\tau=\hbar\Delta t/m_p(\Delta r)^2$. The second difference operator is defined as

$$\Delta^{(2)}u_{lm}^{(j)}(t)=u_{lm}^{(j+1)}(t)+u_{lm}^{(j-1)}(t)-2u_{lm}^{(j)}(t), \tag{7}$$

with $u_{lm}^{(j)}(t)=u_{lm}(r_j,t)$, where r_j is a position in the radial lattice. In Eq. (6), $U=v_C^{12}(r_j)+v_N^{12}(r_j)$ is the Coulomb + nuclear interaction between the two protons as a function of their distance, r_j, and the function $S_{l'm';lm}$ is given by

$$S_{l'm';lm}(r,t)=\sum_{l'm'}\langle Y_{l'm'}|V_C(r,t)|Y_{lm}\rangle u_{l'm'}(r,t). \tag{8}$$

This method of solving the time-dependent equation is the same as used in Ref. [18] for studying reacceleration effects in breakup reactions in nucleus–nucleus collisions at intermediate energies. A grid adequate for our purposes has 5000 spatial mesh points separated by 0.1 fm and 2000 time mesh points separated by 0.5 fm/c.

We use the quantization-axis along the $\mathbf{R}(t)$ center-of-mass radial trajectory. As a consequence, $P_2(\cos\theta)=\sqrt{4\pi/5}Y_{20}(\theta,\phi)$, and one only needs to consider the $m=0$ component of the spherical harmonics implicitly contained in the potential $V_C(t)$. The initial $l=0$ state cannot develop a final $l=1$ component, and only $l=0$ (s-waves) and $l=2$ (d-waves) will be present in the final state. Higher l values will be small and need not be considered.

The proton–proton potential is taken as $v_N^{12}(r)+v_C^{12}(r)=e^2/r+v_0(b/r)\exp(-r/b)$. The set of parameters $v_0=-46.124$ MeV and $b=1.1809$ fm yields the proton–proton scattering length, $a_p=-7.8196$ fm and the effective range $\rho_0=2.790$ fm, in accordance with experimental data. But we choose a higher absolute value of v_0 which allows the presence of a single weakly bound s-wave state. We use this localized wavefunction for the relative motion of the two protons in the initial state: $u_0\equiv u_{l=0}(r,t=0)$. This is an artifact of the numerical method chosen as to allow for a localization of the initial wavefunction. The observables associated with the final state will depend on the binding energy, reflecting the dependence on the initial average separation between them.

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